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Analysis of the coupling between impact and quasistatic field mechanisms in Stark broadening

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Abstract

In this work, the scope of the hypothesis of the physical and statistical independence between ionic (slow variation) and electronic (fast variation) microfields applied to spectral lines calculation in plasmas is analyzed. This hypothesis is used in some spectral profiles calculation methods. This analysis has been done by computer simulation.

Keywords: Plasma diagnosis; Spectroscopy; Stark effect; Computer simulation

1. Introduction

Computer simulation techniques have been successfully applied from the beginning of the 1980s to the Stark broadening calculations [1–8]. These techniques offer an excellent agreement with the experimental data, especially in neutral emitter spectrum. However, up to now the used methods have the serious inconvenience of being very costly in calculation (many hundreds of configurations are needed to obtain a good spectrum and each one of them requires a considerable volume of calculation, especially if we are considering conditions of low density or very high temperature). Because of this, some fast methods to calculate spectral profiles have been proposed lately. Though these techniques simplify the physical and the mathematical problem, they can offer results with sufficient precision. On the other hand, these methods use the simulations to test the quality of the results they offer [9].

Examples of these methods are the 'frequency separation method' [10], 'frequency fluctuation model' (FFM) [11,12], and the PIM PAM POUM method [9,13]. The former consider that the memory loss at the emission process that characterizes the line broadening can be treated as

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a relaxation problem in which appear two superposed contributions: the one of the electrons, with fields of a very rapid temporary variation, and that of the ions, with fields of a very slow variation or even static in the typical times of autocorrelation. As a consequence, it is proposed to separate the two effects as if they were physically and statistically independent, so that the relaxation of emission process that gives rise to the width of the line comes to be the product from the two relaxations treated as independents. In this way, the effects of the fast (of the electrons) and slow (of the ions) electrical fields will be calculated separately. As the author emphasizes, mathematical argument to justify this separation does not exist, and a calculation experiment is necessary to prove the possible independence between these two broadening mechanisms.

An identical separation between these two types of effects is used in the analytical calculation models [14,15] and, some times, the calculation in simulations have been simplified by the same separation. Thus, for example, in [1] only the movement of the ions is simulated to obtain the sequence from slow variation microfield (for studies of ions dynamics) and the effect of the electrons is considered as averaged in each temporary step of the simulation through an *impact operator* that considers a homogeneous relaxation produced by electron collisions.

In this work, we propose a *calculation experiment* to prove the scope of the approximation. Concretely, we have calculated the emission profiles of the Lyman-alpha hydrogen line — the simplest for calculation and, so, the one that reveals with more facility the physical processes that gives rise to the broadening — in a relatively high-density plasma. To study the independence between the effects of fast variation fields and slow variation fields, the configurations have been separated as a function of the electrical field created by the ions, that will be considered static. That is to say, the total electric field over the emitter atom has been separated in a *static* part and a *dynamic* part. The former will be constant, and the latter will correspond to the field produced by the electrons. In our calculation we will consider separately configurations with different static ionic fields, i.e., we will separate, in a certain manner, the effects of the homogeneous broadening due to the electrons and the nonhomogeneous broadening due to the static ions. In that sense, the examination that we accomplish is more exigent than that carried out in Ref. [10] to prove the hypothesis of statistical independence, in which this independence over the set of all the ionic configurations with low-frequency fields and the electronic configuration with high-frequency fields is analyzed.

In our study, we have evaluated the dependency of impact width on the quasistatic ionic field configuration. The obtained results permit to blend the statistics independence hypothesis of the two relaxation processes and lead to a possible modification of the methods of simulation that intend to reduce calculation in this way.

