

An Interpretation of Anomalous Intensity Ratios Within the Hydrogen-like L_α Doublet in Optically Thick Laser-produced Plasmas

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Abstract

It is shown that the anomalous intensity ratios of the fine-structure components of $\text{MgXII } L_\alpha$ and $\text{AlXIII } L_\alpha$ observed in optically thick laser-produced plasmas can be explained in terms of the usual radiative and collisional processes provided allowance is made for the effect of differential motion on the plasma opacity. The Doppler effect associated with differential expansion causes the red component of the doublet in the peripheral plasma region to overlap the blue component in the central region, leading to preferential absorption of the blue component. Model calculations give intensity ratios in good agreement with experiment, and explain such details of the line profile structure as the so-called short-wave satellite and a component spacing in excess of the theoretical optically thin value. This study further demonstrates the major role played by differential motion in shaping the emergent profile of an optically thick line from a laser-produced plasma.

1. Introduction

The intensity ratio of the fine-structure components of the hydrogen-like L_α doublet of third sequence elements (sodium–argon) is currently of interest as a density diagnostic in high temperature plasmas, in density regimes which include laser-produced plasmas on the one hand and astrophysical and tokamak plasmas on the other (Vinogradov *et al.* 1974, 1977; Boiko *et al.* 1978*b*, 1979; Skobelev *et al.* 1978). As summarized in Section 2 below, recent measurements on laser-produced plasmas under optically thin conditions (Boiko *et al.* 1978*b*, 1979), when interpreted with the theory of Vinogradov *et al.* (1977), give values of electron density in good agreement with other independent measurements, thereby lending support to the theory.

Also summarized in Section 2 are recent measurements (for $\text{MgXII } L_\alpha$ and $\text{AlXIII } L_\alpha$) on laser-produced plasmas under optically thick conditions which show the red (and normally weaker) component of the doublet as being up to twice as strong as the blue component (Bayanov *et al.* 1976; Beigman *et al.* 1976, henceforth referred to as BBPF; Boiko *et al.* 1978*a*). This result has been labelled anomalous by BBPF, who argue that it cannot be explained in terms of the usual radiative and collisional processes (an argument which is refuted here); and they claim qualitative agreement with theory once the usually forbidden single-photon transition $2s_{1/2}-1s_{1/2}$ is included in the rate equations (a claim which is also refuted here).

The object of this paper is to show that the so-called anomalous intensity ratios can be explained in terms of the usual radiative and collisional processes *provided allowance is made for the effect of differential motion on the plasma opacity*. While differential motion is recognized as an intrinsic property of laser-produced plasmas (though it was not considered by BBPF), the variety of spectral effects arising therefrom is still being surveyed (Irons 1975, 1976, 1980*a*, 1980*b*; Nicolosi *et al.* 1978,

and references therein; Yaakobi *et al.* 1979). In the present case, as described in Section 3, differential expansion causes the red component of the doublet in the outer region of the plasma to overlap the blue component in the central region, leading to preferential absorption of the blue component. A similar mechanism has previously been proposed by Hewitt and Noerdlinger (1974) to explain intensity anomalies in expanding stellar atmospheres.

Model calculations, in Section 4, put the present work on a semiquantitative basis. They illustrate the effect of differential motion on a doublet line profile for a wide range of velocity conditions. For conditions typical of laser-produced plasmas they give intensity ratios in good agreement with what has been observed, and they furthermore explain such details of the line profile structure as the so-called short-wave satellite (Bayanov *et al.* 1976) and a component spacing in excess of the theoretical (optically thin) value (Aglitskii *et al.* 1974; Boiko *et al.* 1977, 1978a). The model calculations also illustrate the rather intriguing fact (commented upon previously by Irons 1980a) that, under certain conditions (as encountered in laser-produced plasmas), the emergent profile of an optically thick line may be so distorted by differential motion as to resemble the profile of the same line when optically thin.

The only comparable work relating to laboratory sources appears to be that of Braun *et al.* (1970), who have reported anomalous intensity ratios (red component about 1.5 times more intense than the blue component) for deuterium L_α from a microwave-powered lamp. Interestingly enough, Braun *et al.* explain this result in terms of a relative shift in line centre between the central (emitting) region of the source and the outer (absorbing) region, the shift being ascribed to a resonance interaction. The present explanation, for Mg XII L_α and Al XIII L_α in a laser-produced plasma, is similar in principle, except that the shift is here ascribed to the Doppler effect associated with differential motion.

One final point to emerge from the BBPF paper is the question of whether the escape-factor technique (of allowing for self-absorption in rate equations) is valid for the individual components of an overlapping doublet. This is of wider interest than just the present case and is discussed in detail elsewhere (Irons 1980b, henceforth referred to as Paper 1).

2. Current Status of Experiment and Theory, with Critical Comments on BBPF Study

The L_α doublet of hydrogen-like ions is composed of the two transitions $2p_{3/2}-1s_{1/2}$ and $2p_{1/2}-1s_{1/2}$ which we will refer to as transitions 1 and 2 respectively. Define a ratio

$$\beta = E_2/E_1, \quad (1)$$

where E_1 and E_2 denote the photon efflux (i.e. the rate of emergence of photons from the source) for transitions 1 and 2 respectively. The current status of experiment and theory regarding the determination of β will now be reviewed, firstly for the optically thin case and then for the optically thick case.

(a) Optically thin case

In the optically thin case, since the Einstein coefficient for spontaneous emission is the same for both transitions, the definition (1) of β reduces to

$$\beta = \varepsilon \equiv n_2/n_1,$$

where n_1 and n_2 denote the average population densities of quantum states $2p_{3/2}$ and $2p_{1/2}$ respectively.

Hutcheon and McWhirter (1973) discussed the magnitude of ε in qualitative terms, and concluded that any departure of ε from the statistical value (due to collision-induced transitions between the 2s and 2p states) would be slight. Vinogradov *et al.* (1977) set up the appropriate quasi-steady state rate equations and presented calculations of β ($= \varepsilon$) for Mg XII L_α . Additional calculations were subsequently presented by Boiko *et al.* (1978*b*, 1979) and Skobelev *et al.* (1978), for Mg XII L_α -SXVI L_α . These results may be summarized as follows. In the low density (coronal) limit, and in the high density (Boltzmann) limit, the states $2p_{3/2}$ and $2p_{1/2}$ are populated in proportion to their statistical weights, and $\beta = 0.5$. At intermediate electron densities ($10^{13} \lesssim N_e \lesssim 10^{21} \text{ cm}^{-3}$), where there is partial intermixing of the 2s and 2p states, β rises to a plateau value β_{max} . When the external process which populates the state of principal quantum number $n = 2$ is collisional excitation from the ground state, β_{max} is fairly insensitive to temperature and takes values in the range 0.70–0.78 for the ions Mg XII–SXVI. As commented by Vinogradov *et al.* (1977), the values calculated for β suffer from errors in the collisional rate coefficients and from the neglect of cascade population of the state $n = 2$ from higher states. In a laser-produced plasma, cascade population of the state $n = 2$ becomes progressively more important with the progressively lower temperatures and densities encountered with increasing distance from the target (e.g. Irons and Peacock 1974).

Over the range of electron densities $N_e = 10^{19}$ – 10^{23} cm^{-3} , which is of interest with laser-produced plasmas, β (for Mg XII L_α -SXVI L_α) is dependent on N_e , and decreases from $\sim \beta_{\text{max}}$ (at $N_e = 10^{19} \text{ cm}^{-3}$) to ~ 0.5 (at $N_e \approx 10^{23} \text{ cm}^{-3}$). Boiko *et al.* (1978*b*, 1979) have measured β for Mg XII L_α in a laser-produced plasma which contained only 1.5% Mg and was optically thin to Mg XII L_α , and have deduced values of N_e for distances out to $\sim 200 \mu\text{m}$ from the target, which are in agreement with other measurements. On the assumption that the line PXV L_α is optically thin in a pure phosphorous laser-produced plasma, Boiko *et al.* (1978*b*, 1979) have similarly deduced values of N_e which are in agreement with other measurements (see the further discussion in Section 4*a* below). These results appear to be the only confirmation of the optically thin calculations. In particular, they confirm the relevance of calculations based on the assumption that the external process populating the state $n = 2$ is collisional excitation from the ground state. In the other case considered by Vinogradov *et al.* (1977), of a cooling plasma in which the external process populating the state $n = 2$ is photorecombination, β_{max} is sensitive to temperature as well as to density, and takes values in the range 1–2 for Mg XII L_α , at least for the various parameter values considered by Vinogradov *et al.* which are pertinent to laser-produced plasmas.

