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HIGH DENSITY PLASMA EFFECTS ON ATOMIC AND IONIC SPECTRA

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London SW7 2BZ, U.K.*Résumé

La compression d'un plasma par laser produit des champs électriques et magnétiques d'une intensité exceptionnelle. Les perturbations des spectres sont telles qu'il reste peu de niveaux liés, même pour des émetteurs multiples ionisés avec $z > 10$ - et la physique même du plasma devient complexe à cause d'interactions fortes entre les espèces qui le composent ou encore à cause d'effets de dégénérescence.

La physique de base des plasmas ultra-denses se distingue de celle qu'on suppose couramment dans la spectroscopie des plasmas, tant en astrophysique qu'en laboratoire. Nous passons en revue les problèmes théoriques que pose l'interaction émetteur - plasma dans un milieu dense et nous donnons quelques exemples d'intérêt expérimental et théorique actuel.

Abstract

Laser compressed plasmas are sources of exceptionally intense electric and magnetic fields. Consequent perturbations of spectra are such that even for high ionization stage emitters with $z > 10$ only a few bound states may remain - and the physics of the plasma itself is complicated by the strong interactions amongst the constituent species and by near-degeneracy.

The basic physics of such ultra-dense plasmas is contrasted with that usually assumed in longer established areas of laboratory and astrophysical plasma spectroscopy, theoretical problems in treating emitter-plasma interactions in dense plasmas are surveyed, and some examples are given of problems of current experimental and theoretical interest.

1. Introduction

This paper discusses what are probably the most severe examples of the modification of atomic and ionic spectra by strong fields yet realized in the laboratory. Laser compressed plasmas presently under laboratory study reach densities of hot, highly-ionized matter measurable in grams or tens of grams cm^{-3} , i.e. higher than solid, with electron densities of interest being in the range 10^{24} - 10^{25}cm^{-3} . The corresponding plasma microfield due to the inter-particle electric fields is therefore such that even for highly-ionized emitters (with, say, $z > 10$) only the very lowest bound states retain any real identity with those in isolated ions. Since in such plasmas the particles are strongly interacting (i.e. compared with thermal energies), the calculation of the statistical properties of the microfield, allowing not only for its instantaneous magnitude but its fluctuations (which are often more crucial than instantaneous field strength in determining effects on emitted spectra), is a formidable statistical-mechanical problem, even before the problems of the response of an atom/ion to the field, or the reaction of the emitter state back on the microfield are approached.

Complex as it is, this 'strongly-perturbed' spectroscopy of these very dense

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plasmas is of much more than just academic interest. Laser-compressed plasmas are presently the most frequently studied example of the so-called Inertial Confinement (ICF) route to fusion power. Unlike the situation in the longer-established Magnetic Confinement programme, spectroscopy plays a fundamental role in both the diagnostic study and the underlying physics of ICF plasmas. For a laser plasma at an electron density, N_e of, say, 10^{24} cm⁻³, the electron plasma frequency, ω_{pe} , already lies in the soft X-ray spectral region. Consequently, laser-based diagnostics such as Thomson scattering, well established in magnetic confinement work, are impossible, and the prime quantitative diagnostic technique is X-ray emission spectroscopy (together, in some cases, with absorption work based on back-lighting from separately generated heavy-element laser plasma X-ray sources). Also, radiation transport may play a significant role in the physics of the plasma as a whole, for instance in the compression dynamics via preheating of material ahead of the converging shock front, and the radiation transport in turn is controlled by processes such as pressure broadening, series merging and ionization potential depression. In both contexts it is crucial to develop an understanding of the strong perturbations of atomic and spectral properties that occur, even though the complex nature of the medium means that theoretical models are necessarily less elegant and more approximate than those used in other problems involving 'strong fields'.

Interestingly, but of less practical importance, laser-generated plasmas are also the sources of probably the highest magnetic fields yet generated in the laboratory. Scale lengths for electron density and temperature gradients in these plasmas are so short that very intense currents develop and local magnetic fields in the megagauss range are generated. Whilst spectroscopic effects of these B fields are usually masked by the combined effects of both Doppler and pressure broadening, some possible indirect consequences will be mentioned below.

In this paper the basic physical situation in a laser compressed plasma will first be contrasted with that in more conventional magnetically contained and astrophysical plasmas. The observable effects of the plasma environment will be surveyed, and a review will be given of some rather general problems encountered in attempts to develop a sound theoretical base for the spectroscopy of these ultra-dense sources.

FUSION PLASMAS

	<u>Magnetic Containment</u> (Tokamaks etc.)	<u>Inertial Confinement</u> (Laser Compression etc.)
T_e	$10^7 - 10^8$ K	$10^6 - 10^7$ K
n_e	$10^{12} - 10^{14}$ cm ⁻³	$10^{23} - 10^{26}$ cm ⁻³
Duration	0.1 - 10 s.	10^{-10} s.
Size	1 - 10 m	$10^{-5} - 10^{-6}$ m
Ion stages	$X - XXX$	$X - XX$
Kinetic pressure	1 bar	$1 - 10^3$ Mbar (Fusion 10^6 Mbar)
Statistics	Debye ($n\lambda_D^3 \gg 1$) Classical ($n\lambda_D^3 < 1$)	Strongly-coupled ($n\lambda_D^3 \approx 1$) Near degenerate ($n\lambda_F^3 \approx 1$)
Ionization	Non-LTE	Saha fails
Atomic spectra	Unperturbed	Strongly perturbed ($na_0^3 \gg 1$)
Line Shapes	Doppler	Pressure
Radn. Transport	$\tau \ll 1$ (except in far IR)	Optically thick, even in continuum.

TABLE 1: Comparison of physical parameters in low-density, magnetically and high-density, inertially confined, fusion plasmas.

2. Comparison with the spectroscopy of other laboratory plasma sources.

The recent advent of laser-compressed plasmas represents perhaps the most rapid and extreme change in the entire history of laboratory plasma spectroscopy. For the first time the physics of the laboratory sources of interest differs in fundamental rather than merely quantitative ways from the theoretical basis assumed in most of the essentially astrophysical background out of which the subject grew. The nature of this physical change can be seen by contrasting a laser compressed plasma with that generated in a magnetically-contained fusion device such as a Tokamak, Table 1. (For a more general survey of other present-day plasma sources, and their spectroscopy, see the references by Burgess, 1, and by Peacock and Burgess, 2).

Table 1 shows that laser plasmas and Tokamaks are somewhat similar in their electron temperature, T_e , and in that both are capable of generating ions in high ionization states. However, at any deeper level the two sources differ so greatly as to almost represent different states of matter.