2. Mathematical treatment

As it is well known, the profile of a spectral line of dipole emission can be obtained from the Fourier transform of the average $\{\}$ of the autocorrelation function $C(t)$ of the atomic dipole moment $\mathbf{D}(t)$ performed over an ensemble [15]:

$$I(\omega) = \frac{1}{\pi} \mathrm{Re} \int_0^\infty \mathrm{d}t \, \{C(t)\} \mathrm{e}^{\mathrm{i}\omega t}, \tag{1}$$

$$C(t) = \mathrm{tr}(\mathbf{D}(t)\cdot\mathbf{D}(0)), \tag{2}$$

$$\mathbf{D}(t) = U^{+}(t)\mathbf{D}U(t), \tag{3}$$

where $U(t)$ is the evolution operator of the emitter atom that satisfies the Schrödinger equation

$$\mathrm{i}\hbar\frac{\mathrm{d}}{\mathrm{d}t}U(t) = H(t)U(t) \tag{4}$$

in which the Hamiltonian

$$H(t) = H_0 + q\mathbf{E}(t)\cdot\mathbf{R} \tag{5}$$

includes both the structure of the undisturbed states, H_0, and the effects of the charged perturbers on the emitter through the dipole interaction, $q\mathbf{E}\cdot\mathbf{R}$. $\mathbf{E}(t)$ is the temporary sequence of the electric microfield generated surrounding the emitter and $q\mathbf{R}(t)$ is the atom dipole moment operator.

We assume for our calculation the *no-quenching* approximation, that limits the transitions induced by the perturbed field to the states with identical energy without perturbation. In these conditions, the structure of the operators of the $n \to n'$ transition may be written as

$$H_0 = \begin{pmatrix} H_{0n} & 0 \\ 0 & H_{0n'} \end{pmatrix},$$

$$\mathbf{R} = \begin{pmatrix} \mathbf{R}_n & 0 \\ 0 & \mathbf{R}_{n'} \end{pmatrix},$$

$$U = \begin{pmatrix} U_n & 0 \\ 0 & U_{n'} \end{pmatrix}. \tag{6}$$

Then, Eq. (4) can be separated in boxes. Further, in this equation the H_0 operator is not present because it commutes with all other operators. On the other hand, as we are going to apply this treatment to the hydrogen Lyman-α line (transition between levels $n = 2$ and $n' = 1$), the evolution operator box of the lower level is not perturbed by the electric microfield. Accordingly, we write $U_1 = \mathbb{1}$ and then we must only calculate the evolution of the $n = 2$ state.

As it is well known [16], in hydrogen atoms, the restriction of the differential Equation (4) to the group of estates n can be decoupled into two separate differential equations by writing

$$\mathbf{R}_n = \tfrac{3}{2}n\hbar a_0(\mathbf{J}_n(2) - \mathbf{J}_n(1)), \tag{7}$$

where the $\mathbf{J}$ operators commute with each other and satisfy the commutation rules of an angular momentum operator.

$$[J_i, J_j] = \mathrm{i}\varepsilon_{ijk}J_k. \tag{8}$$

Then, for both energy levels n

$$U_n(t) = \exp\left[-\frac{\mathrm{i}}{\hbar}H_{0n}t\right]\Phi_{n1}(t)\Phi_{n2}(t), \tag{9}$$

$$\mathrm{i}\frac{\mathrm{d}}{\mathrm{d}t}\Phi_{nk}(t) = \chi_{nk}\mathbf{E}(t)\cdot\mathbf{J}_n(k)\Phi_{nk}(t), \tag{10}$$

$$\chi_{nk} = (-1)^k\frac{3}{2\hbar}nqa_0, \quad k = 1,2. \tag{11}$$

The solution of Eq. (10) may be written in the form

$$\Phi_{nk}(t) = \exp[-\mathrm{i}\psi_{nk}(t)\mathbf{u}_{nk}(t)\cdot\mathbf{J}(k)] \tag{12}$$

that takes the form of a rotation matrix.

2.1. Separation of static and dynamic fields

When we separate the statistics of ionic field (low-frequency) and electronic field (high-frequency), the complete spectral profile may be written [13,15] as

$$I(\omega) = \int W(\mathbf{E}^{\mathrm{s}})\,\mathrm{d}^3\mathbf{E}^{\mathrm{s}}\,I(\omega,\mathbf{E}^{\mathrm{s}}), \tag{13}$$

where $W(\mathbf{E}^{\mathrm{s}})\,\mathrm{d}^3\mathbf{E}^{\mathrm{s}}$ is the statistical distribution of the static field $\mathbf{E}^{\mathrm{s}}$ and $I(\omega,\mathbf{E}^{\mathrm{s}})$ is the spectral profile corresponding to the emission of several atoms affected by that static field and to the high-frequency electron collisions. Hence the spectrum is an incoherent sum of emission processes, each one of them with the emitter affected by a predetermined static field.