One implication of these calculations, which has not so far been commented upon, is that, if ε takes a value other than the statistical value of 0.5, then the usual wavelengths quoted for the weighted line centre of L_α , which assume $\varepsilon = 0.5$ (Garcia and Mack 1965; Edlén 1966), are inappropriate. Since astrophysical and tokamak plasmas as well as laser-produced plasmas are involved, this is a matter of some concern. For example, for Mg XII L_α the line centre weighted according to $\varepsilon = 0.7$ is 8.42143 \AA ($1 \text{ \AA} \equiv 10^{-10} \text{ m}$), whereas the line centre weighted according to $\varepsilon = 0.5$ is 8.42100 \AA . The difference of 0.0004 \AA , remains approximately constant for other elements in the range magnesium–sulfur and is probably of no practical significance

(any measurement with a resolution approaching 0.0004 \AA would, after all, resolve the doublet as two transitions).

(b) *Optically thick case*

After a summary of the two experiments which report β values in optically thick laser-produced plasmas, we consider the interpretation of these values by BBPF.

(i) *Experiment*

Bayanov *et al.* (1976) have reported β values for $\text{Mg XII } L_\alpha$ as a function of distance x from the surface of a pure magnesium target. The ratio β was observed to increase from 1.3 at $x \approx 0.07 \text{ mm}$ to 2.2 at $x = 0.23 \text{ mm}$, and then to decrease to 1.3 at $x = 0.4 \text{ mm}$.* Such values are well in excess of β_{\max} for optically thin conditions, at least for the case where the state $n = 2$ is populated primarily by collisional excitation from the ground state. Bayanov *et al.* have commented that such values 'could arise if the electron temperature were $T_e \sim 2 \text{ keV}$ in a cooling plasma, when the main mechanism populating the levels with $n = 2$ would be photo-recombination. However, there is as yet no confirmation of the existence of such high temperatures'.

BBPF have reported results which show that:

- (1) for $\text{Mg XII } L_\alpha$, β varies with the magnesium concentration, from 1.7 for a pure magnesium plasma to 0.7 for a plasma produced from a (duraluminium) target containing 1.5% magnesium;
- (2) for $\text{Al XIII } L_\alpha$, from a pure aluminium target, β decreases with distance from the target, having a maximum value of 1.4 at $x = 0$.

Again, values such as $\beta = 1.7$ and 1.4 are well in excess of β_{\max} for optically thin conditions. BBPF interpreted the trends in the cases 1 and 2 as confirming the role of opacity in determining the value of β (this follows in case 2 because the line density decreases with distance from the target). We note also the observation of $\beta \approx 1$ for $\text{Si XIV } L_\alpha$ (Boiko *et al.* 1977) which is in excess of β_{\max} .

We note in passing an observation, by Kononov *et al.* (1971) in a laser-produced plasma, of relative intensities within a multiplet (in each of Al IX , X and XI) which are not proportional to the statistical weight of the emitting substates. However, there appear to be no striking anomalies here, as with the $\text{Mg XII } L_\alpha$ and $\text{Al XIII } L_\alpha$ doublets.

(ii) *A Refutation of the Inequality $\beta < 1.3$ derived by BBPF for the Optically Thick Case*

The optically thick case has not received the same detailed theoretical study as the optically thin case. BBPF have incorporated the effect of opacity into the theory outlined in Section 2a by allowing the spontaneous transition probabilities to take effective values A_1 and A_2 for transitions 1 and 2 respectively, which are smaller than the optically thin value A . In this case equation (1) becomes

$$\beta = A_2 n_2 / A_1 n_1 = (A_2 / A_1) \epsilon. \quad (2)$$

* The densitograms in Fig. 5 of the paper by Belyaev *et al.* (1977), in another experiment involving $\text{Mg XII } L_\alpha$, indicate a value of $\beta < 1$ at $x = 2 \text{ mm}$.

BBPF derived an upper limit to β (summarized by Boiko *et al.* 1978*a*), which we shall now consider in some detail. Following BBPF, we shall firstly consider the case $\varepsilon \leq 0.5$ and then the case $\varepsilon > 0.5$.

To treat the case $\varepsilon \leq 0.5$ BBPF invoked the approximation

$$A_1 = Ag(\tau_{0,1}), \quad A_2 = Ag(\tau_{0,2}) \quad (3a, b)$$

for transitions 1 and 2 respectively, where $\tau_{0,1}$ and $\tau_{0,2}$ ($= \frac{1}{2}\tau_{0,1}$) denote the respective optical depths at line centre and $g(\tau)$ refers to the transmission factor (Holstein 1947, 1951). This latter factor has been tabulated and plotted by a number of authors (for references see Section A.1 of the appendix to the paper by Irons 1979*a*) and, for Doppler broadening and $\tau \gtrsim 2$, it has the form

$$g(\tau) \approx \tau^{-1}(\pi \ln \tau)^{-\frac{1}{2}}, \quad (4)$$

where, for transitions 1 and 2, read $\tau = \tau_{0,1}$ and $\tau_{0,2}$ respectively. It follows from the equations (3) that A_2/A_1 has a maximum value of 2.6 (for $\tau_{0,1} = 3$), in which case (for $\varepsilon \leq 0.5$) we have

$$\beta \leq 1.3. \quad (5)$$

This result is, of course, dependent on the approximation in equations (3), which is based on the assumption that the emitting atoms are concentrated towards the centre of the source (see e.g. Irons 1979*b*). By comparison, for the spatial model in Paper 1, A_2/A_1 has a maximum value of 2.8 (for $\tau_{0,1} = 30$), implying $\beta \leq 1.4$. For a uniformly excited source the value of A_2/A_1 increases monotonically with $\tau_{0,1}$, having a value of 1.9 at $\tau_{0,1} = 100$, implying $\beta \leq 0.95$ for $\tau_{0,1} \leq 100$.

However, the equations (3) are only valid for isolated lines, that is, $g(\tau)$ does not allow for the fact that the two (Doppler-broadened) components of the L_α doublet overlap and that photons emitted in one component may be absorbed in the other, and vice versa. At first sight this effect may appear to be of little consequence, since the overlap only occurs in the component line wings. However, the line wings contain the very photons that escape and contribute to $g(\tau)$ (see e.g. Irons 1979*b*). Model calculations (Paper 1) which allow for overlap show that, while A_1 and A_2 are significantly affected by overlap, the ratio A_2/A_1 nevertheless remains approximately equal to the isolated line value. However, this situation is dramatically altered if the source is moving differentially. Model calculations (Paper 1, and Section 4 below), for conditions typical of Mg XII L_α close to the target surface in a differentially expanding laser-produced plasma, show that A_2/A_1 may readily attain a value of 4, in which case (for $\varepsilon \leq 0.5$) we have $\beta \lesssim 2$.

Now consider the case $\varepsilon > 0.5$. BBPF adopted a basically different approach for this case. They utilized the rate equations describing population density to obtain an expression for the ratio $A_2 n_2/A_1 n_1$ and thence for β (equation 2). Using this expression they showed that, for $\varepsilon > 0.5$,

$$\beta \lesssim 1. \quad (6)$$

Implicit in this procedure, however, is the assumption that the effective transition probabilities A_1 and A_2 which enter into the rate equations (and allow for photo-excitation) are the same as those which appear in the expression (equation 2) for β

(and allow for the loss of photons from the photon efflux). The conditions under which this assumption is valid have been discussed elsewhere (Irons 1979*a*; Paper 1), and may be summarized as follows.