The first point to consider before discussing spectroscopic properties is the general statistical mechanical state of the plasma. A Tokamak typifies the type of plasma classically encountered in 'traditional' laboratory and astrophysical spectroscopy in that it is a high temperature, low density system and is both classical statistically (in the sense that it is very far from Fermi degeneracy) and is an ideal gas (i.e. $p = nkT$) since the mean interparticle interaction energies are much lower than kT . It is therefore 'weakly-interacting' in the sense that:-

$$\frac{e^2}{R_0} < kT_e$$

where $R_0 = \left\{ \frac{3}{4\pi N} \right\}^{1/3}$ is the mean interparticle separation,

$$\text{i.e. } r = \frac{e^2}{R_0 kT_e} = 2.69 \times 10^{-3} \frac{N^{1/3}}{T_e} \ll 1 \quad (1)$$

(with T in Kelvin)

Equivalently, (apart from a numerical constant) the plasma is 'Debye' in the sense that the Debye shielding length is much greater than the interparticle separation), i.e.:-

$$a = R_0 / \lambda_D = 0.19 N_e^{1/3} T_e^{-1/2} < 1 \quad (2)$$

This simply means that there is no basic statistical mechanical problem in working out the properties of such a plasma as they pertain to its spectroscopy. The 'macroscopic' behaviour of a Tokamak plasma may be complex, e.g. in its interaction with magnetic fields, the existence of collective plasma waves, etc., but the frequencies of such plasma modes are much lower than those involved in most spectroscopic processes, the impact parameters of interest in excitation/de-excitation events are much smaller than mean interparticle radii, and as in the classical theory of astrophysical spectroscopy, the spectroscopic properties of the plasma can be defined entirely by a knowledge of a single one-particle parameter, namely the electron velocity distribution, plus a few atomic properties known a-priori - energy levels, binary collision cross-sections and transition probabilities.

In contrast, a laser-compressed plasma with a kinetic pressure of perhaps 10^8 atmospheres (and for laser fusion 10^{12}), in present experiments the radiation pressure due to the laser light alone may reach 10^6 atmospheres) may be near degenerate (for ICF it will have to be close to the Fermi adiabat) and it will be 'strongly-coupled' or 'non-Debye', i.e. :-

$$r \text{ (or } a) > 1 \quad (3)$$

In consequence it is a 'non-ideal' gas, the particle velocity and spatial distributions depend on the mutual interactions of the plasma constituents (and near a z-times ionized emitter on the local emitter-plasma interaction, see Ref.1), and specification of the spectroscopic state of the medium intrinsically requires a treatment of the correlations between its constituent particles. It is also important to note that $\Gamma > 1$ automatically implies that there will be a strong coupling between a z-times ionized emitter and its local plasma environment, so that the plasma properties relevant to the determination of the emitted spectrum are not necessarily those typical of the plasma as a whole, but instead are those of the strongly-interacting local polarization cloud specific to the particular emitter. Thus the 'perturbation' due to the plasma cannot be decoupled from the emitter's properties but will certainly be dependent on the emitter's charge and perhaps also on its atomic state. (As pointed out e.g. in Ref.1, for large z this problem can be important even when $\Gamma < 1$).

3. Mean plasma fields in relation to nuclear Coulomb fields

It is now relevant to consider typical values of the microfield, E_p , in a laser compressed plasma in relation to the nuclear field seen by bound electrons, E_N . As a first approximation we consider the field to be strong if:-

$$E_p / E_N > 0.1 \quad (4)$$

It is important to notice that spectroscopic effects may be observable at much smaller relative values of E , since e.g. for closely spaced levels strong mixing, and hence the appearance of field-induced forbidden-lines, can occur at relatively low values of E . It is also crucial to note that a frequent situation in a plasma is that the fluctuation properties of the field are more important than the effects of the instantaneous field itself, the simplest (second-order) picture being that the high frequency components of the Fourier decomposition of the field, $E(k, \omega)$, induce off-resonant transitions exactly similar to 2-quantum transitions in the laser physics case. If the field fluctuates rapidly enough, line broadening will be dominated by the frequency spread of the fluctuations rather than any direct Stark shifts due to E itself.

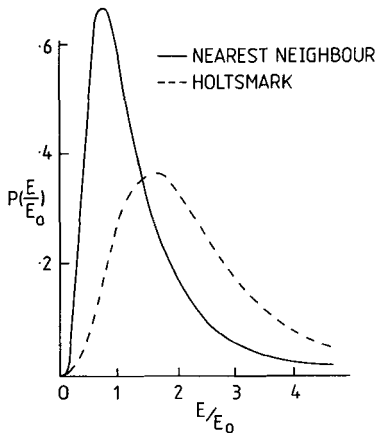


FIG1 THE MICROFIELD DISTRIBUTION $P(E)$ VERSUS NORMALIZED FIELD STRENGTH FOR NEAREST NEIGHBOUR — AND HOLTSMARK ---- THEORY

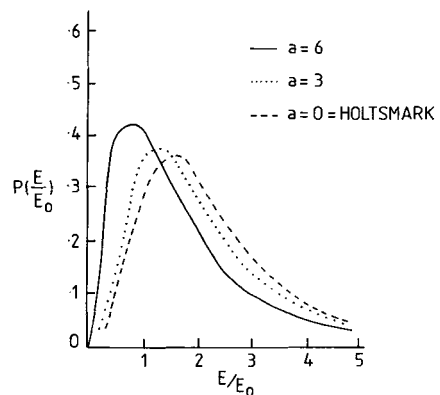


FIG2. THE MICROFIELD DISTRIBUTION VERSUS NORMALIZED FIELD STRENGTH FOR VARIOUS VALUES OF $a = r_0/\lambda_D$

With these reservations, however, it is still useful to consider the mean static field in a plasma of given density. Fig. 1 shows the predictions for the probability distribution $P(E)$ of the two simplest models in which no account is taken of particle correlations, namely the nearest-neighbour approximation and the vector sum of the particle fields due originally to Holtsmark (see e.g. Griem, Ref.3). Fig.2, based on the work of Hooper (4), shows changes in the distribution due to perturber correlations as the parameter a , Eqn.(2), increases.

Weisheit and Pollock have shown computationally that for very large the field distribution is very close to the simple nearest neighbour model. In every case, therefore, the most probable field lies between the Holtsmark and single nearest neighbour values, and even in a very dense plasma a reasonable estimate of the mean field for present purposes is simply

$$E_p \approx E_o = e/R_o^2 = 2.60 e N_e^{\frac{2}{3}} \quad (5)$$

With also the value for E experienced by an orbiting electron of principal quantum number n in a hydrogen-like ion of charge z ,

$$E_N = \frac{Z^3 e^2}{n^4 a_o^2} \quad (6)$$

we obtain:-

$$R_E = E_p / E_N = 7.2 \times 10^{-17} N_e^{\frac{2}{3}} n^4 Z^3 \quad (7)$$

The value of n for $R_E = 0.1$ is plotted against N_e in Fig.3 for various values of z . The contrast between a Tokamak and a laser compressed plasma in terms of the effect of the plasma on atomic spectra is then immediately evident. For an ion of $z=10$ in a Tokamak, $R_E = 0.1$ only occurs for $n > 100$, states which will be quite unobservable since the Doppler broadening will have caused lines to merge for much smaller n . However, for a laser plasma at 10^{24} cm^{-3} , the equivalent value of n is $n=3$ - in other words the entire energy level structure of a hydrogenic ion with $z=10$ will certainly be strongly changed by the surrounding plasma environment, and at most only one or two excited states will be observable as distinct entities.

