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Doppler redistribution of anisotropic radiation and resonance polarization in moving scattering media

I. Theory revisited in the density matrix formalism

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Abstract. Under the light of recent developments of the theory of matter-radiation interaction in the presence of magnetic field applied to non-LTE spectropolarimetry in astrophysics, we have revisited the theory of anisotropic resonance line scattering in moving media by means of the density-matrix formulation. This has led us to present a theoretical method of determination of the matter velocity field vector in the solar wind acceleration region.

The example of the O VI 103.2 nm line has been chosen for putting this theory into operation. It has been observed by the ultraviolet spectrograph SUMER of SOHO in different regions of the solar wind acceleration region; it is partially formed by resonance scattering of the incident underlying transition region radiation which competes (and can predominate) with electron collisional excitation at the low densities which prevail at these high altitudes.

The theory which is developed hereafter not only shows that this line is shifted and its intensity dimmed by the Doppler effect, due to the matter velocity field of the solar wind, but also predicts that it is linearly polarized, owing to the anisotropy of the incident radiation field; its two linear polarization parameters, degree and direction of polarization, are sensitive to the matter velocity field vector.

Our results show that the interpretation of polarimetric data, associated to the shift and the Doppler-dimming effect, may offer a method of diagnostic of the complete velocity field vector, provided that the partial anisotropy of the incident radiation field be taken into account. In fact such a diagnostic is currently missing. Yet its interest is crucial to understand various problems in astrophysics, such as stellar winds, and especially the acceleration mechanisms of the solar wind. It is also essential for a dynamical modelling of solar structures.

Key words: atomic processes – line: formation – line: profiles – polarization – Sun: corona – Sun: solar wind

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1. Introduction

Accurate knowledge of the physical conditions which prevail in the solar corona and in the solar wind acceleration region is crucial for understanding the physical mechanisms which are at the origin of coronal heating, mass and energy transport and acceleration of the solar wind. In particular, determination of both magnetic and matter velocity field vectors is essential since their coupling must be taken into account in magnetohydrodynamical modelling.

Determination of scalar quantities, such as temperatures and densities can be achieved through usual spectroscopic methods based on interpretation of the frequency dependence of the intensity of suitable lines or of the continuum of the observed radiation field. Yet determination of vectorial quantities is more complex: in fact, the complete information on strengths and directions is contained in the polarization parameters of the received radiation. Intensity transports a part of the information only. Consequently, for achieving a complete vector diagnostic, polarimetric measurements are required, and theoretical methods of spectropolarimetric diagnostics leading to all the Stokes parameters of adequate lines must be developed.

During the twenty past years methods of determination of vector magnetic fields in astrophysics have been developed (Bommier 1977; Bommier & Sahal-Bréchot 1978, and further papers), leading to interpretation of the linear polarization parameters of lines affected by the Hanle effect in terms of magnetic field vectors in solar prominences (Leroy et al. 1983, 1984; Bommier et al. 1981, 1986a, 1986b; Bommier et al. 1994; Bommier & Leroy 1997). This has been carried out with the help of the quantum theory of the matter-radiation interaction within the density matrix formalism (Fano 1957; Cohen-Tannoudji 1962, 1968, 1977; Cohen-Tannoudji et al. 1988). The entirely new vectorial data which have been obtained for solar prominences have changed our knowledge on their magnetic structure; they have shown the interest of having developed a method capable of determining the complete vector field and not only its projection on the line-of-sight as a result of conventional Zeeman studies. Likewise it would be certainly interesting to also develop a method of determination of the matter velocity field

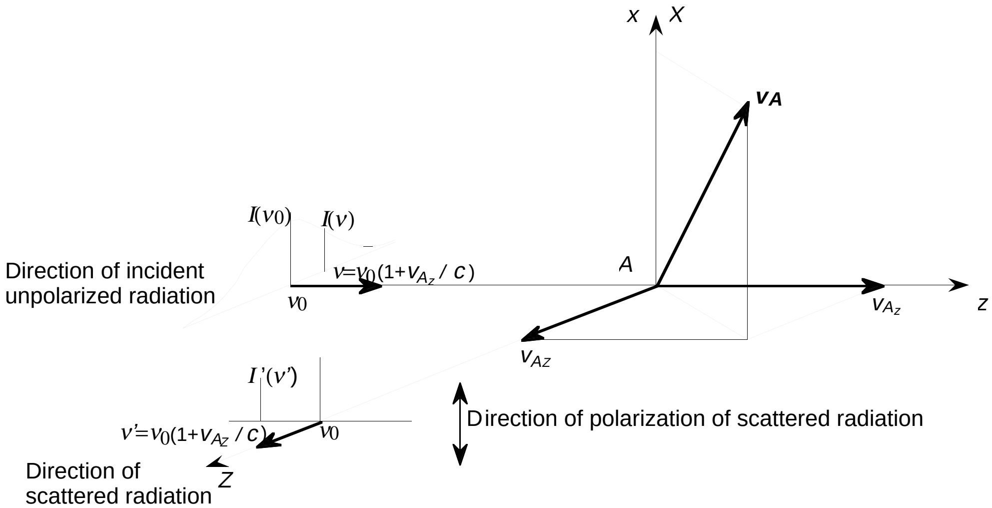

Fig. 1. Resonance scattering in the perfectly directive case: A is the scattering atom, with velocity $\boldsymbol{v}_A$. Az is the direction of incident radiation line. Its intensity $I(\nu)$is frequency-dependent. $AZ$is the line-of-sight. The atom absorbs and reemits the incident light at the eigenfrequency ν_0 in its atomic frame at rest. Due to the Doppler effect, in the laboratory frame, the atom absorbs the frequency $\nu = \nu_0(1 + v_{A_z}/c)$ and reemits the line in the AZ direction at the frequency $\nu' = \nu_0(1 + v_{A_Z}/c)$. Thus the scattered line is shifted and dimmed through $I(\nu)/I(\nu_0)$. The shift is sensitive to v_{A_Z} and the dimming to v_{A_z}. The scattered line is linearly polarized along the perpendicular to the scattering plane but the degree of polarization does not depend on the atomic velocity: for a scattering at right angles, it is equal to 1 for a normal Zeeman triplet and to 0.428 for a $J = 3/2 \rightarrow J' = 1/2$ line.

vector. Usual methods based on the interpretation of the Doppler shift of spectral lines yield only the component along the line-of-sight.

The first attempts for determining one more component of the matter velocity field of the solar acceleration region have been carried out by Gabriel (1971) and Beckers & Chipman (1974) and then by Kohl & Withbroe (1982) and Withbroe et al. (1982a, 1982b). Following Gabriel et al. (1971) who observed high in the corona during an eclipse the Lyα line of hydrogen formed by resonance scattering of the incident Lyα chromospheric radiation, they remarked that a number of other interesting lines of the transition region should also be observed high in the corona: they should be partially formed by resonance scattering of the incident transition region radiation. In particular, Li -like ion lines should be expected, and the O VI 103.2 nm line should be one of the most intense: in fact, these ions have broad abundance curves and thus a sufficient amount of ions can remain at coronal temperatures, allowing their detection. Besides, these ions may also be "frozen in" within several solar radii (Bame et al. 1974; Withbroe et al. 1982a, 1982b; Kohl & Withbroe 1982), which enhances the probability of finding them high in the corona. The EUV spectrometer SUMER of the spatial Solar and Heliospheric Observatory SOHO of ESA-NASA offers a new and important opportunity for such observations (Wilhelm et al. 1995, 1997; Lemaire et al. 1997; Hassler et al. 1997).