It is sufficient to consider the example of transition 1, though the remarks to follow apply equally to transition 2. The effective transition probability, as it enters into the rate equations, may be expressed as $A_1 = A A_1$, while the effective transition probability as it enters into the formula (2) for β may be expressed as $A_1 = A \Theta_1$, where the parameters A_1 and Θ_1 may be referred to as the mean net radiative bracket and the mean escape probability (or escape factor) respectively. The point at issue is the validity of the equality

$$A_1 = \Theta_1. \quad (7)$$

If transition 1 could be considered as an isolated line (i.e. if we could neglect the presence of transition 2), then equation (7) would hold in general, provided that (Irons 1979*a*):

- (i) the rate equations are formulated for the population density averaged over the whole source, and
- (ii) the photon efflux refers to the rate of emergence of photons in all directions from, in general, the whole surface area of the source.

It follows that, even for an isolated line, the validity of equation (7) is doubtful, since an observation usually samples the average density and photon efflux along a specific direction (the line of sight). More significantly, however, transition 1 cannot be considered as an isolated line because of the overlap with transition 2. We would normally expect that a contribution to A_1 from other than transition 1 photons would render equation (7) invalid. However, a detailed study for a doublet such as L_α (Paper 1) has shown that there are conditions under which equation (7) is valid, namely when (in addition to conditions (i) and (ii) above):

- (iii) the upper states $2p_{3/2}$ and $2p_{1/2}$ are populated in proportion to their statistical weights, and
- (iv) the emission and absorption profiles are equal (as is normally assumed to be the case) and constant in space.

It follows from (iii) that, in the present case, the validity of equation (7) is further in doubt because ε is not necessarily equal to the statistical value of 0.5. More significantly, however, is the fact (which follows from (iv)), that equation (7) is rendered invalid by the presence of differential motion, which introduces a Doppler shift into the emission (and absorption) profile which is not constant in space.

On the basis of equations (5) and (6), BBPF concluded that 'for an optically thick plasma allowance for the usual radiative and collisional transitions does not make it possible to obtain $\beta > 1.3$ '. The derivation of this result has now been shown to be invalid. Model calculations, which have been quoted and will be discussed in detail in Section 4, show that for ε as small as 0.5 it is possible for β to attain a value of 2, comparable with what is observed, provided allowance is made for the effect of differential motion on the plasma opacity.

(iii) *Interpretation of Anomalous β Values by BBPF*

Being convinced that the anomalous β values could not be explained in terms of the usual radiative and collisional transitions, BBPF sought an explanation involving

an additional such transition. While we may not share their conviction, their explanation nevertheless deserves close attention.

BBPF proposed additional radiative de-excitation of the $2s_{1/2}$ state via the single-photon $2s_{1/2}-1s_{1/2}$ transition (i.e. additional to the two-photon transition). This transition is normally forbidden (see e.g. Marrus 1977; Marrus and Mohr 1978) but it can be induced through ion-ion interactions, as described by BBPF, who approximated the corresponding transition probability A_3 . The actual magnitude of A_3 is not critical to the argument of BBPF, only that it be sufficiently large. The effect of such an additional transition is as follows. Provided we have

$$A_3 \gg 2b, \quad (8)$$

where b denotes the rate coefficient for collisional de-excitation of $2p_{3/2}$ to $2s_{1/2}$ (and $2b$ denotes the rate coefficient for the inverse process), then the additional transition has the effect of reducing the (quasi-steady state) population density n_3 of the $2s_{1/2}$ state, thereby reducing the rate of collisional excitation from this state. The net effect is an increase in the ratio ϵ and thence in β .

BBPF have argued as follows. Firstly they pointed out that the energy spacing $2s_{1/2}-2p_{1/2}$ is so small (see e.g. Kugel and Murnick 1977) that the transition $2s_{1/2}-1s_{1/2}$ is practically indistinguishable in wavelength from transition 2 and so contributes to an enhancement of the observed β value, which is now defined by

$$\beta = (A_2 n_2 + A_3 n_3) / A_1 n_1. \quad (9)$$

Secondly they utilized the rate equations describing population density to obtain an expression for the ratio on the right-hand side of equation (9). According to BBPF this expression reduces to

$$\beta = \frac{1}{2} \{1 + \mu + b(3 + \mu) / A_1\} \quad (10)$$

when $A_3 \gg b$,* where μ denotes the rate of population of $2s_{1/2}$ by collisional excitation from the ground state expressed as a fraction of the corresponding rate for $2p_{1/2}$. Upon substituting the observed value of β into this equation, BBPF obtained a value for A_1 which, when interpreted by means of equation (3a), gave a value of $\tau_{0,1}$ in the range 3–8. On the basis of this analysis, plus a subsidiary analysis involving targets of magnesium oxide and duraluminium, BBPF considered that they had demonstrated a satisfactory qualitative agreement between theory and experiment.

* The three rate equations describing n_1 , n_2 and n_3 (equation 11 of BBPF) when summed give (using the present subscript notation)

$$A_1 n_1 + A_2 n_2 + A_3 n_3 = (3 + \mu)f.$$

This equation, when combined with the rate equation describing n_1 , leads to the present equation (10) *provided we have* $2f > 2bn_3$, that is, provided the rate $2f$ of population of $2p_{3/2}$ by collisional excitation from the ground state exceeds the rate $2bn_3$ of population of $2p_{3/2}$ by collisional excitation from $2s_{1/2}$. Radiative de-excitation of the state $2s_{1/2}$ to the ground state does not enter explicitly into the derivation of equation (10). The condition $A_3 \gg b$ quoted by BBPF is not a necessary condition for equation (10) to follow. However, it is a sufficient condition, as outlined beneath equation (8), i.e. if A_3 is sufficiently greater than b then n_3 will be sufficiently small so that the necessary condition $2f > 2bn_3$ is satisfied.

Regarding the magnitude of A_3 , BBPF commented that, for de-excitation via the transition $2s_{1/2}-1s_{1/2}$ to influence β , it is necessary to have

$$A_3 > A_1. \quad (11)$$

They argued that this inequality could not be satisfied in an optically thin plasma, where $A_1 = A$, but that it could conceivably be satisfied in an optically thick plasma. BBPF derived a theoretical value for A_3 of

$$A_3 \sim \frac{1}{2}A \quad \text{or} \quad \frac{1}{3}A, \quad (12)$$

which is sufficiently large to suggest that equation (11) might indeed be satisfied in an optically thick plasma.

The above interpretation is now subject to the following criticism. Firstly, we note that theoretical expressions for A_1 and A_2 (such as in equations 3), when substituted into equation (2), provide an expression for β , provided ε is known. The basic reason for using the rate equations to provide an expression for β is that ε is not known. Yet this very uncertainty in ε is sufficient to negate the use of the rate equations for this purpose (see condition (iii) in Section 2b(ii)), and consequently is sufficient to negate the validity of equation (10). In a laser-produced plasma the presence of differential motion has a similar effect, rendering equation (10) invalid (see condition (iv) in Section 2b(ii)), and this is likely to be more significant than the uncertainty in ε .

Secondly, the value deduced above for $\tau_{0,1}$ is likely to be erroneous (apart from the fact that equation 10 is invalid) since the theoretical expressions for A_1 and A_2 in equations (3) allow neither for overlap nor differential expansion. However, this is not an important issue because the value deduced for $\tau_{0,1}$ (in the range 3–8) is not critical to the main argument. BBPF comment only that they expect the optical depth to exceed unity.

Thirdly, equation (11) (namely, $A_3 > AA_1$) is less likely to be satisfied in a differentially expanding plasma than in a stationary plasma (all other parameters being equal), since the presence of differential expansion increases the value of A_1 (see Fig. 7 of Paper 1).