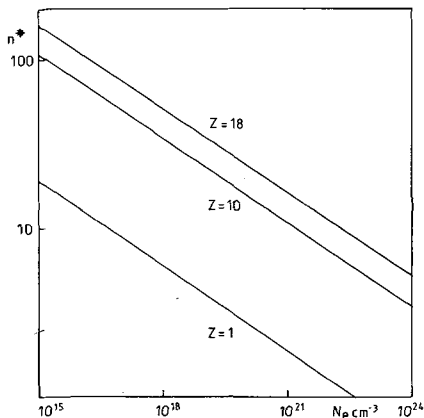


FIG.3. PLOT OF PRINCIPLE QUANTUM NUMBER n^* AT WHICH THE HOLTSMARK MEAN FIELD IS 10% OF THE NUCLEAR COULOMB FIELD VERSUS ELECTRON DENSITY N_e AND EMITTER CHARGE Z

4. Magnetic Fields

In present day Tokamak plasmas, the confining magnetic fields typically lie between 1-10 Tesla, and thus rate as rather moderate fields by the standards of

this Conference. From the practical spectroscopic viewpoint, the total magnitude of the field is actually not of too much interest, since the field is primarily externally applied, and the only part one would want to measure is that generated by the plasma current, typically between 1-10% of the total. As will be shown below, in practical terms the Zeeman effect of the overall field is normally already swamped by Doppler broadening. Although given sensitive enough polarization dependent techniques one could envisage a measurement of the absolute field magnitude it would normally be hard indeed to determine the plasma induced component and hence the pitch angle from Zeeman splitting measurements. Attempts to devise means of measuring the pitch angle of the total field (an important quantity in stability considerations) have therefore concentrated on laser scattering techniques; see the discussion in reference (2).

Laser-generated plasmas on the other hand are sources of very intense magnetic fields indeed spontaneously generating fields in the range of 100-1000 Tesla (1-10 Megagauss), Ref.(11). These fields arise from the very short scale lengths for density and temperature gradients in such plasmas, the most important case perhaps being when electron density and temperature gradients are orthogonal. The resultant " $\nabla n \cdot \nabla T$ " magnetic fields thus essentially arise accidentally, crossed density gradients of this type usually resulting from lack of uniformity in the laser irradiation of the target.

A rough estimate of the size of the B fields generated from density gradients can be obtained from:-

$$B(\text{Tesla}) = 0.9 T_e^{0.5} L^{-1} (A/Z) \quad (8)$$

where T_e is the electron temperature in Kelvin, A and Z are the atomic weight and the mean charge state of the target material respectively and L is the scale length for the gradients measured in microns. With L typically ten microns or less and the ratio of A to Z being roughly unity, values of B reaching one thousand Tesla are easily attained.

Spectroscopically, one then first has to compare the Zeeman effects of the fields in the two types of plasma with those of Doppler and pressure broadening, and secondly to enquire whether any further effects are caused, e.g. on the mechanisms underlying pressure broadening.

Comparison of Zeeman splitting and Doppler broadening is straightforward and with T in Kelvin yields:-

$$\frac{\Delta\nu_Z}{\Delta\nu_D} = 0.67 \times 10^{-2} \lambda(\text{\AA})^{\frac{1}{2}} (A/T) B(\text{Tesla}) \quad (9)$$

A rough estimate of the Zeeman splitting to the linear Stark effect in the Holtzmark mean field similarly yields

$$\frac{\Delta\nu_Z}{\Delta\nu_S} = 8.10^9 B(\text{Tesla})(Z/n)N^{\frac{2}{3}} \quad (10)$$

where n is the principal quantum number of the state of interest, and N is the electron density in cm^{-3} .

It then follows that in a Tokamak, the Zeeman and Stark effects are roughly comparable; but both are normally entirely masked by the Doppler broadening which typically will be about one thousand (or more) times larger for the soft X-ray transitions (100 \AA and below) of prime interest.

In a laser-generated plasma on the other hand very intense magnetic fields will be generated in the surface regions near the critical density layer (where the laser frequency is equal to the electron plasma frequency), i.e. at electron densities typically of the order of 10^{21} - 10^{22}cm^{-3} . With values of B of 100-1000 Tesla, Doppler broadening and Zeeman splitting may well be comparable even for soft X-ray transitions. On the other hand both will normally be exceeded by pressure ('Stark') effects on the line width.

Direct (Zeeman splitting) effects of magnetic fields are therefore often masked by other effects in both Tokamaks and laser plasmas, despite the very large values of B that may be reached in laser plasmas. It should however be noted that

in lower density, lower temperature plasmas simultaneous magnetic field and pressure effects on spectral line shapes may be observable both as modifications of line shapes as a whole, and in causing polarization effects. These have been considered for hydrogen Balmer-alpha by Drawin and his co-workers, Ref.(9). Similarly, in the sub-critical density layer of laser plasmas ($N_e < 10^{21} \text{ cm}^{-3}$) magnetic field effects on spectra of highly ionized emitters may be important and the combination of these with pressure effects has been considered in some detail by Nguyen-Hoe, 10.

There may, also, be a number of subsidiary effects of magnetic fields which also need consideration. These include the curvature of perturber orbits relevant to pressure broadening calculations, see Ref.10, the fact that in very strong magnetic fields the electron velocity distribution as a whole may become strongly anisotropic, and finally that the effects of collective plasma waves on observed line shapes (see below) may be modified by the presence of electron cyclotron and hybrid modes.

5. Observable effects on spectra in dense plasmas

The rather simple considerations already given establish the difference between laser compressed plasmas and other sources in two important respects. Firstly the extent to which the plasma perturbs the properties of bound atomic levels is greater in a laser compressed plasma than perhaps in any other source, even when the emitter charge is high. Secondly, the strongly-interacting (non-Debye) nature of the medium poses statistical mechanical problems which are of themselves of great interest (these latter properties can, however, be simulated in lower density, lower temperature sources given the n_e and T_e dependence of Γ).

Although the present paper is not primarily devoted to experimental problems it is worth noting the differences, Table 1, between laser plasmas and Tokamaks in this respect also. Whilst a Tokamak is rather uninteresting as a source of strong field effects, it is nevertheless close to being a spectroscopist's dream source in practical terms in that the plasma is very long-lived, large, and quasi-stationary (in respect e.g. to properties such as ionization equilibrium, etc.). In contrast a laser-compressed plasma is something of an experimenter's nightmare and for this practical reason comparison of experiment and theory is still at a relatively very early stage. The plasma is only microns rather than metres in size, and lasts for picoseconds rather than the seconds characteristic of modern large Tokamaks, so that transient effects on ionization balance etc. need consideration. Only the extreme brightness of these very small plasmas makes them observable at all!

Qualitatively, the observable changes in the spectra from laser compressed plasmas correspond in general terms to those studied for many years in low density plasma spectroscopy, the difference being firstly that these effects are observable for the first time for high z ions, secondly the severity of the perturbation, and finally the statistical mechanical problems of the strongly-interacting medium. In general spectroscopic terms, the effects which need consideration are:-

- (1) Line-broadening/shift
- (2) Field-induced 'forbidden' ($\Delta L=0, \pm 2$) transitions
- (3) 'Plasmon' satellites and other plasma wave induced features
- (4) Ionization Potential Depression (as distinct from series merging)
- (5) Doppler narrowing.