The interest of detecting these resonance scattered lines lies in the fact that they should be affected by the velocity field of these coronal ions: qualitatively, the moving ion absorbs the incident radiation somewhere in the incident line wing because of the Doppler effect, and the absorbed intensity is smaller. This leads to a decrease of the scattered line intensity, which is called the Doppler dimming effect (Hyder & Lites 1970). Thus the intensity of the scattered line will be sensitive to the projection of the velocity field on the direction from the scattering ion towards the region of incident radiation (the vertical to the surface of the sun in average). The scattered line being also shifted by the Doppler effect (leading to the determination of the projection of the velocity field on the line-of-sight), Kohl & Withbroe (1982) and Withbroe et al. (1982a, 1982b) suggested that the interpretation of the Doppler shift associated to that of the Doppler dimming should offer an opportunity of increasing our knowing on the velocity field of the solar wind acceleration region. They based their analysis on the basic theory of resonance scattering for an incident radiation perfectly directive, the moving scattering atoms having an anisotropic Maxwell distribution of velocities with a hydrodynamical ensemble velocity. The diagnostic can then give two informations (cf. Fig. 1): the velocity field component V_z along the direction of the incident radiation, and that along the line-of-sight V_Z. However, the component perpendicular to the scattering plane is not attained and the complete vector diagnostic cannot be achieved.

The aim of the present paper is to show that the complete information on the three components of the matter velocity field vector is contained in the three first Stokes parameters (I, Q, U: intensity and linear polarization) of the coronal scattered line sensitive to the Doppler dimming effect, provided that the complete geometry of the scattering would be taken into account (*i.e.*, the incident radiation is partially and not perfectly directive, cf. Fig. 2).

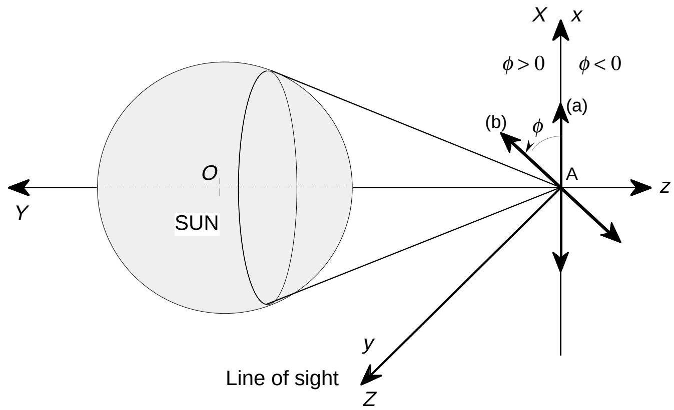

Fig. 2. The direction of polarization of the reemitted line: OAz is the preferred direction of incident radiation. (**a**) Without velocity field, the direction of polarization of the reemitted line is most often perpendicular to the scattering plane zAZ (i.e. along Ax (or AX)). It can be in the scattering plane (along AY) in the case of a strong limb brightening and a scattering atom very close to the limb. (**b**) In the presence of a velocity field, the polarization degree and polarization direction can be modified (rotation of angle ϕ). $-\pi/2 \leq \phi \leq \pi/2$

In Sect. 2 we will focus on the formalism which is general and can be applied to a variety of astrophysical problems where anisotropies of resonance scattering occur.

In Sect. 3 we will apply the formalism to the two-level atom and we will give the expression of the Stokes parameters of a coronal line formed by resonance scattering of the same line originating from the transition region and having a Doppler profile. Other line-broadening mechanisms will be neglected. Doppler redistribution of radiation will be only considered.

Quantitative results will be presented in the next paper on the example of the O VI 103.2 nm line which should be one of the most intense lines that are sensitive to the Doppler dimming effect in the solar wind acceleration region (cf. also Sahal-Bréchot et al. 1992b; Sahal-Bréchot & Choucq-Bruston 1994 where preliminary results have already been given).

2. Basic theory

2.1. General considerations

Our purpose is to calculate the Stokes parameters of a spectral line emitted by a moving ensemble of atoms which scatter the same line (resonance fluorescence) emitted by an underlying surface. As for example, we will treat the case of the O VI 103.2 nm line emitted by an ensemble of atoms located at a point A in the corona, at a height h above the solar surface. They have a macroscopic velocity $\boldsymbol{V}$ and scatter (absorb and reemit) the same line emitted by the underlying transition region of the sun (cf. Fig. 3). Owing to the anisotropy of the incident radiation which comes from the spherical cap limited by the angle $\alpha_{\max}$, the reemitted line is linearly polarized; in the absence of any perturbing effect, the direction of the linear polarization of the scattered line should be parallel to the solar limb, except at very low heights ($h < 0.07$ in units of solar radius) where it becomes radial if the limb-brightening is taken into account (Sahal-Bréchot et al. 1986).

In Sahal-Bréchot et al. (1986), the modification of the polarization of the O VI 103.2 nm line by a magnetic field (the Hanle effect) was studied. However, the Doppler effect due to the macroscopic velocity of the scattering atoms and which induces an angular dependence of the incident radiation intensity was outside the scope of that paper (cf. the end of Sect. 2.1. of Sahal-Bréchot et al. 1986). In the present one the Hanle effect is not taken into account but the Doppler effect is introduced. As we will show in the following, this angular dependence of the incident radiation can also modify the intensity and polarization parameters of the scattered line.

As far as possible, we will use the same notations as Sahal-Bréchot et al. (1992b) and Sahal-Bréchot et al. (1986) which will be referred to as Paper I hereafter.

The density-matrix formalism that we use is based on the pioneer works by Fano (1957) and Cohen-Tannoudji (1962); then it has been reinforced (Cohen-Tannoudji 1977, Omont 1977, Cohen- Tannoudji et al.1988).

This basic density-matrix formalism was extended to a complex medium for the first time for solar prominences studies by Bommier (1977) and by Bommier & Sahal-Bréchot (1978): multilevel atom coupled to an anisotropic beam of photons not perfectly directive and to a magnetic field, the strength and direction of which are unknown and may “a priori” be arbitrary. In addition, excitation by isotropic collisions has been included in Paper I.

In that formalism, rederived from the basic equations of quantum electrodynamics by Bommier & Sahal-Bréchot (1991), the coupling of an ensemble of independent identical radiating atoms with an incident anisotropic radiation field, a magnetic field and collisions (which can be also anisotropic) has been treated in a non-phenomenological manner. In that work the atoms have been assumed at rest and their levels infinitely sharp (*i.e.* the natural and collisional broadening of the lines is negligible).

The theory has been recently revisited by Bommier (1997a, 1997b), who has included line-broadening effects, through a perturbation development of the matter-radiation interaction up to orders higher than 2. Since line-broadening effects are neglected in the present paper, it can be deduced from Bommier (1997a) that effects of orders higher than 2 will vanish in the present case: consequently the lowest non-zero order (order 2) of the matter-radiation interaction will be the only one to play a role as in Bommier & Sahal-Bréchot (1991).

The problem can be reduced in a first step to the study of one atom only interacting with the ensemble (the “bath”) of perturbers. The bath is composed of photons and colliding particles which act independently: this is the impact approximation. The coupling with the particles of the bath leads to the “master equation” for the atomic density matrix, and then to the statistical equilibrium equations at the stationary state, the solution of which being the “populations” of the Zeeman states of the atom and the “coherences” between these Zeeman states (see for instance Bommier 1996 for definitions). These equations can