Fourthly, there is no independent evidence for the presence of the $2s_{1/2}-1s_{1/2}$ transition. Forbidden lines have been observed in laser-produced plasmas (see the review by Bekefi *et al.* 1976), and the observation of the forbidden $3s-2s$ transition in the lithium-like ion FVII by Dewhurst (1974) may prove to be significant (see the discussion by Vinogradov and Yukov 1975). Unfortunately, the $2s_{1/2}-1s_{1/2}$ transition in the hydrogen-like L_α doublet virtually coincides in wavelength with $2p_{1/2}-1s_{1/2}$, and so cannot be observed directly. The presence of this transition can only be deduced indirectly from, for example, observations of anomalous β values. A laser-produced plasma is not a good choice of source in this respect, since the differential motion inherent in such a plasma can itself produce anomalous β values (see Section 4). A better choice of a high density source might be the vacuum spark or exploding wire plasma, though here again differential motion may prove to be an issue. Observations of the fine-structure components of the hydrogen-like L_α doublet in the vacuum spark (Lie and Elton 1971; Kononov 1978), the exploding wire plasma (Burkhalter *et al.* 1977), a gas puff Z-pinch device (Burkhalter *et al.* 1979) and also in the low

density solar corona (Grineva *et al.* 1973) all show values of $\beta < 1$ for the ions Mg XII, A XVIII, Ti XXII and Fe XXVI. Up to now, values of $\beta > 1$ appear to be unique to the laser-produced plasma and to the microwave-powered source mentioned in Section 1.

We conclude that the assumed presence of the single-photon transition $2s_{1/2}-1s_{1/2}$ must be regarded as tentative and awaiting confirmation. If this transition is present then it may well contribute to the anomalous β values. However, by neglecting the effect of overlap and differential expansion, BBPF have failed to demonstrate the qualitative agreement between theory and experiment which they claim.

3. An Alternative Interpretation

In Section 2 we reviewed the observations of anomalous β values in optically thick laser-produced plasmas, and discussed the interpretation by BBPF. An alternative interpretation is now offered, based on the capacity of differential plasma motion to strongly influence the transport of radiation (see e.g. Athay 1972).

Consider a line of sight which typically is transverse to the axis of the expanding laser-produced plasma (see Fig. 1 of Irons 1975). Along this line of sight the component of expansion velocity towards the observer varies continuously from some value $+V$ at the near plasma surface through zero on the axis to $-V$ at the far surface. Observations on optically thin line profiles, Doppler broadened by differential plasma motion, have shown values of $V \sim 10^7 \text{ cm s}^{-1}$ (Irons *et al.* 1972; Irons 1973), which are several times greater than mean thermal velocities. The plasma located between any emitting ion and the observer has a higher component of velocity towards the observer than does the emitting ion, and so absorbs preferentially on the blue side of the line. Observations of asymmetries in optically thick line profiles in laser-produced plasmas provide striking evidence of this effect (see e.g. Valero *et al.* 1969; Galanti *et al.* 1974; Nicolosi *et al.* 1978; Irons 1975, 1980a).

Clearly, in the case of a doublet the absorption will be reinforced when photons emitted in the blue peak of the doublet find, in their passage through the outer region of the expanding plasma, that they coincide with the local frequency of the red peak. Anomalous intensity ratios arising in this way in expanding stellar atmospheres have been previously discussed by Hewitt and Noerdlinger (1974), who have commented that: 'There are many astrophysical objects in which P Cygni-type profiles of a composite type are produced due to doublet structure, or line "blending"; in other words, there are two resonance lines or two components of one line so arranged that the expected rest position of the blueward line falls within the absorption feature of the redward line.' However, to the author's knowledge, this particular effect has not been previously discussed in relation to a laboratory plasma, though we note the closely related effect discussed by Braun *et al.* (1970) and described here in Section 1. It is now claimed that preferential absorption of the blue peak of the hydrogen-like L_α doublet, arising from overlap with the red peak as outlined above, can account for the values of $\beta \approx 2$ observed in laser-produced plasmas.

Unlike the interpretation by BBPF, the precise value of ε is no longer seen as an issue, at least not for the semiquantitative agreement between theory and experiment which we intend to demonstrate. Hence we do not concern ourselves with rate equations. The important issue in the present interpretation is the modification to the line profile during transmission through the plasma, along the line of sight towards the observer; implying that the line-of-sight properties of the plasma, particularly

the velocity field, are the parameters of prime importance. If Ξ_1 (Ξ_2) denotes the average probability that a transition 1 (transition 2) photon emitted towards the observer along the line of sight escapes from the plasma, then we have (equations 2 with $A_1 = A\Xi_1$ and $A_2 = A\Xi_2$)

$$\beta = (\Xi_2/\Xi_1)\epsilon.$$

The model calculations in Paper 1 provide us with values of Ξ_1 and Ξ_2 in a differentially expanding plasma which permit a comparison with experiment to be made, as discussed below.

4. Model Calculations: Comparison with Experiment

The following model calculations serve two purposes. Firstly, they illustrate that, for conditions typical of Mg XII L_α in a differentially expanding laser-produced plasma, β can attain a value of 2 (Section 4a). Secondly, they reveal certain structure within the spectral profile of the emergent radiation which is not apparent from the qualitative discussion in the preceding sections but which, nevertheless, is of the greatest importance to the present interpretation of the anomalous β values (Section 4b).

(a) The Ratio $\beta = (\Xi_2/\Xi_1)\epsilon$

(i) Model Calculations

For details of the following calculations the reader is referred to Section 4 of Paper 1. Briefly, the spacing ζ of the hydrogen-like L_α doublet is $\zeta = 0.00541 \text{ \AA}$, independently of atomic number Z to a first approximation (Garcia and Mack 1965). Because ζ is independent of Z when expressed in wavelength units, so it is convenient to express the following argument in wavelength (rather than frequency) units. We assume that Stark broadening is negligible, and we denote the HWHM of each transition due to thermal Doppler broadening as $\lambda_\frac{1}{2}$. We set

$$\zeta/\lambda_\frac{1}{2} = 3. \quad (13)$$

This implies $\lambda_\frac{1}{2} = 0.0018 \text{ \AA}$, which (for Mg XII L_α) corresponds to a temperature of 740 eV, which is characteristic of Mg XII close to the target surface in a laser-produced plasma (Bayanov *et al.* 1976; BBPF). Also, close to the target, we have a plasma depth of typically 10^{-2} cm and a ground state density, for Mg XII, of $\sim N_e/3(Z-1)$, which becomes $\sim 3 \times 10^{19} \text{ cm}^{-3}$ when $N_e = 10^{21} \text{ cm}^{-3}$ and $Z = 12$. The above values give an optical depth for the plasma, when at rest, of $\tau_{0,1} \approx 14$ for transition 1 at line centre and of ~ 7 ($= \frac{1}{2}\tau_{0,1}$) for transition 2 at line centre.

The one-dimensional spatial model employed in Paper 1 assumes a constant ground state density, an upper state density which decreases steadily towards the surface of the source, and a motional velocity which increases linearly with distance from the centre of the source, to a value V at the surface.* Corresponding to the velocity V , there is a wavelength shift $\lambda_0 V/c$ (where λ_0 denotes the wavelength of the usual line centre) or, in units of $\lambda_\frac{1}{2}$, a shift of

$$\eta = \lambda_0 V/c\lambda_\frac{1}{2}.$$

* The limited available evidence (see Section 2 of Irons 1975) indicates that the component of motional velocity transverse to the expansion axis of a laser-produced plasma increases approximately linearly with distance from the axis.

For a one-dimensional source we have $\Xi_1 = \Theta_1$ and $\Xi_2 = \Theta_2$, and the plot of Θ_2/Θ_1 for $\tau_{0,1} = 10$ in Fig. 6 of Paper 1 is reproduced here (in Fig. 1) as a plot of Ξ_2/Ξ_1 , with η as variable. Additional results are presented for $\tau_{0,1} = 20$. We see that Ξ_2/Ξ_1 increases with η , attaining a maximum value at $\eta \approx 3$, and then decreases towards unity. The increase in Ξ_2/Ξ_1 is a consequence of the increasing overlap of transition 1 by transition 2 as η increases (as discussed in Paper 1). The maximum occurs when

$$\eta \approx \zeta/\lambda_{\frac{1}{2}}, \quad (14)$$

that is, approximately when the central wavelength of transition 2 at the surface of the source overlaps with the central wavelength of transition 1 at the centre of the source. For increasingly large values of η the plasma opacity decreases and both Ξ_1 and Ξ_2 approach the optically thin value of unity, as also does the ratio Ξ_2/Ξ_1 .