These effects will be returned to after a brief survey of the present situation in regard to calculation of the plasma microfield and the emitter response thereto.

6. General Comments on Calculation of the Plasma Microfield and Emitter Response.

There are two basic levels at which the emitter-perturber coupling can be considered.

One, appropriate at moderate densities and typical of classical line-broadening theories (see e.g. Griem (Ref.3)), can be loosely called 'perturbative' in that the plasma is regarded as the source of a fluctuating perturbing field, which can be calculated independently of any knowledge of the atomic state of the emitter. In more formal language the system density matrix can be factored into a plasma ('bath') part and a term describing the emitter. Within this basic approach it is possible to account for some emitter-perturber correlations, for instance in the well-known effects on the static part of the microfield distribution if the emitter is charged (which can be rather accurately handled, see Refs.4, 5, 6 and 7), or even in approximate attempts to treat screening effects on the bound electron-nuclear interaction due to penetration of plasma particles, usually referred to as 'polarization shifts' (see e.g. Refs.1 and 3 for general references and Skupsky, Ref.8, for a specific theoretical treatment). This type of theory of line broadening and related phenomena in moderate density plasmas has been intensively developed - and in low temperature/low density sources experimentally checked - over many years, and for this reason has formed the basis of most theoretical predictions yet used in diagnostic studies on very dense, hot plasmas. However, the extent to which such predictions for high z -emitters in hot, dense matter have yet been checked experimentally in situations in which densities and temperature are known independently a-priori is still extremely limited.

At the highest plasma densities of current interest such a 'perturbative' approach clearly fails for two separate reasons. Firstly, even if the statistics of the plasma remain classical the basic assumption that the plasma 'perturbation' can be treated as independent of the details of the emitter state fails if the emitter-plasma coupling becomes strong and comparable to thermal energies. Secondly, as the plasma becomes part degenerate, and as interaction distances amongst the plasma particles become comparable to atomic radii (which they clearly do at $n > 10^{24} \text{cm}^{-3}$), the effects of the Pauli principle become important both on the 'perturber-perturber' interactions and because the specific occupancy of 'bound' states in a given emitter may strongly affect the statistics of the 'free' ('perturbing') particles. In such a regime the only way forward is to attempt a self-consistent 'first principles' solution of the physics of the system of nucleus and 'bound' and 'free' electrons as a whole. Such treatments, see below, do give some useful information on properties such as mean ionization state (relevant e.g. to conductivity calculations), effective ionization potential depression, and even some (mean) opacity information, but at present cannot provide the detailed spectroscopic information on e.g. shapes/shifts of specific lines needed for diagnostic purposes.

7. 'Moderate' density plasmas

7.1 General comments

Provided the 'system' density matrix can be factored into 'emitter' and 'plasma' parts and $\Gamma < 1$ (and even in some cases if $\Gamma > 1$), then it is a relatively tractable plasma physics problem to calculate either the probability distribution of the instantaneous electric field, $P(E)$, or the Fourier components of the field as a whole, $E(k, \omega)$. Unfortunately, this still leaves us a long way from being able to calculate a practical quantity such say as a line-shape - or even perhaps from making contact with the rest of this conference where mostly 'strong' fields are either static or have just one or two well defined Fourier components! The problem in the plasma case is that calculation of the atomic response requires fundamentally differing approaches for the different frequency components of $E(\omega)$, combined with the fact that the atomic response is 'non-linear' in the sense that the various frequency components cannot be treated independently. The plasma case corresponds to a field which is (a) 'strong', (b) has a broad frequency spectrum and (c) has superimposed strong peaks at well-defined frequencies, corresponding to the underlying collective plasma modes. The problems (which in the plasma case have been encountered for 20 years and

more) are therefore identical to those in attempting to treat the response of atoms to strong multi-mode laser fields in which the modes are not necessarily randomly phased. For the purposes of this Conference it seems worth illustrating the consequent difficulties from first principles, since the extent to which the problem is simplified by the assumption of a field that is static, or has only a few Fourier components may not otherwise be evident.

First, let us note the obvious feature that the difference in response of a system of atoms/ions to a fluctuating field $E(\omega)$ compared to that in a distribution of static fields is not just computational but is observable spectrally, Fig. 4, in that the fluctuating field may generate spectral intensity at frequencies in the line profile where no intensity can ever be caused by purely static effects.

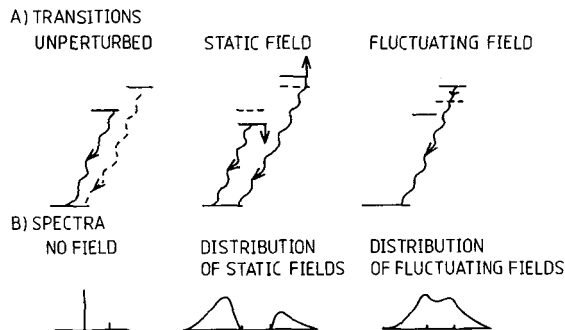


FIG. 4. COMPARISON OF STATIC AND FLUCTUATING FIELD EFFECTS ON A 3-LEVEL ATOM

7.2 Treatments to second order in E

Fig. 4 suggests that to some extent the line profile 'maps' the Fourier components of the field, $E(\omega)$. If second-order perturbation theory were adequate (i.e. for all components of E), this would in fact be true rigorously, as can be shown by writing down the line shape starting either from the time-dependent (auto-correlation function) approach (see e.g. Refs.3,12) or a suitable spectral representation (see e.g. Ref.13). If one then makes (say) an S-matrix type expansion in the perturbation, the result is that the line shape, $L(\Omega)$, is very simply related to the so-called 'dynamic structure factor', $S(k, \omega)$, (which also describes the light scattering properties of the medium) and thence to the spectral density of the field $E(\omega)$. In fact for hydrogenic transitions (see below for other species):-

$$L(\Omega) \propto \text{Im} \frac{1}{\Omega + \Sigma(\Omega)} \quad (11)$$

with $\Omega = \omega - \omega_0$ (12)
 where ω_0 is the line-centre frequency

$$\Sigma(\Omega) = \Delta(\Omega) + i\gamma(\Omega) \quad (13)$$

$$\text{and } \gamma(\Omega) \propto \int dk \frac{S(k, \Omega)}{k^2} \quad (14)$$

(the k^2 arises from the transform of the r^{-2} dependence of the electric field)

$$\text{and } S(k, \Omega) = -\frac{N}{4\pi} \frac{k^2}{k_D^2} \text{Im} \frac{1}{\epsilon(k, \Omega)} \quad (14)$$

where $\epsilon(k, \Omega)$ is the dielectric function of the medium (which is easily calculable provided the random phase approximation applies, i.e. $\Gamma < 1$), N is the density, and k_D is the inverse of the Debye length. (For lines from non-hydrogenic species the approach is marginally more complicated in interpretation in that the relevant frequencies in $S(k, \Omega)$ are shifted from the detuning $(\omega - \omega_0)$ by amounts depending on the difference frequencies of the atomic states coupled by $E(\Omega)$).

