Polarization of Fe XXV $1s^2$-$1s2l$ lines: Collisional resonances and radiative cascade contributions to $1s2l$ magnetic-sublevel excitation rates

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The effects of collisional resonances and radiative cascades on the linear polarization of He-like iron (Fe XXV) lines from $1s2l$ to $1s^2$ levels have been investigated. Detailed calculations have been carried out for the $1s3l3l'$ resonance contributions to electron-impact excitation rates from the $1s^2$ ground level to the individual magnetic sublevels of $1s2l$ configurations. Excitation collision strengths from $1s^2$ to $1s3l$ sublevels as well as radiative cascade transitions from $1s3l$ to $1s2l'$ sublevels have also been computed. The autoionization transition-matrix elements (for the resonance effects) and the collison strengths (for the radiative cascades) have been computed in the distorted-wave approximation using intermediate coupling with fine-structure mixing multiconfiguration bound wave functions. The results indicate that the collisional resonance contributions, when averaged over a small energy range just covering them, have a somewhat significant depolarizing effect on the ($1s^2\,{}^1S_0-1s2p\,{}^1P_1$, electric dipole) w line, the ($1s^2\,{}^1S_0-1s2p\,{}^3P_2$, magnetic quadrupole) x line, and the ($1s^2\,{}^1S_0-1s2p\,{}^3P_1$, fine-structure electric dipole) y line. However, the averaged polarization degree over some single resonances can reach high values. Now for the ($1s^2\,{}^1S_0-1s2s\,{}^3S_1$, relativistic magnetic dipole) z line the individual resonances are found to induce a polarizaton degree less than 15% in absolute value but after averaging over all $1s3l3l'$ resonances z remains practically unpolarized. Concerning radiative cascades, it is shown that they can create a relative small degree of polarization on the z line, the highest value being close to 14%. For w, x, and y lines, the cascades have a weak depolarizing effect, except for x at large incident energy. It is expected that the contributions from $1snl$ cascades as well as $1s3lnl'$ resonances, with $n \geq 4$, would not greatly change the results already obtained.

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I. INTRODUCTION

An accurate knowledge of the polarization of He-like lines emitted by highly charged ions excited by unidirectional electron impact may be important for a spectroscopic diagnostic of the electron distribution anisotropy in high-temperature plasmas. Indeed, such polarization may provide information about both angular and energy distributions of nonthermal electrons existing in astrophysical and laboratory plasmas, particularly in the solar corona [1], in tokamak plasmas [2], and laser-produced plasmas [3]. Due to the geometry of the electron beam, accurate values of the line polarization may also be important for interpreting high-energy resolution measurements carried out on electron-beam ion trap (EBIT) sources. More particularly, experimental data for Fe XXV and Sc XX He-like ions have been obtained at Livermore National Laboratory, for electron beams of well-defined energy [4,5].

In a previous paper (Ref. [6], referred to hereafter as paper I) we reported on calculations of the linear polarization for He-like iron (Fe XXV) lines where the population mechanism of the emitting levels was considered to be only *direct* electron-impact excitation from the ground level $1s^2$ to the $1s2l$ excited levels. For unidirectional incident electrons, some excitation cross sections were found different for different M_J magnetic sublevels belonging to $1s2p$ $J=1,2$ levels. However, taking the quantization axis along the electron beam and assuming these electrons unpolarized, the cross sections were equal for the same $|M_J|$ sublevels. The alignment of the excited levels, i.e., the unequal population of different $|M_J|$ sublevels, gave rise to a relatively high degree of linear polarization of the lines emitted from these levels. In low-density plasmas, for He-like ions where the ground level has only one sublevel and is the unique populated level, the measurements of the intensity and polarization of the He-like lines could therefore be valuable in detecting anisotropic electron distributions.

In the present paper, we improve the former study by inserting the contributions of $1s3l3l'$ collisional resonances on the linear polarization of the same $n=2$ to $n=1$ Fe XXV lines. In the second part of this paper, we also include *indirect* electron-impact excitation of the $1s2l$ levels via radiative cascades from the $1s3l$ levels. It is well known that collisional resonance excitation near threshold may be an important mechanism for populating the upper levels of forbidden transitions. In fact, a number of papers [7–9] have already investigated this resonance excitation process for He-like ions but, by assuming an isotropic Maxwellian electron distribution, they were only concerned with intensities of unpolarized lines. These calculations show that the dominant contribution came from the $1s3lnl'$ autoionizing levels with $n=3$. In Refs. [8,9], where both results almost agree, it is found that for the most sensitive line to resonance effects, the ($1s^2\,{}^1S_0-1s2s\,{}^3S_1$, magnetic dipole) z line (Gabriel notation [10]), the resonant contribution to the excitation rate is only $\sim 15-25\,\%$ at electron temperatures of interest for Fe XXV [i.e., $(1.5-3.5)\times 10^7$ K]. For EBIT experiments, the electron beam being now almost monoenergetic, the polarization results could be quite affected for energies in the $1s3l3l'$ resonance region because the resonance effects are not smeared by averaging over a Maxwellian energy distribution.

Concerning *indirect* excitations of the $1s2l$ levels via radiative cascades from higher levels, their contributions to the *effective* excitation rates for the $1s^2\,{}^1S_0-1s2p\,{}^1P_1$ (w line) and $1s2p\,{}^3P_1$ (y line) in Fe XXV have already been calculated in Ref. [11], using collision strengths of Sampson, Parks, and Clark [12] for $n\geq 3$ levels and the same radiative branching ratios as the ones used in the present paper. The respective differences between direct and effective cross sections were found to be, respectively, less than $\sim 7\%$ and 15% for any energy. Now for the transition to the $1s2p\,{}^3P_2$ level (x line) the cascade contributions is about 12% for energies just above the $n=3$ excitation threshold but becomes very significant for higher energies, reaching 70% at 8 times the $n=3$ threshold energy. Finally, the effective excitation cross section for the $1s2s\,{}^3S_1$ level (z line) is due more to indirect than to direct excitation, the former ones contributing by 55% at energies near the $n=2$ threshold up to 75% near the $n=3$ threshold. We expect, in the present work, that indirect excitations could lead to a significant depolarization of the x line for high impact energies. On the other hand, they could also cause some polarization on the z line which is unpolarized in a population model including only direct excitations.

The theory is described in Sec. II, which is divided into two subsections. In the first we derive the expression for the resonance contribution to the magnetic sublevel excitation rate coefficients for ions initially in $J_j=0$, the target total angular momentum (e.g., $1s^2\,{}^1S_0$). In the second we present the general formula for calculating the radiative cascade contribution to the $1s2l$ sublevel collision strengths. Section III deals briefly with the computation methods, which are basically the same as the ones already used in paper I and in Ref. [8]. The results are discussed in Sec. IV, which is also divided into two subsections. The first presents the average effect of the $1s3l3l'$ resonances on the linear polarization of the w,x,y,z lines. We also examine the effects that the resonances have individually on these polarizations. Particular attention is given to the influence of each contributing partial wave to the autoionization probabilities. The second presents the results of alignment of the $1s3l$ cascading levels for several electron energies: from 582 Ry, the excitation threshold, to 5000 Ry. Finally, we evaluate the contribution of radiative cascades from the $1s3l$ levels, as well as cascades between $n=2$ levels, on the polarization of the $1s^2$-$1s2l$ lines.

II. THEORY

A. Resonant contribution to the magnetic sublevels' excitation rates

As in paper I, we are interested in the polarization of line radiation emitted from degenerate ionic states (i.e., the magnetic sublevels) excited by unidirectional but unpolarized electrons, more precisely, radiation from the sublevels of the $1s2p\,{}^1P_1$, $1s2p\,{}^3P_{1,2}$, and $1s2s\,{}^3S_1$ levels to the ground $1s^2\,{}^1S_0$ level of He-like ions. Though the definition of these states depends on the choice of the quantization axis (i.e., z axis), polarization observations depend only on the angle between the incident-electron and emitted-photon directions. Using the density-matrix formalism for the excited states, we can transform the results by a frame rotation. If we choose the z axis to be along the direction of the incident-electron beam and assume a low-density plasma, this density matrix becomes diagonal and its elements are directly proportional to the sublevels excitation cross sections. In paper I, the general expression for the scattering amplitude of the excitation from an initial state α_j to a final state α_i was given in terms of the transition matrix T. The resulting expression for excitation cross section is

$$Q(\alpha_i\leftarrow\alpha_j)=\frac{\pi}{2k_j^2}\sum_{m_{s_i},m_{s_j}}\int d\hat{\mathbf{k}}_i\left|\sum_{l_i,m_{l_i},l_j}(2l_j+1)^{1/2}i^{l_j-l_i}e^{i(\sigma_i+\sigma_{l_j})}Y_{l_i}^{m_{l_i}}(\hat{\mathbf{k}}_i)T_{\beta_i\beta_j}\right|^2, \tag{1}$$

where $Y_l^{m_l}(\hat{\mathbf{k}})$ denotes a spherical harmonic, σ_l the Coulomb phase shift, and $T_{\beta_i\beta_j}$ the transition matrix element from β_j to β_i, β being the set of quantum numbers describing the total system in the representation of uncoupled angular momenta $\beta\equiv(\alpha l m_l m_s)$. The target states are defined in terms of the quantum numbers $\alpha\equiv(\Delta JM)$ where J and M stand for the target total angular momentum and its z component and Δ represents the additional quantum numbers necessary to uniquely define the state. The incident-electron quantum numbers are l, s $(=\frac{1}{2})$, m_l, m_s (orbital and spin angular momenta as well as their z components), and $\mathbf{k}=(k,\hat{\mathbf{k}})$, its wave vector ($k=|\mathbf{k}|$ and $\hat{\mathbf{k}}$ is the unit vector in the direction $\mathbf{k}$).