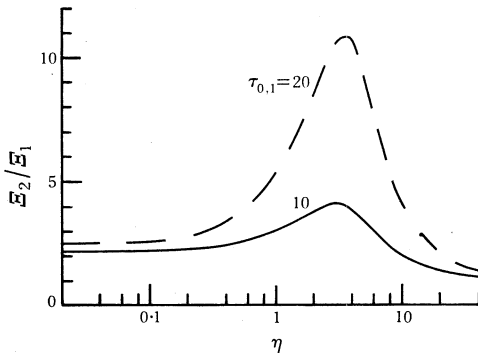


Fig. 1. Model calculations of the ratio Ξ_2/Ξ_1 for the L_{α} doublet with $\zeta/\lambda_{\frac{1}{2}} = 3$, plotted as a function of the wavelength shift η for $\tau_{0,1} = 10$ and 20.

(ii) Comparison with Experiment

As commented upon previously (Irons 1980a), in laser-produced plasmas we may expect V to be of the order of, and perhaps several times greater than, the mean thermal velocity; implying values of η for Doppler-broadened lines (i.e. for lines with negligible Stark broadening) in the approximate range 1–3. Such values are required, for example, to explain the appearance of grossly asymmetrical lines in laser-produced plasmas (Irons 1980a). Hence (see Fig. 1) we may well expect values of $\Xi_2/\Xi_1 \approx 4$ for $\tau_{0,1}$ in the range 10–20; in which case, even for ε as small as 0.5, β may attain a value of ~ 2 , comparable with what is observed. Since ε is likely to exceed 0.5,* so the probability of β attaining a value of 2 is that much greater. Note also that Ξ_2/Ξ_1 is quite sensitive to the value of $\tau_{0,1}$ for η in the range 1–3 (and higher).

With increasing distance from the target surface in a laser-produced plasma, the plasma temperature decreases and η consequently increases. If this were the only effect operating then we would expect the variation of Ξ_2/Ξ_1 with increasing η in Fig. 1 to reflect the variation of β with increasing distance from the target. For example, if we have $\eta \approx 1$ near the target surface then we would expect an increase in β with distance from the target. However, with increasing distance from the target,

* As discussed in Section 2, the interplay of collisional and radiative transitions causes ε to exceed 0.5. The effect of opacity, ignoring overlap, is to reduce the effectiveness of radiative de-excitation compared with collisional de-excitation, thereby reducing the tendency for ε to exceed 0.5. The effect of overlap is that the weaker transition (transition 2) is pumped at the expense of the stronger transition, causing ε to increase; and this is further enhanced by differential expansion.

the plasma opacity also decreases and β must eventually decrease to the appropriate optically thin value. The net variation of β with distance therefore depends on how these two effects combine, and may differ for different plasmas (cascade population of quantum state $n = 2$ may also become an issue with increasing distance from the target; see Section 2a). It is not surprising therefore that Bayanov *et al.* (1976) should report β as initially increasing with distance from the target and then decreasing, whereas BBPF report β as having a maximum value at the target surface and then decreasing.

It is of interest to speculate on the variation of β with Z . Firstly, we note that ζ is practically independent of Z and that λ_0 varies as Z^{-2} , so that the ratio $\zeta/\lambda_{\frac{1}{2}}$ varies as $Z^{5/2}T^{-1/2}$, where T denotes the ion temperature (assumed equal to the electron temperature in the following comments). The temperature might be expected to be approximately constant with Z in experiments conducted with constant laser parameters (e.g. Tonon 1972; Boiko *et al.* 1977). However, if we impose a constraint, such that the hydrogen-like species must be the predominant ion species present, then we would expect T to increase with Z . (Some such constraint is presumably necessary if we are to assume that all hydrogen-like ions are characterized by η in the approximate range 1–3.) If we take T as being proportional to the ionization potential, i.e. to Z^2 (see e.g. Boiko *et al.* 1978*b*, 1979), then $\zeta/\lambda_{\frac{1}{2}}$ varies as $Z^{3/2}$. This implies that η , with values in the range 1–3, drops further below the optimum value (of $\zeta/\lambda_{\frac{1}{2}}$) for anomalous β values with increasing Z . With the same assumption regarding T , and with the ground state density varying as Z^{-1} (see above), we find that $\tau_{0,1}$ varies as $Z^{-7/2}$, that is, it decreases rapidly with increasing Z . For both these reasons we are less likely to observe anomalous β values with increasing Z . For late second sequence elements the fine-structure components may still be resolvable when $N_e = 10^{21} \text{ cm}^{-3}$ (see e.g. Weisheit and Rozsnyai 1976). However, with decreasing Z , Stark broadening becomes dominant and the fine-structure components are no longer observable (see e.g. the profiles of Be IV L_α –O VIII L_α reported by Nicolosi *et al.* 1978). It is no coincidence that anomalous β values should have been reported for magnesium and aluminium, for laser-produced plasmas from plane targets (our expectations would be different, for example, for the imploding microballoon situation).

In addition to the Mg XII L_α and Al XIII L_α results, it would be interesting to have space-resolved observations of Na XI L_α and Si XIV L_α in plasmas produced from targets of pure sodium and silicon. The profile of Si XIV L_α recorded by Boiko *et al.* (1977), apparently without spatial resolution, indicates a value of $\beta \sim 1$, which is in excess of $\beta_{\text{max}} (= 0.73)$. Given that the optical depth varies as $Z^{-7/2}$, and given that $\tau_{0,1} = 14$ for Mg XII L_α , we obtain $\tau_{0,1} = 6.5$ for P XV L_α , close to the target surface. Boiko *et al.* (1978*b*, 1979) have interpreted the β ratios of P XV L_α from a laser-produced phosphorus plasma, assumed optically thin, to give values of N_e which are in good agreement with other measurements (except far from the target where $\beta > \beta_{\text{max}}$, but this does not appear to be an opacity-related problem). This agreement is somewhat surprising given the above estimate of $\tau_{0,1} = 6.5$ for P XV L_α . Space-resolved observations of optically thick doublets of other than hydrogen-like ions would also be of interest.

Except for the above comments regarding second sequence elements, we have not yet considered to what extent the presence of overlap means that we cannot deduce β from an observation of the doublet line profile. This, and other issues associated with the doublet line profile, are considered in the following subsection.