To the extent that second-order perturbation theory is valid, this treatment then provides an elegant formulation of the entire line-shape problem in which collective (possibly non-thermal) properties of the plasma are incorporated. Fig.5 shows a representation of $\gamma(\Omega)$ for a plasma in which electron and ion collective modes are somewhat enhanced (e.g. by a non-Maxwellian electron velocity distribution). The line profile-represented by $L(\Omega)$ then has the following general features:-

- (a) An overall Lorentzian shape for small $(\omega - \omega_0)$.
- (b) Discrete structure representing the ion and electron plasma modes as represented by the frequency dependence of the generalized width, $\gamma(\Omega)$.
- (c) A fall-off in $\gamma(\Omega)$ at large $(\omega - \omega_0)$ because of the corresponding decrease in $S(k, \Omega)$ for frequencies large compared to the electron plasma frequency.

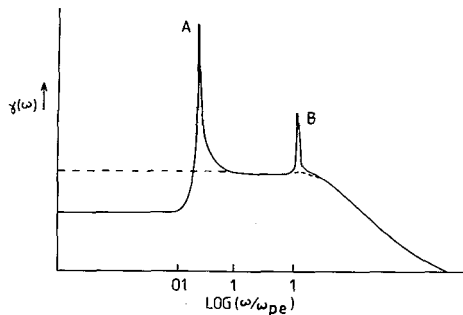


FIG 5 THE WIDTH FUNCTION γ AS A FUNCTION OF FREQUENCY. THE CASE SHOWN IS FOR A NON THERMAL PLASMA WHICH HAS PLASMON ENHANCEMENT B) AND ION ACOUSTIC MODE ENHANCEMENT A) THE THERMAL CONTRIBUTIONS FOR ELECTRONS ONLY ARE SHOWN WITH A DOTTED LINE

The difficulty with this type of approach taken on its own (even in low density plasmas) actually results from (c) above. For large $\Omega = (\omega - \omega_0)$, $L(\Omega)$ becomes dominated by the shift operator, $\Delta(\Omega)$, rather than by $\gamma(\Omega)$. $\Delta(\Omega)$ in turn depends on a Cauchy type integral over $\gamma(\Omega)$, and hence on all frequency components in the field $E(\Omega)$. Examination of the S-matrix type expansion for $\gamma(\Omega)$ shows that second-order is never adequate for the low frequency part. In fact to properly recover the ordinary quadratic Stark shift induced by a static field one has to go to infinite order in the S-matrix type expansion.

The point of this comment is to note the difference with the 'idealized' cases of static or monochromatic fields (whether weak or strong). In both these cases one has the basic information to sum the expansion to all orders in E if need be. In the static case one recovers the ordinary theory of the Stark effect, and in the monochromatic oscillating laser field case the complexity is simply that higher order terms in the S-matrix may generate a wide spectrum of atomic sum and difference frequencies. In contrast, the difficulty in the plasma case is that higher order terms in the S-matrix expansion couple more and more frequency components of E , $E(\Omega)$, $E(\Omega')$, etc. We can calculate the whole of this series only if the fluctuation properties represented by $E(k, \Omega)$ have rather special

properties.

At present there are three main ways of proceeding beyond second order in $E(k, \Omega)$:-

- 1) Fourier Decomposition of E into 'low' (static) and 'high' ('impact') frequency components.
- 2) 'Unified' (collisional) Theories
- 3) 'Model Microfield' Theories.

7.3 Fourier Decomposition Theories

The approach of theories under this heading - which cover a wide variety of different physical approaches, see 3,9,10,14, - is to divide the effects of the Fourier components of E into 2 classes, those ('low' frequency) that can be treated simply as 'static, for which the atomic response can then be calculated to any level required, and those ('high' frequency) for which the fluctuation properties are crucial (which are then normally only treatable to second order in $E(\Omega)$). The approach is then to 'renormalize' the atomic properties perturbed by a specific static field, E , to use the perturbed states as a basis to calculate the dynamic response of this specific emitter to the fluctuating high frequency components $E(k, \Omega)$ through second order (for emission/absorption processes the k dependence will be lost in the integration, but this is not necessarily true in some schemes involving laser scattering off the perturbed atoms. These, however, are feasible only in low density plasmas). The resultant spectrum for a particular static field E is then averaged over the probability distribution $P(E)$. Dependent on the plasma, the particular transition studied in a given species, or the particular region of the line profile of interest, the spectrum may be dominated by the static part alone, by the high frequency fluctuations alone, or by a combination of both.

One great merit of this type of approach is that the static field distribution $P(E)$ can be accurately predicted even in very dense plasmas, i.e. with $\Gamma > 1$. There are physical, mathematical and computational reasons for this. Physically $P(E)$ is well-behaved as the plasma density (and hence Γ) increases. The dominant contribution to the low frequency components of E is due to the ions, and increasing Γ corresponds to increasing ion-ion repulsion, and hence increasing short range order with the separation of all nearest neighbour ions becoming relatively constant (and hence close to R_0). $P(E)$ therefore lies somewhere between the extreme low-density (uncorrelated) Holtsmark limit and the simple nearest neighbour approximation and (when compared with other theoretical and experimental uncertainties) these do not differ all that greatly, (particularly in so far as the value of the most probable field is concerned), see Fig.1. Secondly, cluster-integral type theories - see in particular Hooper, Ref.4 - can include ion-ion and electron-ion correlations analytically very successfully and, as indeed proved by comparison with Monte Carlo type calculations, these analytic predictions remain good to surprisingly high values of Γ . Finally, 'molecular dynamics' codes are capable of directly predicting the $P(E)$ due to a large number of interacting particles regardless of Γ , see Weisheit and Pollock, and these show the close agreement of $P(E)$ with the simplest nearest neighbour model as Γ becomes large.

Theories with the effects of the overall E field split into separate static and fluctuating contributions, which still form the majority of computations are, of course to some extent typified by the old 'static ion'-'electron impact' theories of the type formulated early in the development of line shape theory in plasmas, e.g. by Baranger, 12, Griem and Kolb, 14, etc. However, we have deliberately avoided using the 'static'/'impact' terminology as far as possible above since the historical basis of such theories tended to be tied to a specifically 'collisional' (one-perturber) approach to the high frequency component, and to the view that a given perturber species can only contribute to one or other component of the overall perturbation (hence 'static ion-electron impact') an approach which has been somewhat superceded as particle correlations have become of interest.

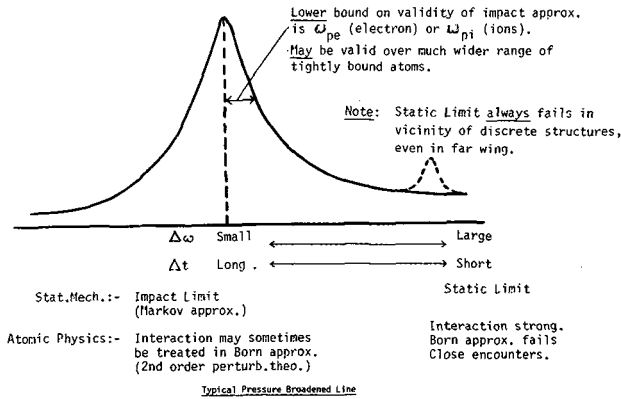
The formal basis for inclusion of electron correlations in the high frequency component as well as in the static contributions to the line shape was derived via a variety of approaches relatively early, (13,,15,16,17) and before the advent of interest in ultra dense laboratory plasmas, and leads directly to the type of approach already outlined in Section 7.2. However rather more controversy has surrounded the inclusion of ionic contributions to the dynamical part of E in addition to their usual incorporation via the static theory.