In a previous work (Ref. [13], referred to hereafter as paper II), the dielectric capture of unidirectional electron was considered. That work was concentrated on the polarization of Li-like dielectronic satellite lines emitted from $1s2lnl'$ autoionizing levels as a possible electron anisotropy diagnostic for the lowest-energy part of the nonthermal electron distribution in plasmas. Taking again the z quantization axis along the incident-electron beam we noted that only the $M_J=\pm\frac{1}{2}$ sublevels of the autoionizing levels could be populated in dielectronic capture transitions originating from ${}^1S_0$ target states. Similarly, in the present work, we shall see that, to calculate the resonance effects on the alignment of the $n=2$ He-like excited levels, we have to determine the alignment transfer accompanying the autoionization of $1s3l3l'$, $M_J=\pm\frac{1}{2}$ sublevels.

In the scattering process of electrons by N-electron ions, the resonances are closely related to the presence of autoionizing states of the $(N+1)$-electron system. For highly ionized atoms, the T matrix elements can be approximated using a perturbational approach (e.g., distorted-wave approximation), which gives a reliable description of the resonance process [8]

$$T_{\beta_i\beta_j}=2i\left[\langle\beta_i|V|\beta_j\rangle+\sum_{\alpha_d}\frac{\langle\beta_i|V|\alpha_d\rangle\langle\alpha_d|V|\beta_j\rangle}{E_j+\varepsilon_j-E_d+i\Gamma_d/2}\right], \tag{2}$$

where the V operator contains the two-electron (Coulombic) interaction, the matrix element $\langle\cdots|V|\cdots\rangle$ being chosen real. The first term on the right-hand side, $\langle\beta_i|V|\beta_j\rangle$, describes the *direct* excitation. This term is a smooth function of the incident-electron energy ε_j $(=k_j^2/2$ in a.u.). The second term, $\sum_{\alpha_d}$, is a function of ϵ_j varying rapidly near the energy E_d (in a.u.) of the autoionizing state α_d and gives the resonant contribution. The matrix element $\langle\alpha_d|V|\beta_j\rangle$ represents the dielectronic capture process, i.e., the capture of a free electron of energy ε_j by the ground-level α_j of the N-electron system forming the α_d doubly excited state of the $(N+1)$-electron system; the matrix element $\langle\beta_i|V|\alpha_d\rangle$ corresponds to the autoionization of the α_d doubly excited state to the continuum associated to the α_i excited or ground state of the N-electron system. The summation over α_d extends over all possible autoionizing states which can decay to the α_i state. The resonance total width Γ_d is the sum of the radiative width Γ_d^r and of the autoionization width Γ_d^a. Taking Eq. (16) as normalization of the electron-continuum wave function, width Γ_d^a is given by $\Gamma_d^a=4\sum_{\beta_k}|\langle\beta_k|V|\alpha_d\rangle|^2$, in atomic units. In the calculations of excited cross sections, reported in paper I, only the direct term $\langle\beta_i|V|\beta_j\rangle$ was taken into account. Here, we examine the resonance contribution which can be treated separately from the direct part since, for highly ionized atoms, this latter part is relatively small in the resonance energy range. Indeed the direct term behaves like z^{-1} while the resonant term is like z^0, where z is the effective (or ionic) charge: $z=Z-N$, Z being the nuclear charge.

Owing to the choice of z axis along the direction $\hat{\mathbf{k}}_j$ of the incident electron $m_{l_j}=0$. For the capture on a target level $J_j=0$ it follows, from the conservation of angular momenta, that $M_d=m_{s_j}$, so that only the matrix elements for which $M_d=\pm\frac{1}{2}$ need to be considered in Eq. (2). Adopting the pair-coupling representation $|J_il_iK_i(s_i);J_d\rangle$, where J_i is coupled to l_i to give K_i, which is then coupled to s_i to give J_d, and using the Wigner-Eckart theorem, the autoionization transition-matrix element can be expressed as

$$\langle\beta_i|V|\alpha_d\rangle=\frac{1}{[J_d]}\sum_{K_i,M_{K_i}}C^{J_i\,l_i\,K_i}_{M_i\,m_{l_i}\,M_{K_i}}C^{K_i\,1/2\,J_d}_{M_{K_i}\,m_{s_i}\,M_d}\times\langle\Delta_iJ_i\varepsilon_i^dl_iK_i;J_d\|V\|\Delta_dJ_d\rangle, \tag{3}$$

where $\varepsilon_i^d=E_d-E_i$ and $[J]\equiv(2J+1)^{1/2}$. $C^{j_1\,j_2\,j_3}_{m_1\,m_2\,m_3}$ is a Clebsch-Gordan (CG) coefficient and $\langle\cdots\|V\|\cdots\rangle$ is the reduced matrix element. E_i is the energy of the α_i state. It should be noted that V is a scalar operator and therefore conserves the angular momentum J_d of the total system. The expression for the dielectronic capture matrix element $\langle\alpha_d|V|\beta_j\rangle$ is similar but, for target ions in initial states $J_j=0$ (e.g., $1s^2\,{}^1S_0$), the algebra is greatly simplified, i.e., $K_j=l_j$ and $M_{K_j}=0$. Moreover, it is worth mentioning that for a given autoionizing level Δ_dJ_d only one value of l_j is allowed; this value is determined from both the triangular relations $l_j=J_d\pm\frac{1}{2}$ and parity conservation.

The resonances considered in this work belong to the $1s3l3l'$ configurations which give rise to autoionizing levels with large energy separation compared to their widths. We can assume the individual resonances to be isolated. This means simply that the total resonant contribution to the excitation rate can be written as a summation over Δ_dJ_d. Now by taking into account the earlier remark of the one-to-one correspondence between l_j and Δ_dJ_d, the summation over l_j in Eq. (1) disappears. Therefore interferences between incident partial waves do not occur.

It is useful to introduce the autoionization probability $A^a(\alpha_d\rightarrow\alpha_i)$ from the α_d sublevel to the continuum relative to the α_i sublevel defined by

$$A^0(\alpha_d\rightarrow\alpha_i)=4\sum_{m_{s_i}}\int d\hat{\mathbf{k}}_i\left|\sum_{l_i,m_{l_i}}i^{-l_i}e^{i\sigma_{l_i}}Y_{l_i}^{m_{l_i}}(\hat{\mathbf{k}}_i)\langle\beta_i|V|\alpha_d\rangle\right|^2, \tag{4}$$

where the integration and summation are over all directions $\hat{\mathbf{k}}_i$ and quantum numbers l_i, m_{l_i}, and m_{s_i} of the *unobserved* scattered electrons. The orthogonality of the spherical harmonics implies the disappearance of the Coulomb phase σ_{l_i}. Inserting Eq. (3) into Eq. (4), the derivation of $A^a(\alpha_d\rightarrow\alpha_i)$ involves a product of four CG coefficients and a four-fold summation over the magnetic quantum numbers m_{s_i}, m_{l_i}, M_{K_i} and $M_{K_i'}$. Now using the rules for nonzero CG coefficient, we have $M_{K_i}=M_i+m_{l_i}$ and $M_d=M_{K_i}+m_{s_i}$ and therefore the 4-summation reduce to only one on m_{l_i}:

$$A^a(\alpha_d \rightarrow \alpha_i) = \frac{4}{[J_d]^2} \sum_{K_i, K_i', l_i} \langle \Delta_i J_i \varepsilon_i^d l_i K_i J_d \| V \| \Delta_d J_d \rangle \langle \Delta_i J_i \varepsilon_i^d l_i K_i' J_d \| V \| \Delta_d J_d \rangle \times \sum_{m_{l_i}} C^{J_i\ l_i\ K_i}_{M_i\ m_{l_i}\ M_i+m_{l_i}} C^{K_i\ 1/2\ J_d}_{M_i+m_{l_i}\ M_d-M_i-m_{l_i}\ M_d} C^{J_i\ l_i\ K_i'}_{M_i\ m_{l_i}\ M_i+m_{l_i}} C^{K_i'\ 1/2\ J_d}_{M_i+m_{l_i}\ M_d-M_i-m_{l_i}\ M_d} . \quad (5)$$

From the unitarity of CG coefficients, the level-to-level autoionization probability $A^a(\Delta_d J_d \rightarrow \Delta_i J_i)$ is related to the reduced matrix elements by

$$A^a(\Delta_d J_d \rightarrow \Delta_i J_i) = \sum_{M_i} A^a(\alpha_d \rightarrow \alpha_i) = \frac{4}{[J_d]^2} \sum_{l_i, K_i} |\langle \Delta_i J_i \varepsilon_i^d l_i K_i; J_d \| V \| \Delta_d J_d \rangle|^2 . \quad (6)$$

In Eq. (5) we note that the relative signs of the reduced matrix elements for the same l_i but different K_i are important in evaluating $A^a(\alpha_d \rightarrow \alpha_i)$. Furthermore, it is worth mentioning that Eq. (6) gives the total autoionization probability since all states M_d of a level decay by autoionization to the continuum relative to some level with the same probability.