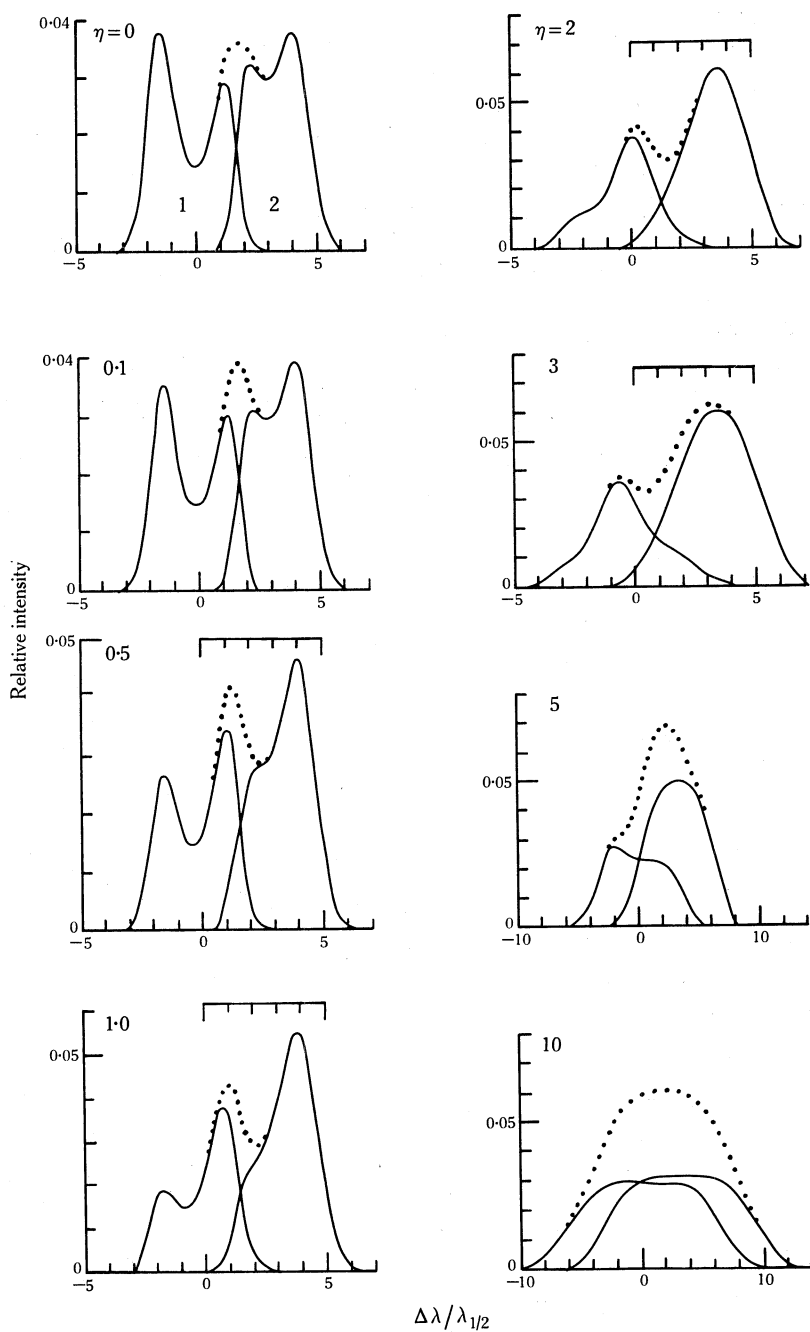


Fig. 2. Model calculations of the emergent profile of transition 1 ($2p_{3/2}-1s_{1/2}$, the lower wavelength component) and transition 2 ($2p_{1/2}-1s_{1/2}$), plus the composite profile (dotted curve), for the L_x doublet with $\zeta/\lambda_{\pm} = 3$ and $\epsilon = 0.5$, plotted as a function of normalized wavelength for $\tau_{0,1} = 10$. The figure shows how the profiles vary with η in the range 0-10 as indicated.

(b) Line profiles

(i) Model Calculations

Fig. 2 shows the spectral profile of emergent radiation, for transition 1 and transition 2 and for the two transitions in combination, computed for the same conditions as described in Section 4a. The areas under the transition 1 and transition 2 profiles divided by the corresponding optically thin values are equal to Ξ_1 and Ξ_2 respectively. In presenting data for both transitions on the one graph it is necessary to make some assumption regarding relative intensities, i.e. regarding the value of ε . In Fig. 2 we have assumed $\varepsilon = 0.5$.

Fig. 3 of Paper 1 shows how the emergent profile varies with $\tau_{0,1}$ for the case of a stationary source ($\eta = 0$). The present Fig. 2 shows how the emergent profile varies with η , for $\tau_{0,1} = 10$. For $\eta = 0$ the profile of each transition is self-reversed, with the adjacent sides of each transition (i.e. the red side of transition 1 and the blue side of transition 2) depressed in magnitude because of overlap, as discussed in Paper 1. With increasing η the blue side of each transition is depressed and the red side enhanced, in accordance with the behaviour of isolated lines, as discussed elsewhere (Irons 1975, 1980a). When $\eta = 1$, the blue peak of transition 2 is wholly depressed, whereas the blue peak of transition 1 is still discernible. For $\eta \lesssim 1$ the doublet appears as three peaks, the red peak of transition 1 and the blue peak of transition 2 having blended into a single central peak. For $\eta \gtrsim 2$ the blue peak of transition 1 is depressed to a point where it is hardly discernible. The doublet now appears as two peaks. For $\eta \approx 2$ the relative Doppler shift between different parts of the source is such that photons emitted in transition 1 on one side of the source may be strongly absorbed in transition 2 on the other side of the source. Transition 1 is now depressed in magnitude relative to transition 2. This effect continues to be present for higher values of η . For $\eta \gtrsim 5$ motional Doppler broadening has become so dominant that the two transitions have merged together, and the doublet now appears as one peak. For $\eta \gtrsim 10$ the shape of each profile becomes strongly dependent on the spatial distributions of emitting and absorbing atoms, as discussed by Irons (1975).

(ii) Comparison with Experiment

(1) Short-wave satellite

For laser-produced plasmas from plane targets, Bayanov *et al.* (1976) and BBPF have reported profiles of Mg XII L_α and Al XIII L_α with two peaks, close to the target surface where these lines are optically thick (see Section 2b(i)). According to the present theory, as represented by the model calculations in Fig. 2, this can only occur when η is of order 2. Indeed, only when η is of this order can β be measured (for $\eta \lesssim 1$ it would be virtually impossible to analyse the central peak of the doublet into the component transitions with any precision, while for $\eta \gtrsim 5$ the doublet appears as only one peak).

For laser-produced plasmas from conical depressions, Bayanov *et al.* (1976) have reported profiles of Mg XII L_α with three peaks close to the target surface. They have labelled the high wavelength peak as transition 2, the central peak as transition 1 and the low wavelength peak as a satellite. According to the present theory, however, the three peaks represent a situation where $\eta \lesssim 1$. This is consistent with the limited evidence (for references see Irons 1980a) that laser-produced plasmas from conical

depressions have transverse velocities approximately three times smaller than laser-produced plasmas from plane targets. Short wavelength satellites have also been reported for the resonance lines of helium-like ions in laser-produced plasmas from conical depressions, and have been similarly explained in terms of a reduced transverse velocity (Irons 1980a). Note that the model for the spatial distribution of emitting and absorbing ions employed here is different from that employed previously (by Irons 1975, 1980a). This serves to illustrate that the asymmetric self-reversal, with the appearance of a resonance line plus satellite, is not model dependent. As commented previously (Irons 1980a), if the satellite is not acknowledged as being part of the line proper, then measurements of plasma density and temperature employing the resonance line intensity are likely to be in error.

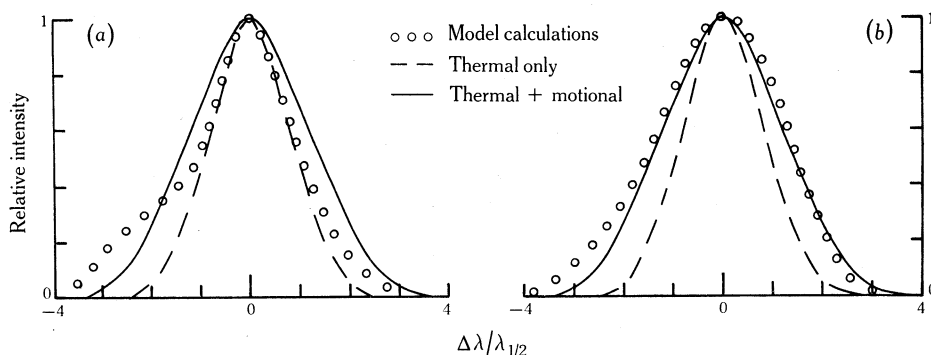


Fig. 3. Model calculations from Fig. 2 for $\eta = 2$ matched against optically thin profiles for thermal Doppler broadening only ($\eta = 0$) and thermal plus motional Doppler broadening ($\eta = 2$): (a) transition 1; (b) transition 2. The profiles are assigned equal peak intensity, and $\Delta\lambda$ is here measured from the centre of each set of matched profiles.

With increasing distance from the target the plasma temperature decreases and η correspondingly increases. If this were the only effect operating then we could expect the variation of the emergent line profile with increasing η in Fig. 2 to reflect the variation with distance from the target. If, for example, we had $\eta \approx 1$ near the target surface then we could expect to see a transition from a three-peaked to a two-peaked profile with increasing distance. However, with increasing distance from the target, the plasma opacity also decreases and the doublet must eventually appear as the sum of two optically thin transitions. The net variation with distance therefore depends on how the above two effects combine. Basically, a profile which initially is two-peaked (as from a plane target) might be expected to retain a two-peaked structure with increasing distance; whereas a profile which initially is three-peaked (as from a conical depression) might be expected to attain a two-peaked structure with increasing distance. These expectations are consistent with the observations of Bayanov *et al.* (1976) and BBPF.