The difficulty here really stems from the whole basis of dividing the total effects of the field into low and high frequency parts, i.e. in deciding whether the field itself or its fluctuation properties cause the dominant effects. The 'obvious' way to carry out this division is to use the Fourier theorem and to compare the time variation of the field with the 'lifetime' of the atomic states of interest, i.e. to consider:-

$$dE(t)/E(t) \lesssim \Delta t^{-1} \quad (16)$$

and then to identify E(t) as being due to a specific set of perturbers (e.g. in the 'static ion'-'electron impact' theories as due to the overall ion field (dressed by the screening due to the electrons) in the case of the static part, and to the electrons in the dynamic part).

The classic approach is then simply to write the line shape in terms of the overall dipole autocorrelation function for the system as a whole, and thus identify $(\Delta t)^{-1}$ as simply Ω , the frequency separation from line-centre, see e.g. Baranger, 12, and also Fig.6.



This leads to the almost universal idea that for large Ω , i.e. in the far line wings the static theory automatically becomes valid regardless of the details of the transition concerned and that fluctuations can then be ignored. From this approach derive the still very widely used static ion, electron impact theories, since for Ω values of most interest in lines from low ionization stage species in low density plasmas, Eqn.(16), turns out to be such that almost always the overall ion field (dressed by the screening of the electrons) can be treated as purely static and the electron contributions as purely dynamic (i.e. dominated simply by the fluctuation properties).

This approach, however, is nowhere like as certain for the very different situations of high z-emitters in hot, dense matter. In this case the characteristic fluctuation time of the 'low frequency' part of the microfield (the ion plasma frequency) is almost always much larger relative to the overall linewidth, and from (16) dynamical effects of fluctuations in the ion microfield are relatively much more important, see e.g. Lee, 18, and also Lee and Freeman,

Reference 19.

However, even before the complications of dense high z plasmas are encountered, considerable caution needs to be adopted before using Eq.(16) as a valid basis for dispensing with the dynamical contributions of a particular component of the field, whether low or high frequency. Firstly, as pointed out several times by one of the present authors (see e.g. Burgess, Ref.1), widely exploited as Eqn.(16) is, it is still wrong! Large Ω simply does not automatically imply that a static field contributes to the profile at all at that point, and consequently dynamical contributions may well dominate even when the mean fluctuation time of the field is long. This is evident in Fig.4, where there would be no static contribution to the profile between the pair of unperturbed states, so that dynamic contributions are dominant regardless of the value of (and of whether Eqn.(19) predicts the static theory to be 'valid').

A second complication of fairly widespread practical consequence arises from the fact that the time of interest can only be defined in terms of the frequency separation from line centre in the case of a line of simple form (e.g. a Lorentzian). However, once discrete wing structures, such as field-induced forbidden components, become important this definition fails. (This was first pointed out by Burgess,20, and the consequent failure of static ion theories demonstrated experimentally by Burgess and Cairns,21, and incorporation of the previously neglected 'ion dynamic' contributions has been treated in a series of papers by Lee,18,22, as well as by many other authors, e.g. 23,24). This restriction on static theories is also important in high density high z plasmas (a) because of the significance of allowed-forbidden line pairs as useful density diagnostics and (b) because the same consideration applies near the discrete structures corresponding to the electron and ion plasma frequencies in Fig.5.

Since the realization that Eqn.(16) was not as universally valid as implied in the original treatments such as that of Baranger,Ref.12, several approaches have been adopted in attempts to incorporate ion dynamic effects, including effects of ion-electron correlations on the 'electron' microfield, whilst retaining as much as possible of the low frequency contributions to $L(\Omega)$ in an essentially static treatment. One approach due to Kogan,25, and also discussed by Hey,26, is to consider the effects of a time-varying field on a single Stark component, and then model these via a modified field distribution. The result obtained is:-

$$P(E/E_0) = P(E/E_0)^{\text{Holtsmark}} \frac{v_F^2}{v_S^2} \left\{ \left(1 + \frac{M_P}{M_R}\right) \Delta + \Delta' \right\} \quad (17)$$

where Δ and Δ' are functions varying with field, M_P and M_R are the perturber and radiator masses, and v_F/v_S is the ratio of the field fluctuation frequency to the static shift. This approach then allows some modelling of the extra broadening due to slow fluctuations in the static field, but is essentially based on a semi-collisional model and does not form a systematic basis for treatment of ion-ion and electron-ion correlations.

On the other hand, Lee,18,22, has demonstrated the application of a formalism originally due to Dufty,27, to the practical inclusion of ionic contributions in the calculation of the effects of the 'high frequency' (fluctuating) part of the field. This approach calculates to second order the dynamical effects of all Fourier components of the field (i.e. all contributions to $S(k, \Omega)$ regardless of frequency) whilst also retaining the full effects of the static field via $P(E)$. The treatment can be shown to be formally correct through second order in the fluctuating part of the field, and has been used to consider field-induced forbidden lines in low density plasmas, Lee,22, and Mahon Lee and Burgess,28, and recently by Lee to treat the broadening of lines of high z species in dense laser plasmas, Ref.18,19. A particularly interesting effect is the discovery (theoretically) of enhancements of broadening in plasmas of high mean perturber z due to increased electron-ion correlations, an effect closely related to predicted enhancements in Thomson scattering in high mean z plasmas,29. The same formalism is also well adapted to calculations of spectroscopic effects of enhanced excitation of electron plasma waves in plasmas containing non-thermal levels of fast electrons, or of ion modes in plasmas with differing ion and electron

temperatures.

7.4 'Unified' collisional theories

If the complete microfield E can be approximated as due to binary encounters with single perturbers occurring in isolation, then sufficient specification of the field fluctuations exists to allow $\gamma(\Omega)$ to be calculated to all orders in an S -matrix expansion. In other words, the atom-plasma interaction is described simply in terms of a binary collision T -matrix, and the complete line shape from line centre to line wings follows directly from Eqn.(11). The theory is thus 'unified' in the sense that it avoids specifically making any static/dynamic division.

'Unified' theories can be derived either classically or quantum mechanically (the classical approach being much more cumbersome!) and classical path approximation calculations of line shapes have been performed particularly by Vidal, Cooper and Smith,³⁰ and by Voslamber,³¹. At low densities the theories work for electron broadening since the region in the line wing at which the second order expansion is inadequate already corresponds to very large k in $S(k, \Omega)$, and thus to very close and hence purely binary encounters. However, such a binary collisional approach is not valid for the low frequency field components dominated by the ions.

In general, 'unified' theories are of little, if any, use in the areas under discussion in the present paper, since under high density conditions line shapes of interest are dominated by those regions of $S(k, \Omega)$ (and hence $\gamma(\Omega)$) in which particle correlations play an important and possibly crucial) role.