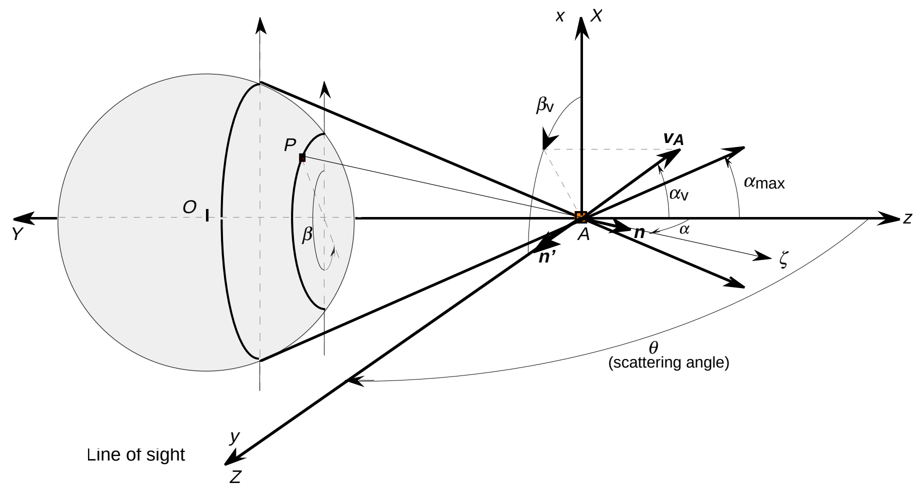

Fig. 3. Coordinates of the scattering atom A located at the distance r (in units of solar radius) from the center of the sun O; h is the height above the limb ($h = r - 1$); Oz is the preferred direction of incident radiation; AZ is the line-of-sight; Ax and AX are identical and perpendicular to the scattering plane zAZ; $\theta = (Az, AZ)$ is the scattering angle; if it is equal to $-90°$, which is the case of this figure where the atom is in the plane of the sky, then $AY = -Az$, and $AZ = Ay$. The incident light comes from the spherical cap limited by the angle $\alpha_{\max}$, each point P of the cap emits a ray $PA\zeta$ referred by the angles (α, β) in the frame $Axyz$. The atom A has the velocity $\boldsymbol{v}_A(v_A, \alpha_v, \beta_v)$; $\boldsymbol{n}$ and $\boldsymbol{n}'$ are the unitary vectors of the directions $A\zeta$ and AZ.

be considered as a generalization, convenient for polarization studies, of the usual non-LTE statistical equilibrium equations for the atomic populations.

In a second step, the polarization matrix ϕ characterizing the intensity and polarization of the emitted photons in the optically thin case is obtained through a trace of the emissivity operator over the states of the upper level of the atomic line studied (Bommier 1977; formula (37) of Bommier & Sahal-Bréchot 1978):

$$\phi = \frac{4\omega^3}{3\hbar c^3} \mathrm{Tr}_A \left[|\boldsymbol{D}\rangle \, \rho_A \, \langle \boldsymbol{D}| \right] \ , \tag{1}$$

where ω is the angular frequency of the radiation, c the velocity of light, $\boldsymbol{D}$ the atomic dipole and ρ_A the atomic density matrix.

Bommier (1997a) has shown that the above expression (1) is valid only within the second order expansion of the matter-radiation interaction.

In the third step, the Stokes parameters are obtained through a projection of that polarization matrix on a plane perpendicular to the line-of-sight and through adequate linear combinations of its matrix elements. Then the Stokes parameters of the emitted line which have thus been obtained for one atom only are summed over the total number of atoms in the element of volume (Bommier 1977 and Bommier & Sahal-Bréchot 1978 for the optically thin case).

In the final step the integration of these Stokes parameters over the line-of-sight is performed.

Bommier (1991, 1997a, 1997b) has studied the optically thick case and has rederived the radiative transfer equations in the presence of a magnetic field for the Stokes parameters through the density-matrix formalism (cf. also Landi Degl'Innocenti 1983, 1984).

In all these papers the ensemble of atoms has been assumed at rest: thus no average over the atomic velocities has been performed.

In the present paper we will extend the formalism to an assembly of moving atoms having an anisotropic distribution of velocities with a macroscopic velocity $\boldsymbol{V}$.

The present formalism will be limited to the optically thin case, which is suitable for the lines of the solar wind acceleration region. We will continue to assume that the absorption and emission profiles are infinitely sharp in the atomic rest frame. This means that the intrinsic widths of the lines are very small compared to the frequency distribution variation of the incident radiation spectrum. This is valid in the physical conditions of the solar corona, because the density is very low and thus the collisional widths are very small. Consequently the intrinsic widths are restricted to the natural ones which are also very small.

2.2. Introduction of the atomic motion in the master equation for the radiating atom

The hamiltonian of the total system (radiating atoms A in interaction with a *bath* of particles consisting of colliding particles C, and photons R, and in the presence of a magnetic field $\boldsymbol{B}$) is given by

$$H = H_A + H_{\mathrm{magn}} + H_R + H_C + V_{AR} + V_{AC} \ , \tag{2}$$

where H_{magn} is that of the interaction of the atom with the magnetic field $\boldsymbol{B}$, H_R that of the photons, and H_C that of the colliding perturbers. V_{AR} is that of the interaction atom-radiation (or atom-photons) and V_{AC} that of the collisional interaction.

We will include H_{magn} in H_A (weak field limit). The atomic wave-functions are unchanged and the internal energies become those of the Zeeman states

$$E_i + \mu_{\mathrm{B}} g_i B \ .$$

The eigenstates of the atom at rest are denoted as

$$|\alpha_i J_i M_i\rangle \ .$$

The total system being isolated and reversible, its evolution can be described by a density matrix ρ, solution of the Schrödinger (or Liouville equation for the density matrix) equation

$$\mathrm{i}\hbar \frac{\mathrm{d}\rho}{\mathrm{d}t} = \Big[H, \rho \Big] \ . \tag{3}$$

The evolution of the subsystem A is described by the density matrix ρ_A, defined by a trace over all the states of the bath

$$\rho_A = Tr_{R,C}\left[\rho\right] \ . \tag{4}$$

Within the usual language of classical mechanics, this trace corresponds to an average over all the particles of the bath.

As shown by Cohen-Tannoudji (1977), (cf. also Cohen-Tannoudji et al. 1988), the subsystem A can be described by the reduced density matrix $\rho_A(t)$. Yet we have now to derive the time evolution of $\rho_A(t)$, which is not given by a Liouville equation, because we have introduced some irreversibility in the above process: A is no longer an isolated system. The time evolution of $\rho_A(t)$ is now given by a "master equation" and describes the evolution of a "mean" atom. In the corresponding classical picture, the time evolution of the classical equivalent to $\rho_A(t)$ is given by the "kinetic equation" (Oxenius 1986). It can be shown that the kinetic equation for the radiating atoms or molecules reduces to the so-called "statistical equilibrium equations" at the steady state.

We have to solve the master equation in a fixed reference frame, the so-called laboratory frame. In this frame, the atomic velocity operator is denoted by $\boldsymbol{v}_A$ and the atomic mass by m_A. Following Smith et al. (1971) and Nienhuis (1976, 1977), we introduce the atomic motion in the atomic hamiltonian H_A by adding the translation hamiltonian H_{tr} to the atomic hamiltonian at rest H_0.

The atomic hamiltonian in the laboratory frame is

$$H_A = H_0 + H_{\mathrm{tr}} \ , \tag{5}$$

with

$$H_{\mathrm{tr}} = \frac{\boldsymbol{p}_A^2}{2m_A} \ , \tag{6}$$

where $\boldsymbol{p}_A$ is the atomic momentum operator. $\boldsymbol{p}_A = m_A \boldsymbol{v}_A$ in the non relativistic limit. The corresponding eigenstate of H_{tr} is the plane wave $|p_A\rangle$, with eigenvalue

$$\frac{1}{2} m_A v_A^2 \ .$$

The eigenstates of H_A take the form

$$|u_i\rangle = |\alpha_i J_i M_i\rangle \otimes |p_A\rangle = |\alpha_i J_i M_i, \, p_A\rangle \ , \tag{7}$$

and the corresponding eigenvalues are

$$E_i + \frac{1}{2} m_A v_A^2 \ . \tag{8}$$

We assume that collisional and radiation interactions do not modify the state of translation of the atom: this implies that the recoil of the atom during an interaction and the radiation pressure effects can be neglected. At the considered temperatures, the atomic momentum is much larger than the atomic electron one $\boldsymbol{p}_{\mathrm{e}} = m_{\mathrm{e}} \boldsymbol{v}_{\mathrm{e}}$ because of the high m_A/m_{e} ratio. It is also much larger than the photon momentum $\hbar k = \hbar\omega/c$, ω being the angular frequency and c the velocity of light. Consequently, the momentum transfer during the interaction process is negligible and the atomic velocity is unchanged (Cohen-Tannoudji et al. 1988, p. 286). Apart from the recoil effect, the atomic velocity can change under the effect of collisions (Oxenius 1986, Sahal-Bréchot et al. 1992a, Landi Degl'Innocenti 1996): this change can be determined by calculating the differential cross-section (Balança & Feautrier 1998, Balança et al. 1998). In fact, owing to the relative mass effect, only elastic collisions with particles having about same masses or higher masses than the radiating atom are to be considered. In solar and coronal conditions, these collisional rates are very weak and the atom has time to emit a photon or to be excited or deexcited by any process before its velocity could change (cf. Sahal-Bréchot et al. 1992a).