When the concentration of the element concerned is decreased (as with the use of composite targets), the plasma opacity decreases and the optically thin profile is approached. Again, if the profile is initially two-peaked we may expect it to retain a two-peaked structure with decreasing concentration, as reported for $\text{Mg XII } L_\alpha$ by BBPF. Hence the various properties reported for the $\text{Mg XII } L_\alpha$ and $\text{Al XIII } L_\alpha$ doublets by Bayanov *et al.* (1976) and BBPF can be explained once allowance is made for the effect of differential motion on the plasma opacity.

(2) Similarity between optically thick and optically thin profiles

Perhaps the single most interesting point to emerge from Fig. 2 is that the optically thick profile of transition 1 and transition 2 can take a shape (for $\eta = 2$) which is deceptively similar to the optically thin profile. This is illustrated in Fig. 3, which shows the transition 1 and transition 2 profiles for $\tau_{0,1} = 10$ and $\eta = 2$ matched against the corresponding optically thin profile. Included for comparison is the optically thin profile for thermal Doppler broadening only, i.e. for $\eta = 0$. If profiles similar to the optically thick profiles in Fig. 3 were observed and interpreted as being optically thin, they would clearly lead to plausible estimates of ion velocity. This, in turn, would strengthen the impression that optically thin profiles were being observed. Comments to this effect have been made previously (Irons 1980a) in connection with the resonance lines of helium-like ions. The fact that an L_α profile may be approximated as the sum of two gaussian profiles (Boiko *et al.* 1978b, 1979) is no guarantee of optical transparency.

(3) Component spacing

In Fig. 2 the spacing between the two component transitions for $\eta = 2$, namely $3.4\lambda_{\frac{1}{2}} (= 0.0061 \text{ \AA})$,* is greater than the optically thin value of $0.00541 \text{ \AA} (= 3\lambda_{\frac{1}{2}}$, see equation 13). This spacing increases to $3.8\lambda_{\frac{1}{2}} (= 0.0069 \text{ \AA})$ for $\eta = 3$, and decreases to $2.8\lambda_{\frac{1}{2}} (= 0.0050 \text{ \AA})$ for the two stronger peaks at $\eta = 1$. These values will vary somewhat with the assumed value of ϵ . Broadly speaking, close to the target surface in laser-produced plasmas from plane targets (where the plasma is optically thick and η is of order 2), we may expect to observe spacings in excess of the theoretical optically thin value of 0.0054 \AA . Unfortunately, neither Bayanov *et al.* (1976) nor BBPF have reported experimental values for this spacing. In spectra recorded without spatial resolution the probability of observing spacings in excess of 0.0054 \AA is diminished, because of the contribution of radiation from optically thin plasma far (i.e. beyond a few $100 \mu\text{m}$) from the target surface. Aglitskii *et al.* (1974; see also Boiko *et al.* 1977, 1978a) have reported spacings of 0.0059 \AA (for Mg XII L_α), 0.0056 \AA (Al XIII L_α), 0.0061 \AA (Si XIV L_α), 0.0057 \AA (P XV L_α) and 0.0060 \AA (S XVI L_α), in spectra recorded without spatial resolution. Interestingly enough these spacings *do* indicate a systematic departure from the optically thin value consistent with the above expectation.†

It is worth emphasizing that the theoretical optically thin value of the L_α spacing ζ is known with great precision. To a good first approximation, $\zeta = 0.00539 \text{ \AA}$ independently of Z (Edlén 1964). The more precise calculations of Garcia and Mack (1965), for H I–Ca XX, indicate that ζ varies from a minimum of 0.0054049 \AA for ${}^7\text{Li III}$ to a maximum of 0.005417 \AA for ${}^{40}\text{Ca XX}$. The even greater accuracy obtainable from the energy level tables of Erickson (1977) need not concern us here. The energy spacing $2p_{3/2} - 2p_{1/2}$ has been measured for hydrogen and deuterium (but not for higher hydrogen-like ions; see Erickson 1977). The results (obtained by the

* The values quoted in the text represent the spacings between the peaks of the doublet; the spacings between the peaks of the individual transitions are approximately 8% greater than these.

† Aglitskii *et al.* (1974) do not quote errors for their measurements of the doublet spacing. However, they do quote errors for absolute wavelength measurements of each transition, of $\pm 0.0015 \text{ \AA}$ for Mg XII L_α , decreasing to $\pm 0.0010 \text{ \AA}$ for S XVI L_α . We may expect the errors in the spacing to be very much smaller.

level-crossing technique), namely 10969·13 MHz for hydrogen (Baird *et al.* 1972) and 10971·59 MHz for deuterium (Dayhoff *et al.* 1953), correspond to $\zeta = 0\cdot0054073$ and $0\cdot0054056$ Å respectively, in agreement with the theoretical values of Garcia and Mack (1965), namely $\zeta = 0\cdot0054071$ and $0\cdot0054056$ Å. There is no reason to doubt the accuracy of the theoretical calculations of ζ for higher Z elements. However, the only reported measurements of ζ for high Z elements appear to be those of Aglitskii *et al.* (1974) and Boiko *et al.* (1977, 1978a) described above; additional reports would obviously be desirable. It would be of interest to have, for example, further observations on laser-produced plasmas, where (1) the spacing is observed as a function of distance from the target and (2) the element in question is present as a trace impurity, so that the plasma is optically thin to the L_α transition concerned, even at the target surface.

The B_α transition is basically a doublet, with a spacing of $0\cdot142$ Å ignoring the Lamb shift and slightly smaller ($\sim 0\cdot138$ Å) with the Lamb shift included (Erickson 1977). This spacing is determined largely by the energy spacing $2p_{3/2}-2p_{1/2}$. If this latter energy spacing for the ions Mg XII–SXVI were actually 10% greater than the accepted theoretical value (this is one interpretation of the difference between the experimental ($\sim 0\cdot0059$ Å) and theoretical ($0\cdot0054$ Å) values of the L_α doublet spacing), then we would expect to observe the B_α doublet with a spacing $\sim 10\%$ greater than $0\cdot14$ Å. This is certainly not the case with the B_α transitions of hydrogen and deuterium which, in the past, have been the subject of extensive study (see e.g. Series 1957), nor is it the case with the B_α transitions of C VI–O VIII which have recently been observed as doublets in laser-produced plasmas under optically thin conditions (Nicolosi *et al.* 1979). (Line profile data kindly supplied to the author show a B_α doublet spacing of $0\cdot125 \pm 0\cdot010$ Å when averaged over the ions C VI–O VIII and over points in the region $0\cdot2$ – $1\cdot4$ mm from the target.)

Given that the difference of $\sim 0\cdot0005$ Å between the experimental and theoretical values of the L_α doublet of Mg XII–SXVI is real (and cannot be explained in terms of experimental errors), then we conclude that the absolute wavelengths measured by Aglitskii *et al.* (1974) and Boiko *et al.* (1977, 1978a) for (at least one of) the component transitions are in error by an amount comparable with $0\cdot0005$ Å. While this is smaller than the quoted experimental uncertainty of $\pm 0\cdot0015$ Å for Mg XII L_α (decreasing to $\pm 0\cdot0010$ Å for SXVI L_α), it nevertheless represents a significant source of error for which no allowance has so far been made. The error would be greater if the wavelengths were measured close to the target surface in spatially resolved spectra.

5. Conclusions

The anomalous intensity ratios exhibited by the fine-structure components of Mg XII L_α and Al XIII L_α in optically thick laser-produced plasmas can be explained in terms of the usual radiative and collisional processes, provided allowance is made for the effect of differential motion on the plasma opacity. Model calculations give intensity ratios in good agreement with experiment, and explain such details of the line profile structure as the short-wave satellite and a component spacing in excess of the theoretical optically thin value. The model calculations also show that the anomaly in the intensity ratio is greatest when the condition in equation (14) is satisfied, as is approximately the case with Mg XII L_α and Al XIII L_α in laser-produced plasmas from plane targets. Our ability to observe anomalous intensity ratios is also

dependent on equation (14) being approximately satisfied, as discussed in Section 4a. Suggestions for further experiments are contained in Sections 4a and 4b. Wavelengths measured for the optically thick L_α doublet of Mg XII–SXVI may be expected to differ from the optically thin values by an amount comparable with 0.0005 \AA .