7.5 'Model Microfield' Theories

The remaining alternative, developed originally by Frisch and Brissaud,³² and also used and discussed in particular by Seidel,^{33,34}, attempts to circumvent the difficulty of the higher order moments of $E(\Omega)$ by a purely model/computational solution. That is, an attempt is made to find a model of the fluctuating microfield which has the right statistical properties at both high and low frequencies and allows a complete calculation of the atomic response. The model microfield adopted is chosen so that it is completely specified by the instantaneous probability distribution, $P(E)$ and the second-order correlation function, i.e. $\langle E(t).E(0) \rangle$ (or equivalently the spectral density function $E^2(\Omega)$). (This amounts to a statement that the exact details of the way the field changes do not matter). Since the model field is entirely specified by these parameters there are no difficulties involving the computation of higher order terms involving several separate field components $E(\Omega)$, $E(\Omega')$ etc. For the random (close encounter) part of the field this approximation may be reasonable, although rather few detailed calculations have yet been made. However, it appears much less clear whether the method can handle collective effects in non-thermal plasmas, particularly for instance those involving both electron and ion modes. Interesting discussions of the relation of the Model Microfield theory to more orthodox techniques have been given by Seidel,³³, and by Dufty,³⁵.

8. Fields in Very Strongly Coupled Plasmas, ($r \gg 1$).

In those plasmas where there is less than one particle per Debye sphere we may need to reconsider the basic approach presented above. We have dealt with the change in the low frequencies microfield due to strongly coupled systems but we have not discussed the basic assumption that the plasma field can be treated as a perturbation to the radiator state. In a strongly coupled system this perturbative approach becomes suspect.

The only alternative is to attempt to compute from first principles a self-consistent set of radiator states in the presence of the plasma, i.e. to account for the plasma by a complete renormalization. The major difficulty with this method is that the plasma is treated as essentially static and we thus obtain a frequency independent set of radiator energies and wavefunctions. Clearly this

is only a starting point for the time history of plasma perturbations plays a crucial role in defining the details of the interactions (as is clear from the importance of fluctuations at lower densities).

In general the method is to treat each nucleus as being at the centre of a cell containing z electrons (whether bound or free) and to solve for the electrostatic potential using the Poisson equation together with a distribution function for the electron in the cell. The result is a self-consistent electrostatic potential and charge distribution on the atom. There are a number of methods used to calculate the self-consistent electron charge distribution which are basically similar.

8.1 Thomas-Fermi Model

The TF theory, see e.g. 36 assumes the atom is at the centre of an ion sphere, of radius $R_0 = (3/4\pi N)^{1/3}$ and other ions are excluded from the cell. The electrostatic potential $V(r)$ is calculated by solving Poisson's equation

$$\nabla^2 V(r) = 4\pi e(N_e(r) - Z\delta(r)) \quad (17)$$

where $N_e(r)$ is the total, bound and free, electron density. The electron density is assumed to be given by

$$N_e(r) = (dp/\hbar^3) \cdot 1/(1 + \text{EXP}((p^2/2m - eV(r) - \mu)/kT)) \quad (18)$$

the Fermi Dirac function. The difficulties with this model are numerous and although it provides for ionization potential depression and mean ionization state by allowing a calculation of the number of 'free' electrons it cannot be used to predict spectroscopic information since it ignores the individual electron wavefunctions and replaces them with a smooth TF function for the charge distribution.

8.2 Quantum Statistical Model

The Quantum Statistical Model, 36,37, attempts to improve the TF theory by inclusion of exchange interactions and effects of gradients. However, the model still provides no really spectroscopic information.

8.3 Self-Consistent Schroedinger Equation Method

Poisson's equation for a hydrogenic ion in a plasma is

$$\frac{1}{r} \frac{\partial^2}{\partial r^2} rV(r) = 4\pi e^2 (N_e(r) + |\psi(r)|^2 - N_i(r)) \quad (19)$$

$$\text{where } N_i(r) = N_0 \text{EXP}(-V/kT) \quad (20)$$

$$N_e(r) = N_0 \int_{p_0}^{\infty} \frac{p^2}{2} \text{EXP}(-\frac{p^2}{2m} + \frac{V}{kT}) dp \quad (21)$$

where $p_0 = \sqrt{2mV}$ is the smallest momentum of a free electron. This is augmented by the Schroedinger equation for $\psi_\eta(r)$ the bound electron wavefunction for state η .

$$-\frac{\hbar^2}{2m} \nabla^2 \psi_\eta + V(r)\psi_\eta = E_\eta \psi_\eta \quad (22)$$

This set of formulae can be iterated until a self-consistent potential and charge distribution are found and for high z hydrogenic ions has been used to provide spectroscopic information on shifts of the energy levels and the ionization potential depression. The results for the ionization potential depression are quite similar to the Stewart-Pyatt model(36,38) qualitatively but distinct

quantitatively. The Stewart-Pyatt potential depression as presented by More,³⁶ is

$$\Delta I = -\frac{3}{2} \frac{Ze^2}{R_0} \{ [1 + (\lambda_D / R_0)]^2 / 3 - (\lambda_D / R_0)^2 \} \quad (23)$$

This model therefore aims to predict self-consistently the shift of levels due to the static plasma perturbations going beyond the perturbative type of microfield approach and these shifts are in some senses similar to the plasma polarization shift discussed in numerous references (see Burgess,¹). Quantitatively the shift turns out to be about 5 eV at a plasma electron density of $2 \cdot 10^{23} \text{ cm}^{-3}$ for the Lyman α line of Ne X. This shift should be large enough to be observable but has not yet been identified in experimental studies of this line and provides a source of some interesting experimental possibilities.

8.4 Average Atom Model

In many situations the possibility of determining the detailed information necessary to predict the spectroscopic emission is beyond theoretical/computational power. For example, hot laser plasmas generated in planar target experiments may emit spectral features which arise from numerous ionization stages and from, say, 10 or more configurations in each stage (39). The only treatment is then the average atom model which has found use at both high and low density. The basis of the average atom theory (36) is to use screening coefficients which allow the calculation of the self-consistent energy levels given the atomic shell populations. The energy levels derived are then used to recalculate the populations in either an LTE or non-LTE model, and thus to re-derive the screening coefficients.

The outcome is a set of shell populations which to some extent describe both bound and free electron states. With this simple model some important aspects of radiation transport and opacity can be included in, for example, hydrogenic simulations. However, specific level shifts/widths and other detailed pieces of spectroscopic information are not available.

9. Effects on Spectra

There is already a very extensive literature on the effects of dense plasmas on the spectra of high z-emitters, and the reader is referred to References 1, 2, 36, 40, 41, 42 for general survey purposes.

The most commonly studied features (both computationally and experimentally) have been line widths/shapes. In fact, quantitative observations of pressure broadening of lines of high z-emitters in dense laser plasmas and comparison with theory now dates back fully 15 years, see Burgess, Fawcett and Peacock, 43, but the present high level of experimental interest has only developed with the need to diagnose conditions in laser compressed plasmas relevant to fusion research. For diagnostic purposes line shape measurements have a number of advantages over line intensities in that line widths are easily measured with none of the difficulties in instrumental calibration inherent in intensity measurements, and in particular are directly sensitive to density, perhaps the most important parameter in laser compression experiments. Densities derived from line profile measurements have now been reported by many authors, see e.g. Ya'akobi 40, 44; Hauer 41; and Kilkenny, Lee, Key and Lunney 46.