It is now possible to derive from Eqs. (1), (2), and (4) the expression of the resonance excitation cross section from a level $^1S_0$ ($J_j = 0$) to a magnetic sublevel $\Delta_i J_i M_i$ through a particular $\Delta_d J_d$ autoionizing level. However, since the energy spread of the incident-electron beam is usually very large compared to resonance widths, it is more convenient to use the corresponding excitation rate C_{res}, i.e., the average over the energy distribution of $v_j Q(\alpha_i \leftarrow \alpha_j)$, where v_j ($= k_j$ in a.u.) is the incident-electron velocity. For an arbitrary (normalized to 1) energy distribution $f(\varepsilon_j)$ of the incident electrons we have (in a.u.)

$$C_{\text{res}}(^1S_0 \rightarrow \alpha_i; \Delta_d J_d) = \frac{\pi^2}{2^{1/2}} \frac{f(E_d)}{E_d^{1/2}} \tilde{F}_2^d(\alpha_i) , \quad (7)$$

where the factor $\tilde{F}_2^d$ is defined by

$$\tilde{F}_2^d(\alpha_i) = \frac{(2J_d+1)}{2} \frac{A^a(\Delta_d J_d \rightarrow {}^1S_0) \sum_{M_d=\pm 1/2} A^a(\alpha_d \rightarrow \alpha_i)}{\Gamma_d} , \quad (8)$$

with

$$\tilde{F}_2^d(\Delta_i J_i) = \sum_{M_i} \tilde{F}_2^d(\alpha_i) = (2J_d+1) \frac{A^a(\Delta_d J_d \rightarrow {}^1S_0) A^a(\Delta_d J_d \rightarrow \Delta_i J_i)}{\Gamma_d} . \quad (9)$$

The total resonant contribution to excitation rate $C_{\text{res}}(^1S_0 \rightarrow \alpha_i)$ is then obtained by summing Eq. (7) over all $\Delta_d J_d$ autoionizing levels.

In Eq. (5), using symmetry relations of the CG, it can be seen that the summation over $M_d = \pm \frac{1}{2}$ of the $A^a(\alpha_d \rightarrow \alpha_i)$, in Eq. (8), does not depend on the sign of M_i. This verifies that the autoionization decay process conserves the alignment. Note also that the determination of the linear polarization of subsequent emitted line from the $\Delta_i J_i$ level requires the knowledge of the matrix elements for the autoionization transition. This is in contrast to the dielectronic satellite lines emitted in a stabilizing decay of the $1s2l3l'$ autoionizing levels, where the linear polarization could be evaluated purely geometrically, as pointed out in Sec. 2.2 of paper II. However, in the case where only the $l_i = 0$ partial wave (s wave) of the scattered electron contributes dominantly to the total autoionization probability Eq. (6), Eq. (5) reduces to

$$A^a(\alpha_d \rightarrow \alpha_i) = A^a(\Delta_d J_d \rightarrow \Delta_i J_i) (C^{J_i\ 1/2\ J_d}_{M_i\ M_d-M_i\ M_d})^2 . \quad (10)$$

The $A^a(\Delta_d J_d \rightarrow \Delta_i J_i)$ appearing as a multiplying factor common to every M_i sublevel cancels in the expression of the linear polarization, which can be predicted only in terms of angular coupling coefficients, as pointed out by Baranger and Gerjouy [14]. Unfortunately, as will be seen below, the importance of d and g waves cannot be neglected in most cases.

For $J_d = \frac{1}{2}$, using CG unitarity in Eq. (5), it is easily seen that the excitation rate, defined in Eq. (7), does not depend on the magnetic quantum number M_i. This is consistent with the fact that the collisional excitation via a resonance $J_d = \frac{1}{2}$ populates equally the magnetic sublevels and tends consequently to reduce the polarization degree of the line originating from these sublevels.

B. Radiative cascade contributions from higher sublevels

The transfer of higher-level alignment accompanying the cascade transition can be evaluated using only the square of a CG coefficient, according to the expression

$$A^r(\Delta_k J_k M_k \rightarrow \Delta_i J_i M_i) = (C^{J_i\ j\ J_k}_{M_i\ M_k-M_i\ M_k})^2 \times A^r(\Delta_k J_k \rightarrow \Delta_i J_i) , \quad (11)$$

where $A^r(\Delta_k J_k \rightarrow \Delta_i J_i)$ is the probability for radiative decay from a cascading level $\Delta_k J_k$ to the level of interest $\Delta_i J_i$, and j denotes the photon angular momentum. Here we shall be concerned only with electric dipole cascade transitions; then $j = 1$ and $|M_k - M_i| \leq 1$ in the above equation. Magnetic dipole and quadrupole transitions were important for $1s^2$-$1s2l$ radiative transitions but they are not for $1s2l$-$1snl'$ transitions, with $n \geq 3$. Note that

Eq. (11) summed over the final sublevels M_i gives the total probability $A^r(\Delta_k J_k \rightarrow \Delta_i J_i)$, which does not depend on M_k. For low-density plasmas ($N_e \leq 10^{16}$ cm^{-3}), the effective collision strength from the ground level $\Delta_j J_j$ to the sublevel α_i can be expressed as the sum of the *direct* and *cascade* contributions.

$$\Omega_{\rm eff}(\Delta_j J_j \rightarrow \Delta_i J_i M_i) = \Omega_{\rm dir}(\Delta_j J_j \rightarrow \Delta_i J_i M_i) + \Omega_{\rm cas}(\Delta_j J_j \rightarrow \Delta_i J_i M_i) \; . \quad (12)$$

In the present work, we have neglected cascades from $1snl'$ with $n \geq 4$ and cascades between $n=3$ levels because they are not as important as cascades inside $n=2$ and from $n=3$. To obtain $\Omega_{\rm cas}(\Delta_j J_j \rightarrow \Delta_i J_i M_i)$ one has to sum the products of the *direct* collision strength from $\Delta_j J_j = 1s^2\,{}^1S_0$ to a specific intermediate state $\Delta_k J_k M_k = 1snl'\; J_k M_k$, by the branching ratio of subsequent radiative decay to the $\Delta_i J_i M_i = 1s2l\; J_i M_i$ state:

$$\Omega_{\rm cas}(\Delta_j J_j \rightarrow \Delta_i J_i M_i) = \sum_k \Omega_{\rm dir}(\Delta_j J_j \rightarrow \Delta_k J_k M_k) \frac{A^r(\Delta_k J_k M_k \rightarrow \Delta_i J_i M_i)}{\Gamma^r(\Delta_k J_k)}$$
$$= \sum_{\Delta_k, J_k} R(\Delta_k J_k \rightarrow \Delta_i J_i) \sum_{M_k} \Omega_{\rm dir}(\Delta_j J_j \rightarrow \Delta_k J_k M_k) (C^{J_i\;1\;J_k}_{M_i\;M_k - M_i\;M_k})^2 , \quad (13)$$

the sum being extended over all $n=2$ and $n=3$ higher states $\Delta_k J_k M_k$ connected to the $n=2$ $\Delta_i J_i M_i$ state by an electric dipole transition. $R(\Delta_k J_k \rightarrow \Delta_i J_i)$ denotes the branching ratio for the level-to-level transition. The natural width $\Gamma^r(\Delta_k J_k)$ of the $\Delta_k J_k$ cascading level is given by

$$\Gamma^r(\Delta_k J_k) = \sum_l A^r(\Delta_k J_k \rightarrow \Delta_l J_l) \; , \quad (14)$$

where the sum includes all lower levels to which level $\Delta_k J_k$ decays. Note that, if the cascading level $\Delta_k J_k$ is not aligned, its contribution to the sum in Eq. (13) does not depend on M_i and can be expressed as $R(\Delta_k J_k \rightarrow \Delta_i J_i)\Omega_{\rm eff}(\Delta_j J_j \rightarrow \Delta_k J_k)/(2J_i+1)$.