Consequently the atomic velocity is conserved, and the matrix elements of H can be considered as diagonal in $|p_A\rangle$.

Therefore, as for the external states $|p_A\rangle$, the off-diagonal terms $\langle p_A | \rho_A(t) | p'_A\rangle$ (with $p_A \neq p'_A$) will be completely decoupled from the diagonal terms $\langle p_A | \rho_A(t) | p_A\rangle$ which represent the density of probability (with respect to the external states $|p_A\rangle$) for the density matrix (with respect to the internal states $|\alpha_i J_i M_i\rangle$) of an atom having the velocity $\boldsymbol{v}_A$. We can denote this density matrix as $\rho_A(\boldsymbol{v}_A, t)$, the elements of which are given by

$$\begin{aligned} &\langle \alpha_i J_i M_i, p_A | \, \rho_A(t) \, |\alpha_{i'} J_{i'} M_{i'}, p_A\rangle \\ &= \langle \alpha_i J_i M_i | \, \rho_A(\boldsymbol{v}_A, t) \, |\alpha_{i'} J_{i'} M_{i'}\rangle \ . \end{aligned} \tag{9}$$

Moreover, there will be no coupling between the reduced density matrices $\rho_A(\boldsymbol{v}_A, t)$ and $\rho_A(\boldsymbol{v}'_A, t)$ for different velocities $\boldsymbol{v}_A$ and $\boldsymbol{v}'_A$.

For the internal states $|\alpha_i J_i M_i\rangle$, $\rho_A(\boldsymbol{v}_A, t)$ will be normalized to the distribution of atomic velocities $F(\boldsymbol{v}_A)$:

$$\mathrm{Tr}_A\left[\rho_A(\boldsymbol{v}_A, t)\right] = F(\boldsymbol{v}_A) \ , \tag{10}$$

with the normalization condition

$$\int F(\boldsymbol{v}_A) \, \mathrm{d}^3\boldsymbol{v}_A = 1 \ . \tag{11}$$

Then, within the *impact approximation*, the interactions with the radiation (the photons) R and the colliding particles C (which are the electrons of the corona in the present paper) are decoupled (cf. Bommier & Sahal-Bréchot 1991 and earlier papers). Using the no-back reaction approximation and the impact approximation, as well as the fact that the evolution of the reduced atomic density matrix $\rho_A(\boldsymbol{v}_A, t)$ is markovian, *i.e.*, it depends only on the instant t and not on the past history of ρ_A, one obtains the master equation for $\rho_A(\boldsymbol{v}_A, t)$ (cf. Bommier & Sahal-Bréchot 1991 and Sahal-Bréchot et al. 1992a for the atom at rest), which describes the evolution of the atom coupled to the bath in the laboratory frame (cf. also Cohen-Tannoudji 1977 and related papers, Omont 1977, Bommier 1997a). The *coarse grained approximation* leads to achieve the average over the states of the bath over a time interval s much larger than the mean duration τ_c of a collision (resp. τ_r for the radiation) and much smaller than the mean interval Γ_c^{-1} between two successive collisions (resp. Γ_r^{-1}, inverse of the lifetime). We obtain

$$\frac{\mathrm{d}\rho_A(\boldsymbol{v}_A, t)}{\mathrm{d}t} = \frac{1}{\mathrm{i}\hbar}\left[H_A, \rho_A(\boldsymbol{v}_A, t)\right] + \left(\frac{\mathrm{d}\rho_A(\boldsymbol{v}_A, t)}{\mathrm{d}t}\right)_{\mathrm{B}} ,$$

Owing to the fact that the interactions of the particles of the bath B are decoupled, we can write

$$\begin{aligned}
&\frac{\mathrm{d}\rho_A(\boldsymbol{v}_A,t)}{\mathrm{d}t} = \frac{1}{\mathrm{i}\hbar}\left[H_A,\rho_A(\boldsymbol{v}_A,t)\right] \\
&\quad+\left(\frac{\mathrm{d}\rho_A(\boldsymbol{v}_A,t)}{\mathrm{d}t}\right)_{\mathrm{R}}+\left(\frac{\mathrm{d}\rho_A(\boldsymbol{v}_A,t)}{\mathrm{d}t}\right)_{\mathrm{C}} \\
&=\frac{1}{\mathrm{i}\hbar}\left[H_A,\rho_A(\boldsymbol{v}_A,t)\right] \\
&\quad+\frac{1}{s}\mathrm{Tr}_R\left[S_R\rho_A(\boldsymbol{v}_A,t)\,\rho_R\,S_R^{-1}-\rho_A(\boldsymbol{v}_A,t)\,\rho_R\right] \\
&\quad+\frac{1}{s}\mathrm{Tr}_C\left[S_C\rho_A(\boldsymbol{v}_A,t)\,\rho_C\,S_C^{-1}-\rho_A(\boldsymbol{v}_A,t)\,\rho_C\right] \\
&=0 \quad \text{at the steady state .}
\end{aligned} \tag{12}$$

S is the scattering matrix calculated for a complete interaction with one colliding perturber (subscript C) or with one photon (subscript R). The density matrix of the photons ρ_R and the density matrix ρ_C of the colliding perturbers take place for averaging over all the particles, i.e. over the states of the bath B (radiation R and collisions C that are decoupled).

Then we obtain the projected master equation averaged over the states of the bath (cf. also Cohen-Tannoudji et al. 1988, p. 486):

$$\begin{aligned}
&\frac{\mathrm{d}}{\mathrm{d}t}\langle i|\,\rho_A(\boldsymbol{v}_A,t)\,|i'\rangle \\
&=-\mathrm{i}\omega_{ii'}\langle i|\,\rho_A(\boldsymbol{v}_A,t)\,|i'\rangle \\
&\quad+\sum_{jj'}\langle j|\,\rho_A(\boldsymbol{v}_A,t)\,|j'\rangle\left\langle S_{j'i'}^{-1}\,S_{ij}-\delta_{ij}\delta_{i'j'}\right\rangle_{\mathrm{B}} .
\end{aligned} \tag{13}$$

where $\omega_{ii'}$ is the Bohr frequency and δ_{ij} and $\delta_{i'j'}$ are Kronecker symbols.

2.3. Coupling with the colliding particles

We show now that the collisional interaction is not affected by the atomic motion for the processes studied in the present paper: electrons are very much faster than the radiating ions which can thus be assumed at rest. This would not have been the case if we were studying collisions between heavy particles such as protons and hydrogen atoms for instance (cf. Sahal-Bréchot et al. 1996). Herewith the collisional S matrix has to be calculated within the usual framework of the theory of collisions. The colliding perturbers are isotropic maxwellian electrons at a temperature T_{e}. Owing to that isotropy, the collisional coupling coefficients in the master equation link only the Zeeman populations $\langle i|\,\rho_A(t)\,|i\rangle$, $\langle j|\,\rho_A(t)\,|j\rangle$. In addition the coronal electronic densities are so low that depolarizing collisions are negligible: the photon is emitted before any occurrence of a collision. The only collisional processes that are to be taken into account are excitation from the lower level towards the upper level of the considered transition. Deexcitation by collisions is negligible. The velocity that takes place in the calculation of the cross-section $\sigma(i\to j,\boldsymbol{v}_{\mathrm{e}})$ is the relative one $\boldsymbol{v}_{\mathrm{e}}-\boldsymbol{v}_A$ which can be considered as equal to the electron velocity $\boldsymbol{v}_{\mathrm{e}}$.