One interpretive problem, affecting both line profile and intensity measurements is the plasma opacity (which in laser compressed plasmas may be high not only in the line core but also in the Bremsstrahlung continuum). Here line shape measurements have advantages over line intensity techniques, in that profile changes along a series are relatively well defined, and comparison of measured widths with predicted values for the first few series members of say a hydrogen-like species allows a unique determination of both optical depth and plasma density, see Ref.45. Closely related is the use of allowed-forbidden line pairs in helium-like spectra, field-induced dipole-forbidden lines being strong features of almost all spectra at these high plasma densities, and theoretical

treatment of such features has been attempted using a variety of theoretical approaches, see 18, 46, 47.

A number of other effects of potential interest in studying the microscopic physics/statistical mechanics of very dense plasmas should also be observable in laser compressed plasmas, but the level of quantitative experimental study of these is still relatively very limited. As discussed in Section 7.2, the form of $S(k, \Omega)$ should give rise to discrete satellite features in spectral line profiles corresponding to the electron and ion plasma frequencies. An alternative way of considering such features is as multi-quantum transitions, involving both photons and quanta of the plasma mode, and such satellites are interesting not only for the direct information provided on plasma conditions, but also because for such discrete features the use of a theory going only to second-order in the fluctuating field is rather suspect. Apparent plasma satellite features have been seen in a number of spectra from dense laser plasmas, see e.g. Fig.7, but detailed comparison with theory is limited by the limited extent to which simultaneous space and time resolution can be obtained in spectra from these very small transient sources, and by the effects on overall the line profile on the strong density gradients in the source.

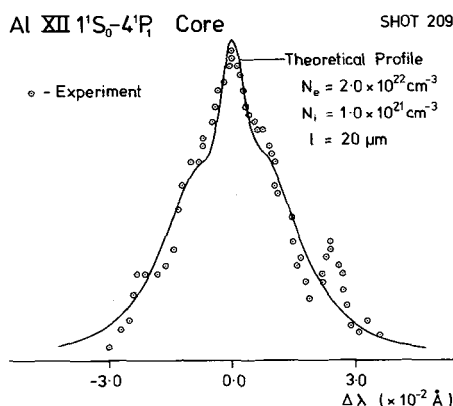


Fig.7. Aluminium XII transition observed in an "exploding pusher" laser plasma (joint Imperial College-Queen's University Belfast-Rutherford Appleton Laboratory experiment, spectrum courtesy of Dr. J. Killkenny).

Similarly, laser compressed plasmas should allow observations of another long-discussed high density effect on spectra, that of ionization potential depression, 36, 38, 48. At lower densities spectroscopic observations of this effect have always been severely limited (a) by the fact that it is small and (b) since at low densities it is then invariably smaller than the Inglis-Teller limit, i.e. the point at which a line series merges. Despite much theoretical speculation, very little is yet known from direct experimental work. Laser plasmas offer two special possibilities in this respect. Firstly the I.P. depression may be very large (whatever model is used to calculate it), for instance in a Neon plasma at 3.5×10^6 Kelvin and a density of only $4 \times 10^{22} \text{ cm}^{-3}$ being large enough on either the low density Debye or high density Stewart-Pyatt models that $n=5$ of Ne X no longer exists (i.e. Lyman-delta should disappear), 49. Secondly, under appropriate conditions the I.P. depression may actually exceed the series merging effects, and hence be observable directly for the first time (see Lee, 49, for a plot of the plasma conditions relevant for this to occur for the

Lyman-alpha and Lyman-gamma lines of Ar XVIII).

Closely related, as already mentioned in Section 8.3, the screening effects on bound orbitals of the local plasma polarization cloud near an emitter should give rise to measurable line shifts for transitions from high z species (even for hydrogenic species whose lines are unshifted by ordinary pressure broadening processes). Theoretical estimates of such screening or polarization shifts vary rather widely, see e.g. 1, 3, 8, and observations would be very important in improving understanding of the various theoretical approaches available, particularly as plasma densities become such that $\Gamma > 1$, see Section 8. Experimental observations have been attempted, but results are presently still limited by restrictions (set by X-ray flux) on space and time resolution, and also in some cases by difficulties in cross-calibration of line wavelengths with standard sources occurring when slitless crystal spectrometers are used (for which the line image position is source geometry dependent).

One other high density effect for high z species is that Doppler narrowing may occur even for soft X-ray transitions, as first pointed out by Burgess, 1, and modelled for specific plasma conditions by Burgess, Everett and Lee, 54. The interest here is that this effect could provide a direct monitor of the ion mean free path against 90 deflection, and hence of ion-ion correlations. The main experimental difficulty is the large competing effect of pressure broadening. At intermediate plasma densities (10^{22} cm $^{-3}$) this can partly be overcome by looking at transitions in high z stages of relatively heavy species chosen such that pressure broadening effects should be small. $n=2$ to $n=2$ transitions in transition elements offer some prospects in this regard, and Burgess, Everett and Peacock, 55, have studied such lines in Vanadium XV through XX, but conclusive identification of Doppler narrowing would require a much more thorough analysis of the effects of plasma geometry and density gradients than presently feasible for these lines.

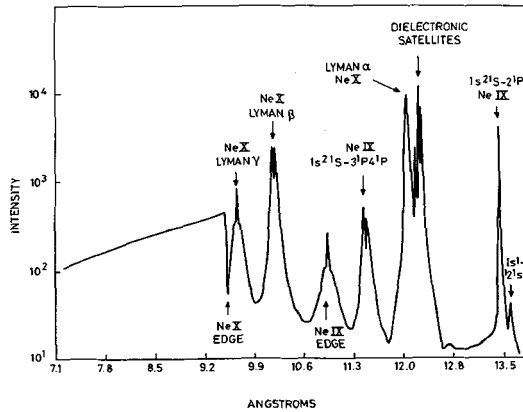


Fig.8. Computed spectrum of a Neon Plasma (see R.W.Lee in Annual Report of the Laser Division, Rutherford Appleton Laboratory 1981).

Finally, it is important to notice the extent to which the information available in the spectra of these sources is complicated by the very extreme plasma conditions. Even at rather modest densities, only a relatively few spectral lines are observable because of the effects of large line width, line overlapping, I.P. depression and series merging, which mean that theoretical predictions need to consider not just isolated lines but the entire spectrum, combining predictions of pressure effects on line shapes and series limits with

computations of radiation transport in the strongly spatially varying plasmas. The extent to which the spectrum is usually limited to only a few lines is shown in Figure 8, which is a simulation of a uniform Neon plasma in the region of the Ne IX and Ne X series limits for a plasma temperature of 3.5×10^6 Kelvin and electron density of 4×10^{22} cm⁻³. This type of simulation (in this case by Lee) has also been extended to consider the effects of detailed plasma geometry and density gradients, see e.g. Ref. 50, and also Refs.51, 52, 53.

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