III. METHOD OF CALCULATION

For autoionization transition-matrix elements, the method of Refs. [8,11,13] has been used. Here only a brief description of the method is presented. The wave functions for the He-like and Li-like states are obtained from the SUPERSTRUCTURE code of Eissner, Jones, and Nussbaumer [15]. In the present work, the configurations included for He-like states are $1s^2$, $1s2s$, $1s2p$, $1s3s$, $1s3p$, and $1s3d$, and for Li-like states, $1s^22s$, $1s^22p$, $1s^23s$, $1s^23p$, $1s^23d$, $1s2s^2$, $1s2s2p$, $1s2p^2$, $1s2s3s$, $1s2s3p$, $1s2s3d$, $1s2p3s$, $1s2p3p$, $1s2p3d$, $1s3s^2$, $1s3s3p$, $1s3s3d$, $1s3p^2$, $1s3p3d$, and $1s3d^2$. The one-electron orbitals nl are calculated in a *scaled* Thomas-Fermi-Dirac-Amaldi potential, with scaling parameters λ_l different for each orbital quantum number l. Multiconfigurational wave functions are first constructed in LS coupling by diagonalizing the nonrelativistic Hamiltonian. Then the Hamiltonian matrix is computed, including the Breit-Pauli relativistic corrections, which are important for highly charged ions such as Fe XXIV and Fe XXV. The eigenlevel wave function for the ΔJ level is expanded on the nonrelativistic eigenterm wave functions $|\Gamma_k L_k S_k J\rangle$,

$$|\Delta J\rangle = \sum_k f_{\Delta J}(\Gamma_k L_k S_k) |\Gamma_k L_k S_k J\rangle \; , \quad (15)$$

where the mixing coefficients $f_{\Delta J}(\Gamma_k L_k S_k)$ are obtained by diagonalizing the Breit-Pauli Hamiltonian matrix. These wave functions are then utilized in SUPERSTRUCTURE to calculate the radiative transition probabilities, branching ratios R, and the radiative part of the total widths Γ_d in the denominator of $\tilde{F}^d_2$. All the radiative transitions to lower levels are included in the calculation of R and Γ^r_d, i.e., 3-3, 3-2, and 3-1 radiative rates. In the computation of autoionization matrix elements, the continuum orbitals kl are generated in a distorted-wave approximation. The calculations are carried out by means of the DISTWAV code developed by Eissner and Seaton [16] and modified by one of us (J.D.) to give LS autoionization matrix elements. These matrix elements as well as $f_{\Delta J}(\Gamma_k L_k S_k)$ are real by the AUTOLSJ program to give LSJ fine-structure autoionization matrix elements.

In the "resonance" part of this paper, to have orthogonal functions we use for l orbitals the same potential to compute the Fe XXIV and Fe XXV nl bound orbitals and εl free orbitals. This is quite valid for highly charged systems. The normalization of the continuum wave function is chosen such that its radial part behaves asymptotically as

$$F_{kl}(r) \underset{r\rightarrow\infty}{\sim} (k)^{-1/2} \sin\left[kr - l\frac{\pi}{2} + \frac{z}{k}\ln(2kr) + \sigma_l + \tau_l\right] , \quad (16)$$

where $z = Z - N = 24$ is the charge of the target ion. The phase shift τ_l due to short-range potentials is almost zero for highly ionized atoms. The energy $\varepsilon = k^2/2$ of the continuum electron is determined by the energy difference between that of the $(N+1)$-electron autoionizing term and that of the N-electron target term, according to energy conservation.

Now for the $1s^2$-$1s3l$ sublevel collision strengths, the method described in paper I has been followed: the DISTWAV code has been used in its original version.

IV. RESULTS AND DISCUSSION

A. Resonance effects

The major resonance effect to the $1s^2$-$1s2l$ excitations is due to the autoionizing Fe XXIV levels issued from

TABLE I. Values of $\tilde{F}_2^d(\Delta_i J_i)$ summed over the five main contributing $1s3l3l'$ resonances with $J_d=\frac{1}{2}$ and divided by the statistical weight $(2J_i+1)$ for each of the Fe XXV $1s^2$-$1s2l$ $\Delta_i J_i \geq 1$ excitations. The five resonances are (1) $1s3s(^1S)$ $3p\ ^2P_{1/2}$, (2) $1s3s(^3S)$ $3p\ ^2P_{1/2}$, (3) $1s3p^2\ ^2S_{1/2}$, (4) $1s3p3d\ ^2P_{1/2}$, and (5) $1s3s^2\ ^2S_{1/2}$. Tabulated values for $\tilde{F}_2^d$ have to be multiplied by 10^{11} s^{-1}.

$\Delta_i J_i$	
$1s2p\ ^3P_2$	4.200
$1s2p\ ^1P_1$	8.614
$1s2s\ ^3S_1$	34.63
$1s2p\ ^3P_1$	8.730

$1s3lnl'$ configurations converging to the $1s3l$ ionization limits [8]. We restrict ourselves to the lowest $1s3l3l'$ levels which lie at energies between 511.0 and 515.6 Ry above the Fe XXV ground state. Among the $1s3lnl'$ resonances, their contribution is dominant [8]. These autoionizing levels are well separated in energy and therefore no interference between them appears.

In intermediate coupling, the $1s3l3l'$ configurations consist of 63 doubly excited levels but the computations show that the dominating contributions to $1s^2 \rightarrow 1s2l$ excitations in Fe XXV come only from 17 levels. Among them, five have total angular momentum $J_d=\frac{1}{2}$. As mentioned above, the resonances with $J_d=\frac{1}{2}$ populate equally the various $\Delta_i J_i$ magnetic sublevels. They do not contribute to the alignment of the $1s2l$ levels and separately they do not induce any polarization of the lines issued from these $1s2l$ levels. It is then convenient in Table I to present the sum over the five resonances $J_d=\frac{1}{2}$ of $\tilde{F}_2^d(\Delta_i J_i)$ divided by the statistical weight $(2J_i+1)$ for each of the excitations, i.e., an equal contribution to each sublevel M_i of $1s2l\ \Delta_i J_i$ level. In the following sections, results will be presented only for the other 12 resonances with $J_d \geq \frac{3}{2}$, the only ones which contribute significantly to the alignment of the $1s2l$ levels.

1. x line: $1s^2\ ^1S_0 - 1s2p\ ^3P_2$

In Table II is listed, for the individual 12 main $1s3l3l'$ resonances with $J_d \geq \frac{3}{2}$, the average polarization degree $\bar{\eta}_3^d$ of line x, the magnetic quadrupole decay of $1s2p\ ^3P_2$ to $1s^2\ ^1S_0$. Also listed are values of the $\tilde{F}_2^d(1s2p\ ^3P_2$ $M_i=1/2)$ defined in Eq. (8) and the individual partial-wave contributions to the autoionization probability

TABLE II. Average linear polarization degree $\bar{\eta}_2^d$ of the Fe XXV x line, following resonance excitation via the main $1s3l3l'$ autoionizing levels with $J_d \geq \frac{3}{2}$. The energies E_d (in rydbergs) of the autoionizing states are also given relative to the Fe XXV ground level. Tabulated values for $A_{l_i}^a$ and $\tilde{F}_2^d$ have to be multiplied by 10^{13} s^{-1} and 10^{11} s^{-1}, respectively.

Autoionizing		$\sum_{M_d=\pm 1/2}$	$A_{l_i}^a(d\ M_d \rightarrow 1s2p\ ^3P_2\ M_i)$		$\tilde{F}_2^d(1s2p\ ^3P_2\ M_i)$		$\bar{\eta}_3^d(x)$
level d	E_d (Ry)	l_i	$M_i=1$	$M_i=2$	$M_i=1$	$M_i=2$	(%)
$1s3s3p\ ^2P_{3/2}$	512.560	0	11.63	0.0	19.61	2.398	−78.2
		2	0.614	1.498			
$1s3s3d\ ^2D_{3/2}$	512.399	1	0.851	0.204	2.766	1.021	−46.1
		3	0.218	0.190			
$1s3s3d\ ^2D_{5/2}$	512.367	1	3.244	0.149	12.61	0.985	−85.5
		3	0.033	0.105			
$1s3p^2\ ^2D_{5/2}$	513.867	1	8.338	1.365	11.68	3.264	−56.3
		3	0.481	1.100			
$1s3p^2\ ^2D_{3/2}$	513.862	1	1.481	1.192	2.860	2.099	−15.4
		3	1.810	1.223			
$1s3p3d\ ^2F_{5/2}$	514.918	2	2.990	4.102	3.818	3.881	+0.8
		4	2.700	1.682			
$1s3p3d\ ^2F_{7/2}$	514.894	2	20.20	4.712	17.82	5.398	−53.5
		4	0.692	1.618			
$1s3p3d\ ^2F_{7/2}$	513.760	2	0.313	0.104	0.660	0.249	−45.2
		4	0.007	0.016			
$1s3p3d\ ^2F_{5/2}$	513.529	0	0.042	0.0	0.195	0.211	+3.8
		2	0.120	0.195			
		4	0.044	0.028			
$1s3p3d\ ^2P_{3/2}$	515.577	0	1.173	0.0	1.231	1.881	+20.9
		2	2.432	5.510			
$1s3d^2\ ^2D_{5/2}$	515.436	3	2.050	4.938	3.046	7.337	+41.3
$1s3d^2\ ^2D_{3/2}$	515.440	3	7.082	5.296	6.944	5.193	−14.4
$\sum_d \tilde{F}_2^d$					83.24	33.92	

TABLE III. Different (l_i and K_i) channel contributions to the total autoionization transition probabilities from $\Delta_d J_d$ to $\Delta_i J_i$. The asterisk denotes that some reduced matrix elements with the same l_i but different K_i have opposite sign. Tabulated values for $A^a_{l_i K_i}$ have to be multiplied by 10^{12} s^{-1}.