Thus the problem is rather straightforward. We obtain after integration over the Maxwell isotropic distribution of electron velocities $f(v_{\mathrm{e}})$ at the temperature T_{e} and multiplication by the electron density N_{e}

$$\begin{aligned}
&\left\langle\left[S_C^{-1}\right]_{j'i'}\left[S_C\right]_{ij}\right\rangle_{\mathrm{C}} \\
&=\delta_{jj'}\delta_{ii'}N_{\mathrm{e}}\int v_{\mathrm{e}}\,f(v_{\mathrm{e}})\sigma(i\to j,v_{\mathrm{e}})\,\mathrm{d}v_{\mathrm{e}} \\
&=\delta_{jj'}\delta_{ii'}N_{\mathrm{e}}\,\alpha(i\to j,T_{\mathrm{e}}) .
\end{aligned} \tag{14}$$

In order to conclude this paragraph, the collisional cross-sections and their average over the electron distribution of velocities are not affected by the atomic motion. The effect of the atomic motion only appears in the matrix elements of the atom-radiation field interaction that we will study in the following.

2.4. Coupling with the radiation field

Ignoring collisions now, the total hamiltonian H of the system is that of the subsystem atom+radiation and can be written as

$$H=H_A+H_R+V_{AR} , \tag{15}$$

where H_R is the hamiltonian of the radiation field and V_{AR} is the atom-radiation interaction

$$V_{AR}=-\boldsymbol{D}\cdot\boldsymbol{E}(\boldsymbol{r}_A) , \tag{16}$$

where $\boldsymbol{D}$ is the atomic dipole which can be calculated in the basis of the atom at rest (Smith et al. 1971), $\boldsymbol{E}(\boldsymbol{r}_A)$ is the electric field of the radiation and $\boldsymbol{r}_A$ is the atomic position operator in the laboratory frame.

Using the second quantization formalism, the electric field can be developed in plane waves as

$$\begin{aligned}
&\boldsymbol{E}(\boldsymbol{r}_A)=\int \mathrm{d}^3\boldsymbol{k}\sum_{\lambda}\mathrm{i}\sqrt{\frac{h\omega}{(2\pi)^3}} \\
&\times\left[a(\boldsymbol{k},\lambda)\,\boldsymbol{e}_{\lambda}\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r}_A}-a^{\dagger}(\boldsymbol{k},\lambda)\,\boldsymbol{e}_{\lambda}^{*}\mathrm{e}^{-\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r}_A}\right] ,
\end{aligned} \tag{17}$$

for the mode $\boldsymbol{k}(k,\hat{k})$, $\hat{k}$ defining the direction of unitary vector $\boldsymbol{n}$, the angular frequency $\omega=ck$ and the polarization λ (polarization vector $\boldsymbol{e}_{\lambda}$). $a(\boldsymbol{k},\lambda)$ and $a^{\dagger}(\boldsymbol{k},\lambda)$ are the operators of creation and annihilation of photons for the radiation propagating in the $\boldsymbol{n}$ direction with the polarization λ. The normalisation chosen is such that

$$\left\langle a^{\dagger}(\boldsymbol{k},\lambda)\,a(\boldsymbol{k},\lambda)\right\rangle=\frac{c^2}{h\nu^3}I_{\nu}(\nu,\hat{k},\lambda) , \tag{18}$$

where $I_{\nu}(\nu,\hat{k},\lambda)$ is the specific intensity.

By expanding the dipole moment $\boldsymbol{D}$ and the polarization vector $\boldsymbol{e}_{\lambda}$ over the standard basis $\boldsymbol{e}_{+}$, $\boldsymbol{e}_{-}$, $\boldsymbol{e}_{0}$ of the laboratory frame (see Bommier & Sahal-Bréchot 1991 or Bommier 1997a for definition), we obtain

$$V_{AR} = -\mathrm{i} \int \mathrm{d}^3\boldsymbol{k} \sum_{\lambda} \sqrt{\frac{h\omega}{(2\pi)^3}} \sum_{p=-1}^{1} \Big[a(\boldsymbol{k},\lambda)(-1)^p D_p [e_\lambda]_{-p} \mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r}_A} - a^\dagger(\boldsymbol{k},\lambda) D_{-p} [e_\lambda]^*_{-p} \mathrm{e}^{-\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r}_A} \Big] . \tag{19}$$

The atom-radiation interaction contains the atomic motion through the factors $\mathrm{e}^{\pm\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r}_A}$.

The developments made by Bommier & Sahal-Bréchot (1978, 1991) for the atom at rest can now be rewritten for a moving atom by including these factors.

By using, in interaction representation,

$$\begin{aligned} \widetilde{\boldsymbol{r}}_A(\tau) &= \mathrm{e}^{\mathrm{i}H_A\tau/\hbar}\, \boldsymbol{r}_A\, \mathrm{e}^{-\mathrm{i}H_A\tau/\hbar} \\ &= \mathrm{e}^{\mathrm{i}\boldsymbol{p}_A^2\tau/2m_A\hbar}\, \boldsymbol{r}_A\, \mathrm{e}^{-\mathrm{i}\boldsymbol{p}_A^2\tau/2m_A\hbar} \\ &= \boldsymbol{r}_A + \frac{\boldsymbol{p}_A}{m_A}\tau \\ &= \boldsymbol{r}_A + \boldsymbol{v}_A\tau , \end{aligned} \tag{20}$$

it can be shown that the results of Bommier & Sahal-Bréchot (1991), obtained for an atom at rest at $\boldsymbol{r}_A = 0$, can be applied to the present case by replacing any radiation angular frequency ω which appears in the $\mathrm{e}^{\pm\mathrm{i}\omega\tau}$ terms of that calculation by $(\omega - \boldsymbol{k}\cdot\boldsymbol{v}_A)$, $\boldsymbol{k}$ being the corresponding wave vector.

Due to the Markov approximation which leads to limit the perturbation expansion of the matter-radiation interaction to the lowest non-zero order (the second one), the absorption and emission line profiles are infinitely sharp: this is valid if the intrinsic line-widths are very small compared to the frequency distribution variation of the incident radiation spectrum. All this is valid in the solar corona.

Thus the scattering is monochromatic in the atomic rest frame. Following Hummer (1962) and Mihalas (1978, p. 412), this corresponds to Case I of the redistribution process (Doppler redistribution).

Therefore, as the line profile is now given by a δ-function

$$\delta\left[\omega_0 - (\omega - \boldsymbol{k}\cdot\boldsymbol{v}_A)\right] , \tag{21}$$

the only change is to replace the angular atomic frequency ω_0 (frequency corresponding to the energy difference between the levels u (upper) and l (lower) which are assumed infinitely sharp) by

$$\omega_0 + (\boldsymbol{k}\cdot\boldsymbol{v}_A) = \omega_0 \left[1 + \frac{(\boldsymbol{v}_A\cdot\boldsymbol{n})}{c}\right] \tag{22}$$

in the laboratory frame. This is expected: the atom absorbs the radiation at its eigenfrequency ω_0 in its own 'comoving' frame (the atomic frame), and due to the Doppler effect, this corresponds to the angular frequency

$$\omega_0 \left[1 + \frac{(\boldsymbol{v}_A\cdot\boldsymbol{n})}{c}\right]$$

of the incident radiation in the laboratory frame.

We are now able to write the system of statistical equilibrium equations leading to the populations and coherences for the atom having the velocity $\boldsymbol{v}_A$, interacting with an incident flux of radiation characterized by its polarization matrix φ. We will solve this system in the atomic frame.