To obtain $I(\omega,\mathbf{E}^{\mathrm{s}})$ we must resolve Eq. (4) with

$$H = H_0 + q\mathbf{E}(t)\cdot\mathbf{R} = H_0 + q(\mathbf{E}^{\mathrm{s}} + \mathbf{E}^{\mathrm{d}}(t))\cdot\mathbf{R}, \tag{14}$$

where $\mathbf{E}^{\mathrm{d}}(t)$ is the dynamic field created by the electrons. All the electronic configurations (different temporary sequences of the $\mathbf{E}^{\mathrm{d}}(t)$) with the same static field produce the emission of a line shifted by static Stark effect of the $\mathbf{E}^{\mathrm{s}}$ field and broadened by electron collisions. Hence, we write:

$$U(t) = U^{\mathrm{s}}(t)U^{\mathrm{d}}(t), \tag{15}$$

where $U^{\mathrm{s}}(t)$ gives the static Stark shift:

$$\mathrm{i}\hbar\frac{\mathrm{d}}{\mathrm{d}t}U^{\mathrm{s}}(t) = [H_0 + q\mathbf{E}^{\mathrm{s}}\cdot\mathbf{R}]\;U^{\mathrm{s}}(t) \;\Rightarrow \tag{16}$$

$$U^{\mathrm{s}}(t) = \exp\left[-\frac{\mathrm{i}}{\hbar}(H_0 + q\mathbf{E}^{\mathrm{s}}\cdot\mathbf{R})\,t\right], \tag{17}$$

and $U^{\mathrm{d}}(t)$ is the rest of the evolution operator:

$$\mathrm{i}\hbar\frac{\mathrm{d}}{\mathrm{d}t}U^{\mathrm{d}}(t) = [U^{\mathrm{s}+}(t)q\mathbf{E}^{\mathrm{d}}(t)\cdot\mathbf{R}U^{\mathrm{s}}(t)]U^{\mathrm{d}}(t). \tag{18}$$

The average of the Hamiltonian that can be read in (18), gives the width produced by fast electron collisions. As will be seen, this width depends on the static field configurations that we are studying. Griem [15] calculated the averages that give the impact width operator, but the static field

Table 1
Phase change of the dynamic field modulator during an individual collision and during the mean correlation time

N_e (m^{-3})	T (K)	γ (s^{-1})	$\gamma\Delta t_{col}$ (rad)	$\gamma\Delta t_{1/2}$ (rad)
10^{22}	5000	4.2×10^{11}	3.0×10^{-2}	6.8
10^{22}	20000	4.2×10^{11}	1.6×10^{-2}	10.2
10^{23}	20000	1.9×10^{12}	3.3×10^{-2}	5.6
10^{24}	20000	9.0×10^{12}	7.2×10^{-3}	3.4
10^{24}	30000	9.0×10^{12}	5.9×10^{-3}	4.0

configuration[1] is not considered. In this work, we will contemplate on the influence of $U^s(t)$ on the Hamiltonian of (18).

Now, we separate the evolution operator into two parts (remember that the **J** operators commute with each other):

$$U^s(t) = e^{-i\omega_0 t}\exp[-i(\mathbf{u}^s\cdot\mathbf{J}(2))\gamma_2 t]\exp[-i(\mathbf{u}^s\cdot\mathbf{J}(1))\gamma_1 t]$$

$$\mathbf{u}^s = \frac{\mathbf{E}^s}{|\mathbf{E}^s|}, \qquad \gamma_k = \chi_k|\mathbf{E}^s|, \tag{19}$$

$$U^d(t) = \Psi_1(t)\Psi_2(t)$$

$$i\frac{d}{dt}\Psi_k(t) = \exp[+i(\mathbf{u}^s\cdot\mathbf{J}(k))\gamma_k t]\chi_k\mathbf{E}^d(t)\cdot\mathbf{J}(k) \tag{20}$$

$$\exp[-i(\mathbf{u}^s\cdot\mathbf{J}(k))\gamma_k t]\Psi_k(t), \quad k=1,2. \tag{21}$$