$\Delta_d J_d$-$\Delta_i J_i$	l_i	K_i	$A^a_{l_i K_i}(\Delta_d J_a - \Delta_i J_i)$
$1s3s3d\ ^2D_{5/2}-1s2p\ ^3P_2$	1	3	71.74*
	1	2	1.400
	3	3	1.727
$1s3p3d\ ^2F_{5/2}-1s2p\ ^1P_1$	2	3	130.2
	4	3	22.92
$1s3p3d\ ^2F_{7/2}-1s2p\ ^1P_1$	2	3	134.0
	4	3	23.55
$1s3p3d\ ^2F_{7/2}-1s2p\ ^1P_1$	2	3	4.227
	4	3	0.585
$1s3p3d\ ^2F_{5/2}-1s2p\ ^1P_1$	2	3	4.121
	2	2	1.258*
	4	3	0.655

$\sum_{M_d=\pm 1/2} A^a(dM_d \to 1s2p\ ^3P_2\ M_i=1,2)$. It is not necessary to tabulate $\tilde{F}^d_2(2\,^3P_2\ M_i=0)$, since it can be shown, with some angular-momentum algebra, that

$$\tilde{F}^d_2(2\,^3P_2\ M_i=0)=[4\tilde{F}^d_2(2\,^3P_2\ M_i=1) -\tilde{F}^d_2(2\,^3P_2\ M_i=2)]/3 \ . \tag{17}$$

In calculating $\bar{\eta}^d_3(x)$ by the expression

$$\bar{\eta}^d_3(x)=\frac{\tilde{F}^d_2(2\,^3P_2\ M_i=2)-\tilde{F}^d_2(2\,^3P_2\ M_i=1)}{\tilde{F}^d_2(2\,^3P_2\ M_i=2)+\tilde{F}^d_2(2\,^3P_2\ M_i=1)} \ , \tag{18}$$

the line is assumed to be observed at right angles to the incident electron beam; see Eq. (28) of paper I.

We first note that certain resonances excite preferentially the sublevel $M_i=1$ and others the sublevel $M_i=2$, leading to a change of the sign of the polarization degree for the line x. However, taking into account the weighted alignment from the different resonances, it appears that the $M_i=1$ sublevel is on average much more preferentially excited. The largest cross section for exciting the $M_i=1$ sublevel relative to exciting $M_i=2$ is due to the even-parity $1s3s3d\ ^2D_{5/2}$ autoionizing state. This resonance induces a negative polarization degree as high as -85.5%. Comparatively, the background (nonresonant) polarization is at the same impact energy about -52%, according to paper I. To understand this surprisingly large polarization, we refer to Table III. Here we show separately the contributions to $A^a(1s3s3d\ ^2D_{5/2}-1s2p\ ^3P_2)$ of each channel with given l_i and K_i. It can be seen that for the partial wave p, which is dominant, there is no need to take into account the interferences between channels $K_i=3$ and $K_i=2$ [represented by the cross terms of different K_i in Eq. (5)], due to the negligible contribution of the latter channel. The next highest degree of polarization amounts to -78.2% and occurs for the odd-parity $1s3s3p\ ^2P_{3/2}$ resonance. This large value is not surprising since, as shown in Table II, it results from the domination of the s wave in the autoionization decay probability, this wave resulting in the excitation of only the $|M_i|\le 1$ sublevels. The strongest depolarization of line x is caused by the resonance $1s3d^2\ ^2D_{5/2}$ which overpopulates the $M_i=2$ sublevel by a factor of 2.4 relative to $M_i=1$. For this resonance the main contribution to the autoionization comes from the f wave. The average polarization degree $\bar{\eta}_3(x)$ taken over the 12 resonance $J_d \ge \frac{3}{2}$ is found close to -42%. It decreases to $\sim -39\%$ when including the five unpolarizing resonances $J_d=\frac{1}{2}$ (Table I).

In Table IV we presented the direct and resonant contributions to excitation rates for the $1s2l\ J_i \ge 1$ magnetic sublevels. These results were calculated by assuming the electron energy distribution $f(\varepsilon_j)$ to be constant over the energy interval 511–516 Ry, which just covers the $1s3l3l'$ resonances, and to be zero outside this interval. Also given are the polarization degrees η_3 of the corresponding lines obtained without including and including the resonant contribution. The results for the x line show that the collisional resonances do not have a severe depolarizing effect although they contribute as much as the direct excitation to the upper level population. For a thermal electron distribution at temperatures in the range of interest for Fe XXV, the increase of the direct excitation rate of x by resonances is not substantial (less than 15%). It is clear that the resonance effects on the polarization of x may be quite small for plasmas in which anisotropic electron flux has an energy distribution very large compared to the region containing the resonances.

TABLE IV. Direct and resonant contributions to excitation rates, $C_{\rm dir}$ and $C_{\rm res}$, respectively, for $1s^2\ ^1S_0-1s2l\ \Delta_i J_i M_i$ transitions (in 10^{-13} cm^3 s^{-1}). The electron energy distribution is assumed to be constant in the interval 511–516 Ry, which just contains the $1s3l3l'$ resonances, and zero elsewhere. Also shown is the degree η_3 of polarization of the corresponding lines without (upper entries) and with (lower entries) inclusion of resonances.

	$1s^2\ ^1S_0-1s2p\ ^3P_2$ (x)		$1s^2\ ^1S_0-1s2p\ ^1P_1$ (w)		$1s^2\ ^1S_0-1s2p\ ^3P_1$ (y)		$1s^2\ ^1S_0-1s2s\ ^3S_1$ (z)	
	$M_i=1$	$M_i=2$	$M_i=0$	$M_i=1$	$M_i=0$	$M_i=1$	$M_i=0$	$M_i=1$
$C_{\rm dir}$	2.31	0.726	12.7	3.21	1.70	2.39	0.959	0.959
$C_{\rm res}$	2.26	0.987	2.82	1.42	1.77	1.86	2.851	2.864
η_3 (%)	-52.5		59.6		-16.9		0.0	
	-45.5		54.0		-10.1		0.2	

2. *w line:* $1s^2\,{}^1S_0-1s2p\,{}^1P_1$

Results of the calculations for the transition $1s^2\,{}^1S_0-1s2p\,{}^1P_1$ corresponding to w line are presented in Table V, which is similar to Table II. It is seen, in contrast to the previous table, that all resonances populate preferentially the sublevel $M_i=0$. Moreover, it can be seen that the s partial wave causes the highest alignment of the $1s2p\,{}^1P_1$ level, which results in a polarization of 60% for w. Unfortunately, this wave does not contribute dominantly to the total autoionization probability. It can even be neglected with respect to the d wave for the resonance $1s3p3d\,{}^2P_{3/2}$. The largest polarization found on w is due to the $1s3s3d\,{}^2D_{3/2}$ resonance, and is slightly lower than the background polarization, which is $\sim 59\%$ in the energy region under consideration; see paper I. It is interesting to note the unexpectedly substantial polarization of w induced by the four resonances $1s3p3d\,{}^2F_{5/2,7/2}$ (about 40%) for which the s wave does not occur because of angular momentum coupling rules. In Table III the results indicate that for the three first resonances only one value of K_i ($K_i=3$) dominantly contributes to $l_i=2$ and 4 waves. For the fourth resonance there are constructive interference effects between channels $K_i=2$ and $K_i=3$ in the d wave, owing to the opposite sign of the corresponding reduced matrix elements. The average polarization degree $\bar{\eta}_3(w)$ taken over the 17 resonances $J_d \geq \frac{1}{2}$ is found to be close to 37%. We note that the five unpolarizing resonances $J_d=\frac{1}{2}$ make a contribution of only 12% to the total $\sum_d \tilde{F}_2^d(1s2p\,{}^1P_1)$.

It can be seen from Table IV that the effect of neglecting the collisional resonances on the polarization of w is $\sim 10\%$. We expect such effect to be very weak for applications to plasma diagnostic of nonthermal electrons.

3. *y and z lines:* $1s^2\,{}^1S_0-1s2p\,{}^3P_1$ *and* $1s^2\,{}^1S_0-1s2s\,{}^3S_1$ *(respectively)*

Table VI lists for the values of $\tilde{F}_2^d(M_i=0,1)$ the relative populations in magnetic sublevels of both $1s2p\,{}^3P_1$ and $1s2s\,{}^3S_1$, following excitation via the 12 resonances $J_d \geq \frac{3}{2}$. The resonances are seen to be weakly selective in populating the sublevels of $1s2s\,{}^3S_1$, as might be expected because of the isotropy associated with an S level. It is to be recalled that the direct-excitation cross sections of these sublevels are strictly identical; see paper I. The examination of the autoionization data reveals that most resonances with $J_d \geq \frac{3}{2}$ decay to the continuum $1s2s\,{}^3S_1$, ε_i predominantly through the d (f) wave rather than the s (p) wave. In addition, the contributions to an l_i wave from the channels $K_i=J_d\pm\frac{1}{2}$, which are often compara-

TABLE V. The same as Table II but for the Fe XXV w line. Tabulated values for $A_{l_i}^a$ and $\tilde{F}_2^d$ have to be multiplied by $10^{13}\ \mathrm{s}^{-1}$ and $10^{11}\ \mathrm{s}^{-1}$, respectively.