The first step is to obtain the elements of the polarization matrix which enter the coupling coefficients of the elements of the master equation leading to the atomic density matrix $\rho_A(\boldsymbol{v}_A)$.

2.4.1. Calculation of the components of the polarization matrix φ of the incident radiation

We will use the same notations as those of previous papers (Paper I and Sahal-Bréchot et al. 1992b).

We will assume hereafter that the incident line profile is the same for all the points of the solar surface: we will neglect inhomogeneities and take only into account the limb-brightening effect. Inclusion of inhomogeneities of the incident radiation was studied in Paper I and is outside the scope of the present paper.

The atomic frame (the comoving frame) $Axyz$ and the line-of-sight frame (the laboratory frame) $AXYZ$ are defined in the following manner (cf. Fig. 3):

- AZ is the line-of-sight oriented towards the observer, its unitary vector is denoted by $\boldsymbol{n}'$.
- Az is the axis of symmetry of the system atom A + incident photons, that is the vertical to the surface of the Sun. It is oriented towards the outside.
- The scattering plane contains the vertical Az and the line-of-sight AZ.
- The perpendicular to the scattering plane defines the common axis Ax and AX. Since there are "a priori" two possible senses for Ax (or AX), the angle θ between Az and AZ is oriented. The sense of Ax (or AX) gives the sign of θ which is thus not a polar angle.
- Ay and AY follow.

The polar angle α_v and the azimuth β_v define the direction of the atomic velocity $\boldsymbol{v}_A$ in the atomic frame. Likewise the polar angle α_V and the azimuth β_V define the direction of the velocity $\boldsymbol{V}$ of the ensemble of atoms in the atomic frame.

The polarization matrix φ of the incident photons which was obtained for the atomic frequency only in Paper I is now frequency-dependent through the Doppler effect.

The radiation incoming from an element of volume centered on a point P of the solar surface in the direction $PA\zeta$ (unitary vector $\boldsymbol{n}$) is characterized by the angles α and β in the atomic $Axyz$ frame. $J(\alpha, \beta, \nu)$ is the angular distribution of the local intensity at point A for the frequency ν (erg cm^{-2} s^{-1} sr^{-1} Hz^{-1}) and incoming from point P:

$$\nu = \nu_0 \left[1 + \frac{(\boldsymbol{v}_A\cdot\boldsymbol{n})}{c}\right] \tag{23}$$

and (cf. Fig. 3)

$$(\boldsymbol{v}_A \cdot \boldsymbol{n}) = v_A \left[\cos\alpha\cos\alpha_v + \sin\alpha\sin\alpha_v\left(\cos\beta\cos\beta_v + \sin\beta\sin\beta_v\right)\right] . \quad (24)$$

The polarization matrix is obtained through an angular average over all the directions of incoming radiation as follows:

$$\varphi\left(\nu\right) = \frac{1}{4\pi}\int_0^{\alpha_{\max}} \sin\alpha\,\mathrm{d}\alpha \int_0^{2\pi}\mathrm{d}\beta \times \mathcal{M}\left(\alpha,\beta\right) J\left(\alpha,\beta,\nu\right) \ . \quad (25)$$

The matrix $\mathcal{M}\left(\alpha,\beta\right)$ characterizes the angular behaviour of a dipolar unpolarized radiation incoming in the direction $P\zeta$ and is normalized to unity (Eq. (43) of Paper I). It can be expanded in multipoles terms over the basis of the irreducible tensorial operators T_q^k relative to the quantization axis Az which is also adapted to the present problem:

$$\mathcal{M}\left(\alpha,\beta\right) = \sum_{k=0}^{2}\sum_{q=-k}^{k} \mathcal{M}_q^k\left(\alpha,\beta\right) T_q^k \ . \quad (26)$$

The results are the same as in Paper I, except that β is changed into $\beta+\pi$. In fact, in Paper I, β was the azimuth of point P, referred in cylindrical coordinates in the atomic frame $Axyz$. In the present paper, β is the azimuth of the $PA\zeta$ direction in the atomic frame.

$$\begin{aligned}
\mathcal{M}_0^0\left(\alpha,\beta\right) &= \frac{1}{\sqrt{3}}\,, \\
\mathcal{M}_0^1\left(\alpha,\beta\right) &= \mathcal{M}_{\pm1}^1\left(\alpha,\beta\right) = 0\,, \\
\mathcal{M}_0^2\left(\alpha,\beta\right) &= \frac{1}{2\sqrt{6}}\left(3\cos^2\alpha - 1\right) \ , \quad (27)
\end{aligned}$$

$$\begin{aligned}
\mathcal{M}_1^2\left(\alpha,\beta\right) &= -\left[\mathcal{M}_{-1}^2\left(\alpha,\beta\right)\right]^* \\
&= \frac{1}{4}\sin 2\alpha\,\left(-\cos\beta + \mathrm{i}\sin\beta\right) \ ,
\end{aligned}$$

$$\begin{aligned}
\mathcal{M}_2^2\left(\alpha,\beta\right) &= \left[\mathcal{M}_{-2}^2\left(\alpha,\beta\right)\right]^* \\
&= \frac{1}{4}\sin^2\alpha\,\left(\cos 2\beta - \mathrm{i}\sin 2\beta\right) \ .
\end{aligned}$$

Thus, using relations (23) and (24), the polarization matrix of the incident radiation becomes velocity-dependent and we can write

$$\varphi_q^k\left(\boldsymbol{v}_A\right) = \frac{1}{4\pi}\int_0^{\alpha_{\max}} \sin\alpha\,\mathrm{d}\alpha \int_0^{2\pi}\mathrm{d}\beta \times \mathcal{M}_q^k\left(\alpha,\beta\right) J\left(\alpha,\beta,\boldsymbol{v}_A\right) \ . \quad (28)$$

The angular dependence of $(\boldsymbol{v}_A \cdot \boldsymbol{n})$ on the angles α and β leads to appearance of coherences ($k = 2$, $q = \pm1$, ±2) in the matrix elements of $\varphi\left(\boldsymbol{v}_A\right)$ (in the atomic $Axyz$ frame). This will lead to a rotation of polarization of the scattered line because the cylindrical symmetry about the Az axis is broken. We can also verify that the coherences vanish when the atomic velocity is directed along the vertical (*i.e.*, $\alpha_v = 0$), because the cylindrical symmetry is recovered.

Formula (28) is general and does not assume anything on the incident radiation. Thus we have obtained a general expression for the elements of the polarization matrix of the incident radiation that enter the statistical equilibrium equations that have to be solved for each atomic velocity $\boldsymbol{v}_A$. In fact, it is also not restricted to the two-level atom.

2.4.2. Expression of the statistical equilibrium equations for the atomic density-matrix and for a two-level atom

We refer to Eq. (5) of Paper I in zero-magnetic field, and we use Eq. (28) of the present paper for the polarization matrix of the incident radiation. We obtain for a two-level atom (u = upper and l = lower level) in the atomic reference frame in the basis of the irreducible tensorial operators

$$\begin{aligned}
&A\left(J_u \to J_l\right)\ {}^{J_uJ_u}\left[\rho_A\right]_q^k\left(\boldsymbol{v}_A\right) = {}^{J_lJ_l}\left[\rho_A\right]_0^0\left(\boldsymbol{v}_A\right) \\
&\times\Bigg[N_{\mathrm{e}}\,\alpha\left(J_l \to J_u, T_{\mathrm{e}}\right)\sqrt{\frac{2J_l+1}{2J_u+1}}\,\delta\left(k,0\right) \\
&\quad + B\left(J_l \to J_u\right)\varphi_q^k\left(\boldsymbol{v}_A\right) \\
&\quad \times 3\sqrt{2J_l+1}(-1)^{J_u+J_l+k+1}\begin{Bmatrix} 1 & 1 & k \\ J_u & J_u & J_l \end{Bmatrix}\Bigg] \ . \quad (29)
\end{aligned}$$

In the above equation k can be equal to 0 (population component) or 2 (alignment component). $A\left(J_u \to J_l\right)$ and $B\left(J_l \to J_u\right)$ are the Einstein cofficients for spontaneous emission and absorption, $N_{\mathrm{e}}\,\alpha\left(J_l \to J_u, T_{\mathrm{e}}\right)$ the collisional excitation coefficient (obtained from Eq. (14)). $\delta\left(k,0\right)$ is the Kronecker symbol: in fact the collisional processes are isotropic and cannot create alignment and we recall that they are also insensitive to the atomic velocity.