(ω_0 is the frequency of the Lyman-alpha transition). Considering the commutation rules of the $\mathbf{J}_n(k)$ operators, the Hamiltonian that is seen in Eq. (20) is the result of rotating the field $\mathbf{E}^d(t)$ with axis $\mathbf{u}_s$ and angle $\gamma_k t$:

$$i\frac{d}{dt}\Psi_k(t) = \chi_k\mathbf{E}'^d_k\cdot\mathbf{J}(k)\Psi_k(t),$$

$$\mathbf{E}'^d_k(t) = \cos(\gamma_k t)\mathbf{E}^d(t) - \sin(\gamma_k t)(\mathbf{u}^s\times\mathbf{E}^d(t)) + [1-\cos(\gamma_k t)](\mathbf{E}^d(t)\cdot\mathbf{u}^s)\mathbf{u}^s. \tag{22}$$

We must remember that in Eq. (22) we are studying two uncoupled differential equations. Both are formally identical; there is only a change of sign in χ and hence of γ. When we do not take into account the coupling between the static and dynamic fields, we are assuming $\cos(\gamma t)\approx 1$ during the mean time of an individual electron collision: $\Delta t \approx b/v$. Table 1 shows the value of this phase for the

[1] When in Ref. [15], the calculation of the $I(\omega, \mathbf{E}^s)$ spectrum is accomplished, it is really considered the influence of the value of $\mathbf{E}^s$: it is necessary to invert a matrix whose principal diagonal depends on the value of the static field. Nevertheless, elements from outside of the diagonal, those which give the effect of the electron collisions, are considered independents of this field.

Lyman-alpha line in different cases of density and electronic temperature. For the calculation of γ we take the value of a typical static field. As will be seen, the rotation performed during a collision is small.

2.2. Autocorrelation function

The expression (22) may be expressed in terms of the Euler–Rodrigues parameters [15] and, if $n = 2$ we have

$$\Phi_{21}(t) = B_0 + 2\mathrm{i}\mathbf{B}\cdot\mathbf{J}(1), \tag{23}$$

$$\Phi_{22}(t) = A_0 + 2\mathrm{i}\mathbf{A}\cdot\mathbf{J}(2). \tag{24}$$

$$A_0 = \cos(\psi_{22}/2), \qquad \mathbf{A} = \mathbf{u}_{22}\cdot\sin(\psi_{22}/2), \tag{25}$$

$$B_0 = \cos(\psi_{21}/2), \qquad \mathbf{B} = \mathbf{u}_{21}\cdot\sin(\psi_{21}/2). \tag{26}$$

The autocorrelation function 2 for the Lyman-alpha may be written in terms of these parameters in the following way [15]:

$$C(t) = A_0 B_0 - \tfrac{1}{3}\mathbf{A}\cdot\mathbf{B}. \tag{27}$$

If we write the Ψ_k operators of Eq. (20) in terms of the Euler–Rodrigues parameters and introduce them in Eq. (27) we obtain

$$\begin{aligned} C(t) = {} & \tfrac{2}{3}[\alpha_0\beta_0 - (\mathbf{u}^{\mathrm{s}}\cdot\boldsymbol{\alpha})(\mathbf{u}^{\mathrm{s}}\boldsymbol{\beta})]. \\ & + \tfrac{1}{3}[\alpha_0\beta_0 - \boldsymbol{\alpha}\cdot\boldsymbol{\beta} + 2(\mathbf{u}^{\mathrm{s}}\cdot\boldsymbol{\alpha})(\mathbf{u}^{\mathrm{s}}\cdot\boldsymbol{\beta})]\cos(\gamma t) \\ & + \tfrac{1}{3}[\alpha_0(\mathbf{u}^{\mathrm{s}}\cdot\boldsymbol{\beta}) - \beta_0(\mathbf{u}^{\mathrm{s}}\cdot\boldsymbol{\alpha}) - \mathbf{u}^{\mathrm{s}}(\boldsymbol{\alpha}\times\boldsymbol{\beta})]\sin(\gamma t), \end{aligned} \tag{28}$$

where $[\alpha_0;\boldsymbol{\alpha}]$ and $[\beta_0;\boldsymbol{\beta}]$ are, respectively, the Euler–Rodrigues parameters corresponding to the $\Psi_2(t)$ and $\Psi_1(t)$ operators, and γ is the positive value $\gamma = \gamma_2$.