Autoionizing level d	$\sum_{M_d=\pm 1/2}$ l_i	$A_{l_i}^a(d\ M_d \to 1s2p\,{}^1P_1\ M_i)$ $M_i=0$	$M_i=1$	$\tilde{F}_2^d(1s2p\,{}^1P_1\ M_i)$ $M_i=0$	$M_i=1$	$\bar{\eta}_3^d(w)$ (%)
$1s3s3p\,{}^2P_{3/2}$	0	5.000	1.250	11.68	5.177	+38.6
	2	2.298	1.984			
$1s3s3d\,{}^2D_{3/2}$	1	4.648	1.305	12.15	3.466	+55.6
	3	0.047	0.035			
$1s3s3d\,{}^2D_{5/2}$	1	4.040	1.347	15.77	5.307	+49.6
	3	0.033	0.024			
$1s3p^2\,{}^2D_{5/2}$	1	6.876	2.292	12.10	5.196	+39.9
	3	2.264	1.632			
$1s3p^2\,{}^2D_{3/2}$	1	5.158	3.942	6.183	4.686	+13.8
	3	1.956	1.467			
$1s3p3d\,{}^2F_{5/2}$	2	14.63	5.704	11.14	4.706	+40.6
	4	1.965	1.310			
$1s3p3d\,{}^2F_{7/2}$	2	15.31	5.742	14.80	6.038	+42.0
	4	2.038	1.336			
$1s3p3d\,{}^2F_{7/2}$	2	0.482	0.181	1.103	0.443	+42.7
	4	0.051	0.033			
$1s3p3d\,{}^2F_{5/2}$	2	0.610	0.233	0.630	0.256	+42.2
	4	0.056	0.037			
$1s3p3d\,{}^2P_{3/2}$	0	0.379	0.095	2.232	1.740	+12.4
	2	6.158	5.002			
$1s3d^2\,{}^2D_{5/2}$	3	5.070	3.656	7.534	5.432	+16.2
$1s3d^2\,{}^2D_{3/2}$	3	4.890	3.668	4.795	3.596	+14.3
$\sum_d \tilde{F}_2^d$				100.1	46.05	

ble, interfere destructively for all resonances. These are quite consistent with the very weak efficiency with which the autoionizing level alignment is transferred to $^3S_1$. For each of the resonances, the average polarization degree $\overline{\eta}_3^d$ of line z emitted in the decay $1s2s\ ^3S_1 \rightarrow 1s^2\ ^1S_0$ does not exceed 15% in absolute value. As long as cascade effects are ignored, the polarization of z remains practically zero when the average effect of all the resonances is taken into consideration. In Table IV, it should be noted that for z the resonance contribution is about three times larger than the *direct* excitation rate.

In examining Table VI for $1s2p\ ^3P_1 - 1s^2\ ^1S_0$, we note that the alignments produced by the resonances give rise to a polarization for the intercombination line y that is often small, either positive or negative. With the exception of the very weak resonance $1s3p3d\ ^2F_{5/2}$, it can be seen that the largest alignment of the $1s2p\ ^3P_1$ level is, as for $1s2p\ ^3P_2$, due to the $1s3s3d\ ^2D_{5/2}$ resonance. We mention that the relative smallness of the alignment transfer of the autoionizing levels to $1s2p\ ^3P_1$ can be attributed to the spin-orbit coupling with $1s2p\ ^1P_1$. The average polarization degree $\overline{\eta}_3(y)$ taken over the 17 resonances is found to be about -2.5%. Comparatively, the background polarization degree is, according to paper I, about -17% in the energy region considered. Adopting for the electron beam an energy spread $\Delta E = 5$ Ry large enough to cover the resonances, the resulting polarization of y decreases to $\sim -10\%$; see Table IV. However, for a broader energy distribution, the average background polarization may be expected to be nearly zero.

B. Cascade effects

We have restricted calculations of the cascade effects to the ten $1s3l$ levels. Of course, the important cascades from the higher $1s2p\ ^3P_{0,2}$ levels into the $1s2s\ ^3S_1$ level are also taken into account. The $1s2s\ ^3S_1$ level is by far the most affected by cascades. In Table VII are listed the results of collision strengths for the transitions $1s^2$-$(1s3l)\ \Delta_k J_k M_k$ in Fe XXV ($M_i = 0,1,2$ in the table) as well as the total collision strengths to the $\Delta_k J_k$ levels (i.e., results summed over M_k, denoted $M_i = $T in the Table). Also given in these tables are the calculated transition energies in rydbergs. To our knowledge for the magnetic sublevels there are no results available for comparison. For the level-to-level collision strengths we can compare with the Coulomb-Born-exchange calculation of Sampson, Goett, and Clark [17] (denoted $M_i = $S) and the distorted-wave calculation of Mann [18] (denoted $M_i = $M). Since their calculations were not carried out at the same energies as ours we used a linear interpolation to derive comparable data. In the case of Sampson, Goett, and Clark, we used their Tables V and VII–X as well as formula (23). Mann does not tabulate fine-structure collision strengths but collision strengths summed over the initial and final levels. In the case considered here, i.e., the He-like case, it is not difficult to come back to the original fine-structure data by using the mixing coefficients tabulated in Mann's Table II, pp. 447 and 448, as well as his collision strengths tabulated in Table I, p. 431. As can be seen from Table VII, our results are in a fairly good agreement with their values. It can be noted that the degree of alignment of the $1s3p\ ^1P_1$ level is nearly similar to the degree of the $1s2p\ ^1P_1$ (for incident energies in threshold units). This is also valid for the $1s3p\ ^3P_2$ level, but not for $1s3p\ ^3P_1$, which appears to be more aligned than $1s2p\ ^3P_1$. The reason for this is that the effect of intermediate coupling on the $1s^2\ ^1S_0 - 1snp\ ^3P_1$ collision strength is smaller for $n = 3$ than for $n = 2$. The calculated branching ratios for the cascade transitions of interest are shown in Table VIII. The cascade effects between the $n = 3$ levels can be neglected since the radiative transition probabilities between these levels are very weak. The cascade effects from $n \geq 4$ are not negligible but nevertheless they are much less important than the cascades between the $n = 2$ and from the $n = 3$ levels.

TABLE VI. Values of $\tilde{F}_2^d(\Delta_i J_i M_i = 0,1)$ (in 10^{11} s^{-1}) obtained for $\Delta_i J_i = 1s2s\ ^3S_1$ and $1s2p\ ^3P_1$ due to the major contributing $1s3l3l'$ resonances with $J_d \geq \frac{3}{2}$.

Autoionizing	$\tilde{F}_2^d(1s2s\ ^3S_1\ M_i)$		$\tilde{F}_2^d(1s2p\ ^3P_1\ M_i)$	
state d	$M_i=0$	$M_i=1$	$M_i=0$	$M_i=1$
$1s3s3p\ ^2P_{3/2}$	9.462	11.80	8.102	5.219
$1s3s3d\ ^2D_{3/2}$	11.01	12.08	6.595	7.879
$1s3s3d\ ^2D_{5/2}$	19.53	16.24	16.98	6.662
$1s3p^2\ ^2D_{5/2}$	13.71	14.07	4.189	3.822
$1s3p^2\ ^2D_{3/2}$	9.242	9.154	3.870	7.693
$1s3p3d\ ^2F_{5/2}$	3.987	3.984	6.035	13.20
$1s3p3d\ ^2F_{7/2}$	5.155	5.161	6.366	4.363
$1s3p3d\ ^2F_{7/2}$	0.262	0.347	negligible	negligible
$1s3p3d\ ^2F_{5/2}$	0.189	0.184	0.088	0.615
$1s3p3d\ ^2P_{3/2}$	negligible	negligible	1.160	2.016
$1s3d^2\ ^2D_{5/2}$	1.740	1.796	4.390	10.21
$1s3d^2\ ^2D_{3/2}$	1.161	1.148	1.913	1.434
$\sum_d \tilde{F}_2^d$	75.45	75.96	59.69	63.11

1. $1s^2$-$1s2p$ transitions (w,x,y lines)

The effects of cascades on the magnetic sublevels of the three $(1s2p)$ $J=1$ and 2 levels are shown in Table IX. This gives a comparison between the contributions to the collision strengths due to the direct excitation (column A) and due to cascade transitions from the $n=3$ levels (column B) for energies from the $n=3$ threshold (582 Ry)

TABLE VII. Collision strengths versus incident electron energy k_j^2 for transitions from $1s^2\,{}^1S_0$ ground level ($n=1$) to the ($n=3$) $\Delta_i J_i M_i$ sublevels ($M_i=0,1,2$). The collision strengths for $1s3s\,{}^3S_1$ $M_i=0,1$ are not listed because they are equal. For the level-to-level transitions we compare present results (M_i=T) to those of Sampson, Goett, and Clark [17] (M_i=S) and Mann [18] (M_i=M). The threshold energies ΔE are shown. Numbers in brackets denote powers of 10.