The analysis in this, and a previous paper (Irons 1980a), clearly demonstrates that, in a laser-produced plasma, differential motion is a major factor in determining the profile and intensity of optically thick Doppler-broadened lines. Much more attention needs to be given to this factor than in the past, in all areas involving radiation transport. This includes the interpretation of line profiles and intensities for plasma diagnostic purposes, the calculation of excitation, ionization and recombination rate coefficients, and the calculation of gain for a given population inversion.

An important point to emerge from this, and the previous paper (Irons 1980a) is that the profile of an optically thick line may be distorted by differential plasma motion in such a way that it may be deceptively similar to the optically thin profile, and may not be recognized as being optically thick.

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References

- Aglitskii, E. V., Boiko, V. A., Zakharov, S. M., Pikuz, S. A., and Faenov, A. Ya. (1974). *Sov. J. Quantum Electron.* **4**, 500–13; English transl. of *Kvantovaya Elektron. (Moscow)* **1**, 908–36.
- Athay, R. G. (1972). 'Radiation Transport in Spectral Lines' (Reidel: Dordrecht).
- Baird, J. C., Brandenberger, J., Gondaira, K., and Metcalf, H. (1972). *Phys. Rev. A* **5**, 564–87.
- Bayanov, V. I., et al. (1976). *Sov. J. Quantum Electron.* **6**, 1226–33; English transl. of *Kvantovaya Elektron. (Moscow)* **3**, 2253–65.
- Beigman, I. L., Boiko, V. A., Pikuz, S. A., and Faenov, A. Ya. (1976). *Sov. Phys. JETP* **44**, 511–15; English transl. of *Zh. Eksp. Teor. Fiz.* **71**, 975–83.
- Bekefi, G., Deutsch, C., and Yaakobi, B. (1976). In 'Principles of Laser Plasmas' (Ed. G. Bekefi), Ch. 13 (Wiley: New York).
- Belyaev, L. M., et al. (1977). *Sov. J. Quantum Electron.* **7**, 67–70; English transl. of *Kvantovaya Elektron. (Moscow)* **4**, 129–35.
- Boiko, V. A., Faenov, A. Ya., and Pikuz, S. A. (1977). *Mon. Not. R. Astron. Soc.* **181**, 107–20.
- Boiko, V. A., Faenov, A. Ya., and Pikuz, S. A. (1978a). *J. Quant. Spectrosc. Radiat. Transfer* **19**, 11–50.
- Boiko, V. A., Pikuz, S. A., and Faenov, A. Ya. (1979). *J. Phys. B* **12**, 1889–1910.
- Boiko, V. A., Vinogradov, A. V., Pikuz, S. A., Skobelev, I. Yu., Faenov, A. Ya., and Yukov, E. A. (1978b). *Sov. J. Plasma Phys.* **4**, 54–8; English transl. of *Fiz. Plazmy* **4**, 97–103.
- Braun, W., Bass, A. M., and Davis, D. D. (1970). *J. Opt. Soc. Am.* **60**, 166–70.
- Burkhalter, P. G., Dozier, C. M., and Nagel, D. G. (1977). *Phys. Rev. A* **15**, 700–17.
- Burkhalter, P. G., Shiloh, J., Fisher, A., and Cowan, R. D. (1979). *J. Appl. Phys.* **50**, 4532–40.
- Dayhoff, E. S., Triebwasser, S., and Lamb, W. E. (1953). *Phys. Rev.* **89**, 106–15.
- Dewhurst, R. J. (1974). *Opt. Commun.* **12**, 60–2.
- Edlén, B. (1964). In 'Handbuch der Physik', Vol. 27 (Ed. S. Flügge), pp. 80–220 (Springer: Berlin).
- Edlén, B. (1966). *Ark. Fys.* **31**, 509–10.
- Erickson, G. W. (1977). *J. Phys. Chem. Ref. Data* **6**, 831–69.
- Galanti, M., Peacock, N. J., Norton, B. A., and Puric, J. (1974). Proc. 5th Int. Conf. on Plasma Physics and Controlled Thermonuclear Reactions, Tokyo, Paper No. CN-33/F3-4 (IAEA: Vienna).
- Garcia, J. D., and Mack, J. E. (1965). *J. Opt. Soc. Am.* **55**, 654–85.

- Grineva, Yu. I., *et al.* (1973). *Solar Phys.* **29**, 441–6.
- Hewitt, T. G., and Noerdlinger P. D. (1974). *Astrophys. J.* **188**, 315–26.
- Holstein, T. (1947). *Phys. Rev.* **72**, 1212–33.
- Holstein, T. (1951). *Phys. Rev.* **83**, 1159–68.
- Hutcheon, R. J., and McWhirter, R. W. P. (1973). *J. Phys. B* **6**, 2668–83.
- Irons, F. E. (1973). *J. Phys. B* **6**, 1562–81.
- Irons, F. E. (1975). *J. Phys. B* **8**, 3044–68.
- Irons, F. E. (1976). *J. Phys. B* **9**, 2737–53.
- Irons, F. E. (1979a). *J. Quant. Spectrosc. Radiat. Transfer* **22**, 1–20.
- Irons, F. E. (1979b). *J. Quant. Spectrosc. Radiat. Transfer* **22**, 37–44.
- Irons, F. E. (1980a). *Aust. J. Phys.* **33**, 25–38.
- Irons, F. E. (1980b). The escape factor in plasma spectroscopy. IV. The case of an overlapping doublet. *J. Quant. Spectrosc. Radiat. Transfer* (in press).
- Irons, F. E., McWhirter, R. W. P., and Peacock, N. J. (1972). *J. Phys. B* **5**, 1975–87.
- Irons, F. E., and Peacock, N. J. (1974). *J. Phys. B* **7**, 2084–99.
- Kononov, E. Ya. (1978). *Phys. Scr.* **17**, 425–32.
- Kononov, E. Ya., Koshelev, K. N., and Ryabtsev, A. N. (1971). *Opt. Spectrosc.* **30**, 534–6.
- Kugel, H. W., and Murnick, D. E. (1977). *Rep. Prog. Phys.* **40**, 297–343.
- Lie, T. N., and Elton, R. C. (1971). *Phys. Rev. A* **3**, 865–71.
- Marrus, R. (1977). *Commun. Atom. Mol. Phys.* **7**, 53–8.
- Marrus, R., and Mohr, P. J. (1978). *Adv. Atom Mol. Phys.* **14**, 181–224.
- Nicolosi, P., Garifo, L., Jannitti, E., Malvezzi, A. M., and Tondello, G. (1978). *Nuovo Cimento B* **48**, 133–51.
- Nicolosi, P., Tondello, G., and Jannitti, E. (1979). *Phys. Rev. A* **20**, 2574–81.
- Series, G. W. (1957). 'Spectrum of Atomic Hydrogen' (Oxford Univ. Press).
- Skobelev, I. Yu., Vinogradov, A. V., and Yukov, E. A. (1978). *Phys. Scr.* **18**, 78–86.
- Tonon, G. F. (1972). *IEEE Trans. Nucl. Sci.* NS **19**, 172–83.
- Valero, F. P. J., Goorvitch, D., Fraenkel, B. S., and Ragent, B. (1969). *J. Opt. Soc. Am.* **59**, 1380–1.
- Vinogradov, A. V., Skobelev, I. Yu., and Yukov, E. A. (1977). *Sov. J. Plasma Phys.* **3**, 389–94; English transl. of *Fiz. Plazmy* **3**, 686–96.
- Vinogradov, A. V., Sobel'man, I. I., and Yukov, E. A. (1974). *Sov. J. Quantum Electron.* **4**, 149–54; English transl. of *Kvantovaya Elektron. (Moscow)* **1**, 268–78.
- Vinogradov, A. V., and Yukov, E. A. (1975). *Sov. J. Plasma Phys.* **1**, 472–3; English transl. of *Fiz. Plazmy* **1**, 860–2.
- Weisheit, J. C., and Rozsnyai, B. P. (1976). *J. Phys. B* **9**, L63–7.
- Yaakobi, B., *et al.* (1979). *Phys. Rev. A* **19**, 1247–62.