Now, we define from Eq. (28) three new autocorrelation functions: C_{C}, $C_{\mathrm{L}}^{\mathrm{c}}$ and $C_{\mathrm{L}}^{\mathrm{s}}$, so that

$$C(t) = \tfrac{2}{3}C_{\mathrm{C}}(t) + \tfrac{1}{3}C_{\mathrm{L}}^{\mathrm{c}}(t)\cos(\gamma t) + \tfrac{1}{3}C_{\mathrm{L}}^{\mathrm{s}}(t)\sin(\gamma t) \tag{29}$$

or, in another form

$$C(t) = \tfrac{2}{3}C_{\mathrm{C}}(t) + \tfrac{1}{6}\mathrm{e}^{\mathrm{i}\gamma t}[C_{\mathrm{L}}^{\mathrm{c}}(t) + C_{\mathrm{L}}^{\mathrm{s}}(t)] + \tfrac{1}{6}\mathrm{e}^{-\mathrm{i}\gamma t}[C_{\mathrm{L}}^{\mathrm{c}}(t) - C_{\mathrm{L}}^{\mathrm{s}}(t)]. \tag{30}$$

This expression shows that the autocorrelation function $C_{\mathrm{C}}(t)$ gives the central part of the line, whereas the lateral part depends on $C_{\mathrm{L}}^{\mathrm{c}}(t)$ and $C_{\mathrm{L}}^{\mathrm{s}}(t)$ functions.

3. Results

To analyze the possible statistical independence between both the broadening mechanisms cited, we have carried out some spectrum calculations for the hydrogen Lyman-alpha transition, in which we separate the configurations of the emitter according to the static electrical field over the emitter. This separation will permit us to analyze, in each case, the form and width of the line due to the impact broadening mechanisms. We have already seen (Eq. (22)) that, initially, both broadening

mechanisms are not independent, since the electrical field that gives rise to the impact broadening in each static field configuration is *modulated* by the quasistatic field. However, the influence that could cause such a modulation is something difficult to predict, since the typical variation frequencies of the electronic field and the variation that is induced by the quasistatic field are very different. It could happen that this modulation has no practical consequence and remains hidden by the fast variation of the dynamic field due to the electrons. Then, the statistical separation between these two fields would be justified. The simulation that we have used is, in a certain way, a *calculation experiment*.

In each configuration of static electrical field $\mathbf{E}^{s}$, we have registered the three components of the autocorrelation function: $C_{C}(t)$, $C_{L}^{C}(t)$ and $C_{L}^{S}(t)$, that, together, provide the total autocorrelation function 30. The results show that the temporary behavior of the central component leads to a Lorentzian-type spectral line, as we would expect in an impact model. However, the lateral component shows a very different behavior.

The expression of the autocorrelation function establishes the center of the two lateral components in connection with the static Stark shift frequency $\xi = \pm 3qa_0E^{s}$. We must remember that static ion models with a treatment of impact broadening for the electron collisions predict for each one of those components an impact width independent of the value of the electric field of each configuration. This is not what is observed in the simulation. In addition to the conventional Stark shift that we would expect, the lateral components present a greater width while the electrical field static gets greater. This result is shown in Fig. 1, where the width at half-maximum of the lateral component as a function of the static field has been represented.