Excited state			k_j^2 (Ry)					
$\Delta_i J_i$	M_i	ΔE (Ry)	582	700	900	1200	2000	500
$1s3s\,{}^3S_1$	T	578.64	9.30[−5]	7.62[−5]	5.52[−5]	3.66[−5]	1.64[−5]	3.33[−6]
	S		9.61[−5]	7.62[−5]	5.48[−5]	3.66[−5]	1.76[−5]	3.48[−6]
	M		9.27[−5]	7.51[−5]	5.46[−5]	3.64[−5]	1.65[−5]	3.40[−6]
$1s3p\,{}^3P_2$	0		9.51[−5]	7.11[−5]	4.48[−5]	2.52[−5]	8.25[−6]	9.25[−7]
${}^3P_2$	1		7.71[−5]	5.74[−5]	3.63[−5]	2.05[−5]	6.92[−6]	8.65[−7]
${}^3P_2$	2		2.34[−5]	1.66[−5]	1.06[−5]	6.51[−6]	2.93[−6]	6.84[−7]
${}^3P_2$	T	579.59	2.96[−4]	2.19[−4]	1.39[−4]	7.93[−5]	2.80[−5]	4.02[−6]
	S		2.95[−4]	2.14[−4]	1.38[−4]	8.17[−5]	3.29[−5]	4.71[−6]
	M		2.97[−4]	2.17[−4]	1.39[−4]	8.02[−5]	2.87[−5]	3.93[−6]
$1s3p\,{}^3P_1$	0		4.35[−5]	4.77[−5]	5.54[−5]	6.49[−5]	7.85[−5]	8.95[−5]
${}^3P_1$	1		7.45[−5]	5.95[−5]	4.50[−5]	3.75[−5]	4.10[−5]	7.26[−5]
${}^3P_1$	T	579.28	1.92[−4]	1.67[−4]	1.45[−4]	1.40[−4]	1.61[−4]	2.35[−4]
	S		1.96[−4]	1.67[−4]	1.48[−4]	1.44[−3]	1.65[−4]	2.40[−4]
	M		2.00[−4]	1.73[−4]	1.55[−4]	1.53[−3]	1.79[−4]	2.63[−4]
$1s3p\,{}^3P_0$	T	579.20	5.80[−5]	4.32[−5]	2.74[−5]	1.57[−5]	5.57[−6]	8.04[−7]
	S		5.89[−5]	4.28[−5]	2.76[−5]	1.63[−5]	6.58[−6]	9.41[−7]
	M		5.95[−5]	4.35[−5]	2.78[−5]	1.60[−5]	5.74[−6]	7.85[−7]
$1s3s\,{}^1S_0$	T	579.19	1.34[−4]	1.59[−4]	1.85[−4]	2.10[−4]	2.41[−4]	2.73[−4]
	S		1.37[−4]	1.56[−4]	1.78[−4]	1.99[−4]	2.25[−4]	2.51[−4]
	M		1.40[−4]	1.58[−4]	1.78[−4]	1.99[−4]	2.24[−4]	2.48[−4]
$1s3d\,{}^3D_3$	0		1.86[−5]	1.18[−5]	5.90[−6]	2.52[−6]	5.07[−7]	3.07[−8]
${}^3D_3$	1		1.47[−5]	9.39[−6]	4.85[−6]	2.19[−6]	5.10[−7]	4.05[−8]
${}^3D_3$	2		5.81[−6]	3.88[−6]	2.35[−6]	1.32[−6]	4.57[−7]	5.70[−8]
${}^3D_3$	3		2.40[−7]	2.99[−7]	3.30[−7]	2.86[−7]	1.64[−7]	4.10[−8]
${}^3D_3$	T	580.02	6.01[−5]	3.89[−5]	2.10[−5]	1.01[−5]	2.77[−6]	3.08[−7]
	S		6.33[−5]	4.09[−5]	2.27[−5]	1.15[−5]	3.81[−6]	4.08[−7]
	M		6.04[−5]	3.92[−5]	2.14[−5]	1.04[−5]	2.85[−6]	3.86[−7]
$1s3d\,{}^3D_2$	0		1.08[−5]	9.20[−6]	8.24[−6]	7.32[−6]	5.67[−6]	6.01[−6]
${}^3D_2$	1		1.10[−5]	8.72[−6]	7.80[−6]	8.62[−6]	1.10[−5]	1.13[−5]
${}^3D_2$	2		2.24[−6]	1.86[−6]	2.04[−6]	2.85[−6]	5.40[−6]	1.12[−5]
${}^3D_2$	T	579.90	3.74[−5]	3.04[−5]	2.79[−5]	3.03[−5]	3.85[−5]	5.10[−5]
	S		3.45[−5]	2.88[−5]	2.74[−5]	3.05[−5]	3.90[−5]	5.26[−5]
	M		3.66[−5]	3.02[−5]	2.83[−5]	3.11[−5]	4.03[−5]	5.31[−5]
$1s3d\,{}^3D_1$	0		1.51[−5]	9.69[−6]	5.04[−6]	2.29[−6]	5.39[−7]	4.21[−8]
${}^3D_1$	1		5.26[−6]	3.48[−6]	1.97[−6]	1.02[−6]	3.24[−7]	4.49[−8]
${}^3D_1$	T	579.89	2.56[−5]	1.67[−5]	8.98[−6]	4.34[−6]	1.19[−6]	1.32[−7]
	S		2.71[−5]	1.75[−5]	9.73[−6]	4.93[−6]	1.63[−6]	1.75[−7]
	M		2.59[−5]	1.68[−5]	9.16[−6]	4.44[−6]	1.22[−6]	1.65[−7]
$1s3d\,{}^1D_1$	0		1.55[−5]	1.40[−5]	1.39[−5]	1.34[−5]	1.12[−5]	1.24[−5]
${}^1D_2$	1		7.36[−6]	8.30[−6]	1.13[−5]	1.58[−5]	2.24[−5]	2.34[−5]
${}^1D_2$	2		1.34[−6]	1.60[−6]	2.78[−6]	5.04[−6]	1.08[−5]	2.31[−5]
${}^1D_2$	T	580.03	3.29[−5]	3.38[−5]	4.21[−5]	5.51[−5]	7.76[−5]	1.05[−4]
	S		2.27[−5]	2.83[−5]	3.97[−5]	5.48[−5]	7.88[−5]	1.10[−4]
	M		3.01[−5]	3.23[−5]	4.11[−5]	5.45[−5]	7.80[−5]	1.05[−4]
$1s3p\,{}^1P_1$	0		2.42[−4]	3.56[−4]	5.00[−4]	6.44[−4]	8.28[−4]	9.78[−4]
${}^1P_1$	1		6.59[−5]	9.01[−5]	1.36[−4]	2.07[−4]	3.79[−4]	7.91[−4]
${}^1P_1$	T	579.95	3.74[−4]	5.37[−4]	7.73[−4]	1.06[−3]	1.59[−3]	2.56[−3]
	S		2.76[−4]	5.34[−4]	7.64[−4]	1.04[−3]	1.54[−3]	2.49[−3]
	M		3.78[−4]	5.32[−4]	7.59[−4]	1.03[−3]	1.54[−3]	2.46[−3]

TABLE VIII. Fe XXV radiative branching ratios R for radiative decay to $1s2s\,^3S_1$ and $1s2p\,^3P_{(0,1,2)},^1P_1$ levels.

	Branching ratio to:				
From:	$1s2s\,^3S_1$	$1s2p\,^3P_0$	$1s2p\,^3P_1$	$1s2p\,^3P_2$	$1s2p\,^1P_1$
$1s2p\,^3P_0$	1.0				
$1s2p\,^3P_2$	0.180				
$1s3s\,^3S_1$		0.103	0.292	0.575	0.029
$1s3s\,^1S_0$		0.0	0.063	0.0	0.936
$1s3p\,^3P_0$	1.0				
$1s3p\,^3P_1$	0.408				
$1s3p\,^3P_2$	1.0				
$1s3d\,^3D_1$		0.557	0.385	0.027	0.031
$1s3d\,^3D_2$		0.0	0.708	0.166	0.126
$1s3p\,^1P_1$	0.005 36				
$1s3d\,^3D_3$		0.0	0.0	1.0	0.0
$1s3d\,^1D_2$		0.0	0.067	0.084	0.849

to 5000 Ry. The cascade effects are seen to be very weak for the $1s2p\,^1P_1$ $M_i=0$ sublevel and to increase the collision strengths for the $1s2p\,^3P_1$ $M_i=0$, $1s2p\,^{1,3}P_1$ $M_i=1$ sublevels by less than 13%. As a consequence, the cascades have a negligibly small effect on the polarization of the w and y lines. This is clearly seen in Table X, where the polarization degree (in percent) as a function of impact energy is given. In the case of the y line, one can notice that near the $n=3$ threshold the $M_i=0$ sublevel is preferently populated relative to $M_i=1$ by the cascades from $1s3d\,^1D_2$ and $1s3d\,^3D_2$, in contrast to direct excitation. Now for the $1s2p\,^3P_2$ level, the cascade contribution is substantial for the $M_i=2$ sublevel near the $n=3$ excitation threshold, and becomes important for all the sublevels at energies larger than approximately three times this threshold. Nevertheless, the cascades do not have a severe depolarization effect on the x line at not too high impact energies since they cause an average alignment on the $1s2p\,^3P_2$ level with the same sign as the one due to direct excitation. In Table X, it is seen that the polarization of x is less negative by only 3–6 % at energies in the range 582–1200 Ry. Near the $n=3$ threshold the main cascades contributions come from $1s3s\,^3S_1$ and $1s3s\,^3D_3$. The population ratio between the $1s2p\,^3P_2$ $M=1$ and $M=2$ sublevels, due to the latter cascade, is found to be close to 4.9, which is much greater than the ratio 3.3 due to direct excitation. Therefore the cascade from $1s3d\,^3D_3$ has the effect to increase the polarization of x, but the more important cascade comes from the unaligned $1s3s\,^3S_1$ level and tends to decrease this polarization. At high energies the alignment due to cascades changes sign, but the collision strengths become vanishingly small at these energies.