The symbol between brackets is a "6j" coefficient. Stimulated emission is negligible and is not introduced. Likewise collisional deexcitation is negligible. The only relevant processes are excitation by collisions and radiation followed by spontaneous emission.

2.4.3. Expression of the polarization matrix of the reemitted photons

Following Eq. (49) of Paper I and using the present equation (28), the polarization matrix $\phi\left(\boldsymbol{v}_A\right)$ of the photons reemitted by an atom having the velocity $\boldsymbol{v}_A$ can now be written. $\phi\left(\boldsymbol{v}_A\right)$ being expressed in number of emitted photons per second, per unit of volume and per unit of distribution of atomic velocities $F\left(\boldsymbol{v}_A\right)$, we obtain in the $Axyz$ atomic comoving reference frame

$$\begin{aligned}
\phi_0^0\left(\boldsymbol{v}_A\right) &= N_A\,n_l\left(\boldsymbol{v}_A\right)\frac{1}{\sqrt{3}}N_{\mathrm{e}}\,\alpha\left(J_l \to J_u\right) \\
&\quad + N_A\,n_l\left(\boldsymbol{v}_A\right)B(J_l \to J_u)\varphi_0^0\left(\boldsymbol{v}_A\right) \ , \\
\phi_0^2\left(\boldsymbol{v}_A\right) &= N_A\,n_l\left(\boldsymbol{v}_A\right)B(J_l \to J_u)c_{ul}\varphi_0^2\left(\boldsymbol{v}_A\right) \ , \quad (30)
\end{aligned}$$

$$\begin{aligned}
\phi_1^2\left(\boldsymbol{v}_A\right) &= -\left[\phi_{-1}^2\right]^*\left(\boldsymbol{v}_A\right) \\
&= N_A\,n_l\left(\boldsymbol{v}_A\right)B(J_l \to J_u)c_{ul}\varphi_1^2\left(\boldsymbol{v}_A\right) \ ,
\end{aligned}$$

$$\begin{aligned}\phi_2^2(\boldsymbol{v}_A) &= \left[\phi_{-2}^2\right]^*(\boldsymbol{v}_A) \\ &= N_A\, n_l(\boldsymbol{v}_A)\, B(J_l \to J_u) c_{ul} \varphi_2^2(\boldsymbol{v}_A)\ ;\end{aligned}$$

where N_A is the number of radiating ions per unit of volume.

$n_l(\boldsymbol{v}_A)$ is the population of the lower level l (the ground state), $n_u(\boldsymbol{v}_A)$ is that of the excited state u:

$$n_l(\boldsymbol{v}_A) = \sqrt{2J_l+1}\ {}^{J_lJ_l}[\rho_A]_0^0(\boldsymbol{v}_A)\ , \tag{31}$$

$$n_u(\boldsymbol{v}_A) = \sqrt{2J_u+1}\ {}^{J_uJ_u}[\rho_A]_0^0(\boldsymbol{v}_A)\ . \tag{32}$$

The normalization condition for the internal states (Eq. (10)) leads to

$$n_l(\boldsymbol{v}_A) + n_u(\boldsymbol{v}_A) = F(\boldsymbol{v}_A) \tag{33}$$

for the two-level atom. Since the population of the excited level is very small compared to that of the ground state, (case of the solar corona), the preceding Eq. (33) reduces to

$$n_l(\boldsymbol{v}_A) = F(\boldsymbol{v}_A)\ , \tag{34}$$

$F(\boldsymbol{v}_A)$ being the distribution of atomic velocities.

c_{ul} is defined in Paper I, according to Bommier & Sahal-Bréchot (1982). It can be expressed in terms of "6j" coefficients as

$$c_{ul} = 3\ (2J_u+1)\begin{Bmatrix}1 & 1 & 2\\ J_u & J_u & J_l\end{Bmatrix}^2 = \frac{\begin{Bmatrix}1 & 1 & 2\\ J_u & J_u & J_l\end{Bmatrix}^2}{\begin{Bmatrix}1 & 1 & 0\\ J_u & J_u & J_l\end{Bmatrix}^2}. \tag{35}$$

It can be noticed that c_{ul} is equal to the coefficient

$$W_2(J,J') = \left[w_{J'J}^{(2)}\right]^2$$

defined in Landi Degl'Innocenti (1984) where these coefficients are tabulated. The notation of Stenflo (1994, pp. 184-187 and p. 91) for c_{ul} is W_2. We have (Paper I):

$$\begin{cases} c_{ul} = \frac{1}{2} \ \text{for}\ \ J_u = \frac{3}{2} \text{ and } J_l = \frac{1}{2}\\ c_{ul} = 1 \ \text{for}\ \ J_u = 1 \text{ and } J_l = 0\\ c_{ul} = 0 \ \text{for}\ \ J_u = \frac{1}{2} \text{ and } J_l = \frac{1}{2}\\ c_{ul} = 0 \ \text{for}\ \ J_u = 0 \text{ and } J_l = 1\end{cases},$$

In order to obtain the polarization matrix of the reemitted photons in the line-of-sight fixed frame $AXYZ$ (the laboratory frame), we have to rotate the coordinates system and to use the proper atomic eigenfuctions including translation states.

Firstly, we have to perform a rotation through the angle θ about the Ax-axis (or AX). In fact we will modify Eq. (50) of Paper I which was obtained for $\theta = -\pi/2$. The resulting polarization matrix is denoted by $\Phi(\boldsymbol{v}_A)$.

The Euler angles of the rotation are

$$\left(-\frac{\pi}{2}, \theta, +\frac{\pi}{2}\right)\ .$$

We obtain after calculations

$$\Phi_0^0(\boldsymbol{v}_A) = \phi_0^0(\boldsymbol{v}_A)\ ,$$

$$\begin{aligned}\Phi_0^2(\boldsymbol{v}_A) &= \tfrac{1}{2}\left(3\cos^2\theta - 1\right)\phi_0^2(\boldsymbol{v}_A)\\ &-\sqrt{\frac{3}{2}}\sin 2\theta\,\mathrm{Im}\phi_1^2(\boldsymbol{v}_A) - \sqrt{\frac{3}{2}}\sin^2\theta\,\mathrm{Re}\phi_2^2(\boldsymbol{v}_A)\ ,\end{aligned} \tag{36}$$

$$\begin{aligned}\Phi_2^2(\boldsymbol{v}_A) &= -\sqrt{\frac{3}{8}}\sin^2\theta\,\phi_0^2(\boldsymbol{v}_A)\\ &-\frac{1}{2}\sin 2\theta\,\mathrm{Im}\phi_1^2(\boldsymbol{v}_A) + \mathrm{i}\sin\theta\,\mathrm{Re}\phi_1^2(\boldsymbol{v}_A)\\ &+\frac{1}{2}\left(1+\cos^2\theta\right)\mathrm{Re}\phi_2^2(\boldsymbol{v}_A) + \mathrm{i}\cos\theta\,\mathrm{Im}\phi_2^2(\boldsymbol{v}_A)\ .\end{aligned}$$

The other components are not written because they do not enter the expression of the following Stokes parameters.

Secondly, the atom reemits the line at the frequency ν_0 in the comoving atomic frame. We introduce the translation motion dependence in the atomic wavefunctions of the laboratory frame.