For a better explanation of this phenomenon we return to the expression of the dynamic field (Eq. (22)). The *effective electrical field* $\mathbf{E}_k^{d}$ has a very fast variation since $\mathbf{E}^{d}$ has it too. Consequently, its effect over the emitter gives rise to the typical impact broadening, in which the middle number of collisions by time unit that receives the emitter is fundamental. In the analytical models, the coupling between static and dynamic parts of the electrical field is ignored. This is equal to

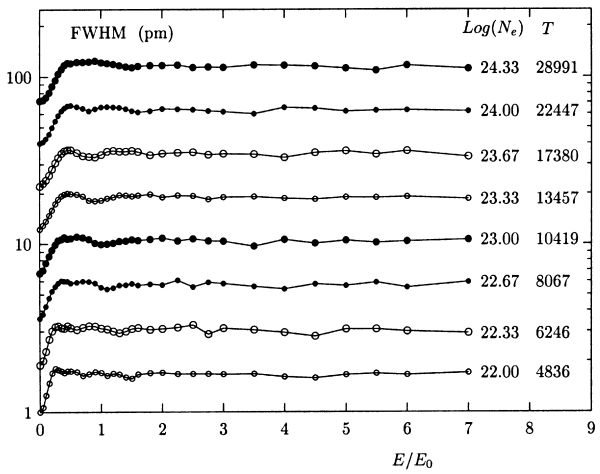

Fig. 1. Full-width at half-maximum of the Stark lateral component as a function of the static electrical field for several cases of density and temperature. E_0 is the value of the typical field (N_e in m^{-3}, T in K).

considering that, in the lifetime of the emission process, $\cos(\gamma t) \simeq 1$ (and then, $\sin(\gamma t) \simeq 0$). We have already seen in the mathematical treatment that the effect of the modulator during an individual collision is negligible. But we should take into account that the rotation effect is accumulated as the correlation time elapses. If we consider the rotation effect over the dynamic field during all the correlation time, a variation in the line width begins to appear. In fact, for a plasma density $N_e = 10^{23}$ m^{-3} and $T = 20000$ K, the Stark width of the Lyman-alpha line is, according to the calculations of Griem [14] for static ions, $\Delta\lambda_{1/2} = 3.2 \times 10^{-4}$ nm, which supposes a mean life $\Delta t_{1/2} = 2.4 \times 10^{-11}$ s. Taking into account that for this plasma and for the most probable value of the static field, we have $\gamma \Delta t_{1/2} \simeq 10.2$ rad, therefore, in the correlation loss typical time at the emission process, the function $\cos(\gamma t)$ has done about two complete cycles. This modulation causes that the effective dynamic field has a considerable change during the emission process. It would seem that this slow rotation is superposed to a field that change very quickly and, as we will average forward many microfield sequences, that slow rotation has very little importance. However, we must take into account that when we calculate the Euler–Rodrigues parameters, $[\alpha_0, \boldsymbol{\alpha}]$ parameters are calculated when the dynamic field rotates in one direction, and $[\beta_0, \boldsymbol{\beta}]$ parameters are calculated when the field rotates in the opposite direction. This additional dynamic effect will increase the loss of correlation between vectors $\boldsymbol{\alpha}$ and $\boldsymbol{\beta}$ and then it will increase the line width.

To evaluate somehow the effect of the field rotation over the autocorrelation function, we calculate the value of this function taking a temporal interval small enough to be able to consider that the electronic field remains constant. In this case, the Euler–Rodrigues parameters acquire the values

$$[\alpha_0; \boldsymbol{\alpha}] = [\cos(\Gamma\tau); \sin(\Gamma\tau)\mathbf{u}_2], \tag{31}$$

$$[\beta_0; \boldsymbol{\beta}] = [\cos(\Gamma\tau); -\sin(\Gamma\tau)\mathbf{u}_1], \tag{32}$$

where $\Gamma = \chi_2 |\mathbf{E}^d|$, τ is the temporal interval, and $\mathbf{u}_1$ and $\mathbf{u}_2$ are unitary vectors in the direction of the dynamic field after the rotation with angles $-\gamma$ and γ, respectively.

Then, the correlation loss in that temporal interval will be

$$C_C(\tau) = \cos^2(\Gamma\tau) + \sin^2(\Gamma\tau)(u_{||}^2), \tag{33}$$

$$C_L^c(\tau) = 1 - 2\sin^2(\Gamma\tau)[1 - (1 - u_{||}^2)\cos^2(\gamma\tau)], \tag{34}$$

$$C_L^s(\tau) = -\sin(2\Gamma\tau)u_{||} - \sin(2\gamma\tau)[1 - u_{||}^2], \tag{35}$$

where $u_{||} = \mathbf{u}^d \cdot \mathbf{u}^s$.