2. z line: $1s^2\,^1S_0-1s2s\,^3S_1$

The calculated direct and cascade contributions to the collision strengths, $\Omega_{\rm dir}$ and $\Omega_{\rm cas}$, for the $1s^2\,^1S_0-1s2s\,^3S_1$ $M_i=0$ and 1 magnetic transitions are presented in Table XI together with the degree of polarization of the z line induced by cascades. At electron energies in the range 495–578 Ry the cascade effects from

TABLE IX. Contributions to the collision strengths for populating magnetic sublevels of Fe XXV $1s2p\,^3P_{1,2}$ and $^1P_1$ by direct collisional excitation from the ground level (column A) and by radiative cascade transitions from higher $n=3$ levels (column B). Numbers in brackets indicate powers of 10.

Excited		k_j^2 (Ry)											
		582		700		900		1200		2000		5000	
level	M_i	A	B	A	B	A	B	A	B	A	B	A	B
$1s2p\,^3P_1$	0	2.11[−4]	2.80[−5]	2.38[−4]	2.38[−4]	2.88[−4]	2.08[−5]	3.38[4]	1.96[−5]	4.00[−4]	1.93[−5]	4.58[−4]	1.91[−5]
	1	2.42[−4]	2.31[−5]	1.98[−4]	1.94[−5]	1.67[−4]	1.65[−5]	1.58[−4]	1.56[−5]	2.06[−4]	1.68[−5]	3.80[−4]	2.12[−5]
$1s2p\,^3P_2$	0	2.76[−4]	3.63[−5]	2.02[−4]	2.56[−5]	1.29[−4]	1.61[−5]	7.44[−5]	1.03[−5]	2.52[−5]	6.32[−6]	3.00[−6]	4.28[−6]
	1	2.24[−4]	2.84[−5]	1.64[−4]	2.07[−5]	1.05[−4]	1.36[−5]	6.10[−5]	8.72[−6]	2.13[−5]	4.73[−6]	2.83[−6]	3.38[−6]
	2	6.83[−5]	1.51[−5]	4.96[−5]	1.21[−5]	3.28[−5]	8.95[−6]	2.09[−5]	6.62[−6]	9.71[−6]	4.68[−6]	2.32[−6]	4.26[−6]
$1s2p\,^1P_1$	0	1.85[−3]	6.02[−5]	2.38[−3]	6.73[−5]	3.06[−3]	7.75[−5]	3.76[−3]	8.87[−5]	4.68[−3]	1.03[−4]	5.62[−3]	1.14[−4]
	1	4.46[−4]	5.07[−5]	5.84[−5]	5.84[−5]	8.06[−4]	6.84[−5]	1.35[−3]	7.99[−5]	2.36[−3]	9.72[−5]	4.66[−3]	1.19[−4]

TABLE X. Degree of polarization (in percentage) of the $1s2p$-$1s^2$ lines of Fe XXV as a function of the incident-electron energy k_j^2 considering only direct excitation (first entries) and including cascade contributions from $n=3$ higher sublevels (second entries). There is no cascade contributions for $k_j^2=495$ and 550 Ry because there is no $n=2$ cascade. Similar results but without $n=3$ cascades, derived from collision strengths of Zhang, Sampson, and Clark [19], are also displayed (third entries) for comparison with first entries.

	k_j^2 (Ry)							
Line	495	550	582	700	900	1200	2000	5000
w	58.4	60.4	61.1	60.6	58.3	47.2	33.0	9.3
	58.4	60.4	58.7	58.4	56.4	45.8	32.1	9.1
		60.3		59.0	55.5	49.5	36.5	
x	−51.8	−52.8	−53.3	−53.6	−52.4	−49.0	−37.4	−9.9
	−51.8	−52.8	−50.3	−49.9	−47.9	−43.4	−28.8	+2.9
		−52.1		−52.6	−51.6	−48.2	−37.2	
y	−19.6	−11.1	−6.8	9.2	26.6	36.3	32.0	9.3
	−19.6	−11.1	−5.2	9.3	25.5	34.6	30.6	8.6
		−12.1		7.3	23.6	33.2	32.1	

the $1s2p\ ^3P_{0,2}$ levels are seen to increase the collision strengths by more than a factor of 2. On the other hand, the z line becomes polarized by about −8% owing to the cascade transition from $1s2p\ ^3P_2$. We note that this cascade contributes to the $1s2s\ ^3S_1$ population as much as the unpolarizing cascade from $1s2p\ ^3P_0$. Indeed, the collision strength to the $1s2p\ ^3P_2$ level is five times larger than that to $1s2p\ ^3P_0$ according to the ratio of statistical weights but 1/5.5 of the $1s2p\ ^3P_2$ population goes down to $1s2s\ ^3S_1$ while the whole of the $1s2p\ ^3P_0$ population radiates to $1s2s\ ^3S_1$. It is also interesting to note that the polarization of the radiation emitted in the $1s2p\ ^3P_2-1s2s\ ^3S_1$ cascade transition is found about 28%. Table XI shows that for energies just above 582 Ry, where the $n=3$ cascading levels occurs, the cascades become the dominant mechanism for populating the $1s2s\ ^3S_1$ magnetic sublevels and cause an $\sim-14\%$ polarization on z. This polarization decreases with increasing impact energy from negative to positive values; the change of sign is at approximately 1800 Ry. It should be mentioned that at energies greater than ~900 Ry the increasingly important cascade resulting from $1s3p\ ^3P_1$ tends to induce an alignment on $1s2s\ ^3S_1$, which is of opposite sign to that induced by the cascades from $1s2p\ ^3P_2$ and $1s3p\ ^3P_2$.

It is to be expected that the amount of polarization of z due to cascades will become very insignificant when performing an energy and angular average over a realistic electron distribution met in plasma and when taking into account the contribution of the unpolarizing process of inner-shell ionization from the $1s^22s$ state which begins at $k_j^2=640$ Ry.

V. CONCLUSIONS

We have calculated the resonance and cascade contributions to the linear polarization of the Fe XXV $n=2\rightarrow n'=1$ lines. For the resonance effects, by assuming the electron energy spread just large enough to include the $1s3l3l'$ resonance group, it is found that the decrease of the polarization degree by the resonances is small for the w, x, and y lines varying from 5.6% to 6.8% in absolute value. It is clear that for plasma applications, where the energy distribution covers a large energy region including all the resonances, they have almost negligible effect on these polarizations, and we expect the higher resonances $1s3lnl'$, with $n\geq 4$, will do the same. Of course, for the line intensities these resonances play a significant role. For example, for the z line, the average effect of the $1s3l3l'$ resonances does not induce any polarization, although it can result in a very significant enhancement of the excitation rate. The present polarization results could be used in analyzing Fe XXV He-like spectra excited by a higher-energy-resolution electron beam, such as those obtained by EBIT sources because it could be possible to tune the electron energy to be in the resonance energy range (see, for example, for He-like titanium, the recent measurements of Chantrenne *et al.* [20], with ~50 eV energy resolution). A variation of the polarization versus energy ought to be seen in the region between 500 and 530 Ry (see Table IV). Now, with regard to radiative cascades, they have a weak depolarizing effect on w and y. For the x line the depolarization becomes increasingly important as the energy of the electron beam increases, reaching, for example, ~25% at four times the threshold. Below the $n=3$ excitation threshold (Table XI, $E<582$ Ry), a degree of polarization of almost −8% is found on the z line due to the cascade from the $1s2p\ ^3P_2$ level. At the $n=3$ threshold, z becomes $\sim-14\%$ polarized.

As the electron-impact energy increases, more and more excited nl levels are populated and cascade to $n'=2$

TABLE XI. Direct excitation contribution to any sublevel $M_i=0$ or 1 of Fe XXV $1s2s\ ^3S_1$ [first row, i.e., $\Omega_{dir}(1s2s\ ^3S_1)/3$]. Radiative cascade contributions from the upper $n=2$ and $n=3$ sublevels to the sublevels $M_i=0$ and 1 (second and third rows). For $k_j^2=495$ and 550 Ry the cascade contribution comes only from $n=2$ upper sublevels. In the fourth row is given the polarization degree $\eta_3(z)$, which is zero if cascades are not included. Numbers in brackets denote powers of 10.

	k_j^2 (Ry)							
	495	550	582	700	900	1200	2000	5000
$\Omega_{dir}(1s2s\ ^3S_1)/3$	1.17[−4]	1.04[−4]	9.77[−5]	7.83[−5]	5.68[−5]	3.82[−5]	1.76[−5]	3.75[−6]
$\Omega_{cas}(1s2s\ ^3S_1\ 0)$	1.68[−4]	1.43[−4]	3.39[−4]	2.52[−4]	1.64[−4]	9.96[−5]	4.81[−5]	3.86[−5]
$\Omega_{cas}(1s2s\ ^3S_1\ 1)$	1.27[−4]	1.07[−4]	2.34[−4]	1.77[−4]	1.21[−4]	8.11[−5]	4.97[−5]	4.23[−5]
$\eta_3(z)$ (%)	−7.8	−7.7	−13.7	−12.8	−10.8	−7.2	+1.2	+4.2

levels, their individual contribution can be approximately scaled as $1/n^3$, where the electron energy is normalized in excitation threshold energy units. There will have also recombination processes involving FeXXVI recombining to Fe XXV (radiative, dielectronic, and charge exchange). The angular distribution of cascade x rays following radiative recombination from an electron beam has been recently studied by Scofield [21] for H-like and He-like titanium and iron. For charge exchange, it is still an uneasy problem.

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