As it can be seen, the expression of the autocorrelation function $C_L^c(t)$ (Eq. (35)) contains a part that depends on the static field value through the γ value in the cosine. Because of this part the value of $C_L^c(t)$ function decays quickly, and the line width will increase. This is the phenomenon that is observed in Fig. 1. Not considering the influence of the static field in accomplishing the calculations, would mean that the value of $C_L^c(t)$ is overvalued.

In the expression of $C_L^S(t)$ there also appears a term that depends on the static field. Note that if we do not consider this term, the function $C_L^S(t)$ would be proportional to $u_{||} = \mathbf{u}^d \cdot \mathbf{u}^s$, whose value is zero on average. Therefore, if we do not take into account the static field influence over the autocorrelation function, the average of $C_L^S(t)$ taken in a set of electronic configurations would be zero. As it can be seen in Eq. (30), this function is the cause of the asymmetry of the spectral line

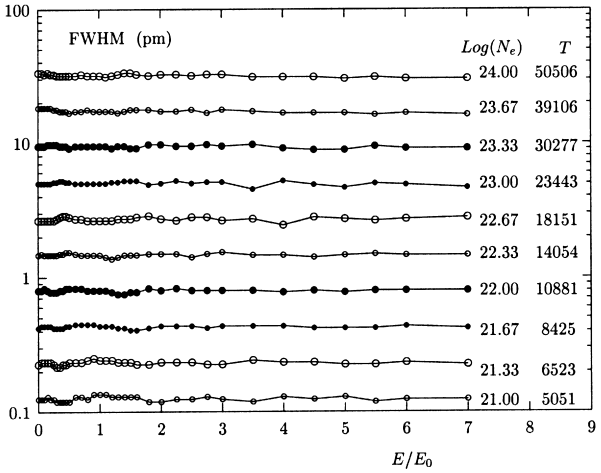

Fig. 2. Full-width at half-maximum of the Stark central component as a function of the electrical static field (N_e in m^{-3}, T in K).

lateral components. A calculation based on the relaxation methods would give a lorentzian spectral profile for each ionic field configuration, but through a detailed calculation an asymmetrical lines overlapping is obtained.

Concerning the central component (C_C), Eq. (34) shows that it is independent of the static field value. We must take into account that the central component is carried on with the temporary evolution of a state that does not suffer the static Stark effect. In fact, if we remember the expression of the autocorrelation function central component (Eq. (28)), we will see that in its expression only that part of the Euler–Rodrigues parameters that is parallel to the static field appears, that is not affected by the rotation and for this reason it is independent of the static field (see Fig. 2).

4. Conclusions

The accomplished study shows that a coupling is produced between the quasistatic Stark displacement of the emission frequencies and the emission phase memory loss caused by the electron collisions. This phenomenon, proved very simply through a traditional treatment, could not be considered by a treatment based on two independent relaxation mechanisms. Then, we can conclude that the two mechanisms dominant in the Stark broadening in plasmas cannot be considered as either physically or statistically independent.

It is true that in our treatment we have been very demanding. We have requested that both broadening mechanisms must be statistically independent in each collectivity with the same quasistatic ionic field. If we reduce this exigency requesting that they should be statistically independent in global form, the result would be different. In fact, the practical calculations of Ref. [10] show it. But we must take into account that our calculations have been accomplished with relatively high-density plasmas, in which the individual emission components are considerably overlapped, masking in this way their individual form; however, we have been able to show the

coupling in the two phenomena. In lower density plasmas, the Stark width due to the dynamic fields is very small so that the individual details of each one of them are shown quickly, and then, the coupling will be seen with more clarity. This example cautions one against excessive confidence in the relaxation methods, because they can hide interesting physical processes.

Anyway, a calculation method that intends to reduce the volume of the simulations calculation must take into account this phenomenon considering an electron impact width (or a relaxation time) whose magnitude depends on the module of the ionic field of each configuration.

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