This is equivalent to write that, due to the Doppler effect, the observer sees a line emitted at the shifted frequency

$$\nu' = \nu_0\left[1 + \frac{(\boldsymbol{v}_A\cdot\boldsymbol{n}')}{c}\right]\ ,$$

along the line-of-sight AZ characterized by the direction $\boldsymbol{n}'$. We have

$$\begin{aligned}(\boldsymbol{v}_A\cdot\boldsymbol{n}') &= v_{A_Z}\\ &= v_A\left(\cos\theta\cos\alpha_v - \sin\theta\sin\alpha_v\sin\beta_v\right)\ ,\end{aligned} \tag{37}$$

and

$$\nu' - \nu_0 = \nu_0\frac{v_{A_Z}}{c}\ . \tag{38}$$

Then, the Stokes parameters of the line emitted at the frequency ν' in the AZ direction by the atoms having the given projected velocity v_{A_Z} can be obtained.

Firstly, for the atoms having the velocity $\boldsymbol{v}_A$ $(v_{A_X}, v_{A_Y}, v_{A_Z})$, we have (Paper I)

$$\begin{aligned}I(\boldsymbol{v}_A) &= \frac{3}{8\pi}\left[\frac{2}{\sqrt{3}}\Phi_0^0(\boldsymbol{v}_A) + \frac{2}{\sqrt{6}}\Phi_0^2(\boldsymbol{v}_A)\right]\ ,\\ Q(\boldsymbol{v}_A) &= -\frac{3}{8\pi}2\mathrm{Re}\Phi_2^2(\boldsymbol{v}_A)\ ,\\ U(\boldsymbol{v}_A) &= \frac{3}{8\pi}2\mathrm{Im}\Phi_2^2(\boldsymbol{v}_A)\ ,\\ V(\boldsymbol{v}_A) &= 0\ ,\end{aligned} \tag{39}$$

in number of photons per steradian.

Secondly, we have to sum over all the atoms having the given v_{A_Z} projected value for obtaining the frequency distribution of intensity and polarization of the reemitted radiation. Therefore we have to integrate over the atomic velocities in the XAY plane perpendicular to AZ. We obtain

$$\begin{aligned}I(\nu')\mathrm{d}\nu' &= \mathrm{d}v_{A_Z}\iint I(\boldsymbol{v}_A)\,\mathrm{d}v_{A_X}\mathrm{d}v_{A_Y}\ ,\\ Q(\nu')\mathrm{d}\nu' &= \mathrm{d}v_{A_Z}\iint Q(\boldsymbol{v}_A)\,\mathrm{d}v_{A_X}\mathrm{d}v_{A_Y}\ ,\end{aligned} \tag{40}$$

$$U(\nu')\mathrm{d}\nu' = \mathrm{d}v_{A_Z} \iint U\left(\boldsymbol{v}_A\right) \mathrm{d}v_{A_X} \mathrm{d}v_{A_Y} \ ,$$

$$V(\nu')\mathrm{d}\nu' = 0 \ ,$$

and

$$\mathrm{d}v_{A_Z} = \frac{c}{\nu_0} \, \mathrm{d}\nu' \ .$$

This is the general formula giving the Stokes parameters, as a function of frequency, of the reemitted line (in number of photons per unit of volume of matter, per unit of time, per unit of frequency and per unit of solid angle in the direction $\boldsymbol{n}'$) for any atomic velocity distribution. In fact, the formal formula (40) is not restricted to the two-level atom.

The degree of linear polarization p and the angle ϕ of the polarization direction with respect to the AX direction follow (cf. Fig. 2):

$$p = \frac{\sqrt{Q^2+U^2}}{I} \ , \tag{41}$$

$$\cos 2\phi = \frac{Q}{\sqrt{Q^2+U^2}} \ , \tag{42}$$

$$\sin 2\phi = \frac{U}{\sqrt{Q^2+U^2}} \ . \tag{43}$$

3. Application to the O VI line 103.2 nm of the solar corona

We begin by giving the expression of the polarization matrix of the incident radiation $J(\alpha, \beta, \nu)$ which is that of the same O VI line but emitted by the underlying transition region.

Neglecting inhomogeneities of the solar surface and using the same notations as Paper I, we can write $J(\alpha, \beta, \nu)$ as

$$J\left(\alpha, \beta, \nu\right) = J\left(\alpha\right) j_{\mathrm{i}}\left(\nu\right) \ , \tag{44}$$

where $j_{\mathrm{i}}\left(\nu\right)$ is the incident normalized profile.

$$J\left(\alpha\right) = I_{\mathrm{C}} f\left(\alpha\right) \ , \tag{45}$$

where $f\left(\alpha\right)$ is the limb-brightening coefficient ($f\left(0\right) = 1$) and I_{C} is the specific intensity emitted from the center of the disk in the line studied and integrated over the incident line profile.

The incident line profile is a gaussian Doppler profile, with a linewidth $\Delta\nu_{D_{\mathrm{i}}}$

$$j_{\mathrm{i}}(\nu) = \frac{1}{\sqrt{\pi}} \frac{1}{\Delta\nu_{D_{\mathrm{i}}}} \exp\left[-\left(\frac{\nu-\nu_0}{\Delta\nu_{D_{\mathrm{i}}}}\right)^2\right] \ , \tag{46}$$

and

$$\begin{aligned} \varphi_q^k\left(\boldsymbol{v}_A\right) &= \frac{I_{\mathrm{C}}}{4\pi} \int_0^{\alpha_{\max}} f(\alpha) \sin\alpha \ \mathrm{d}\alpha \int_0^{2\pi} \mathrm{d}\beta \\ &\times \frac{1}{\Delta\nu_{D_{\mathrm{i}}}} \frac{1}{\sqrt{\pi}} \exp\left[-\left(\frac{\nu_0}{c} \frac{\left(\boldsymbol{v}_A \cdot \boldsymbol{n}\right)}{\Delta\nu_{D_{\mathrm{i}}}}\right)^2\right] \mathcal{M}_q^k\left(\alpha, \beta\right) \ . \end{aligned} \tag{47}$$

Then we assume that the atomic velocity distribution can be written as

$$\begin{aligned} F\left(\boldsymbol{v}_A\right) \, \mathrm{d}^3\boldsymbol{v}_A &= \mathrm{d}^3\boldsymbol{v}_A \\ &\times \left(\frac{m_A}{2\pi kT}\right)^{3/2} \exp\left[-\frac{m_A}{2kT} \left|\boldsymbol{v}_A - \boldsymbol{V}\right|^2\right] \ , \end{aligned} \tag{48}$$

which corresponds to a Maxwellian distribution at a temperature T with an hydrodynamical drift of velocity $\boldsymbol{V}$ (V_X, V_Y, V_Z) or (V_x, V_y, V_z). This drift represents the velocity of the solar wind.

We can write, each subdistribution being normalized to 1

$$\begin{aligned} &F\left(\boldsymbol{v}_A\right) \, \mathrm{d}^3\boldsymbol{v}_A = \\ &F_X\left(v_{A_X}\right) \mathrm{d}v_{A_X} \, F_Y\left(v_{A_Y}\right) \mathrm{d}v_{A_Y} \, F_Z\left(v_{A_Z}\right) \mathrm{d}v_{A_Z} \ , \end{aligned} \tag{49}$$

with

$$\begin{aligned} &F_Z\left(v_{A_Z}\right) \mathrm{d}v_{A_Z} = \mathrm{d}v_{A_Z} \\ &\times \left(\frac{m_A}{2\pi kT}\right)^{1/2} \exp\left[-\frac{m_A}{2kT}\left(v_{A_Z} - V_Z\right)^2\right] \ , \end{aligned} \tag{50}$$

and analogous expressions for the two other distributions.

The integrations over v_{A_X} and v_{A_Y} can be achieved analytically by means of elementary methods and are not detailed here.

The Stokes parameters $I\left(\nu\right)$, $Q\left(\nu\right)$, $U\left(\nu\right)$ of the frequency distribution of intensity and polarization of the reemitted line at the frequency ν are given in Appendix A: Eq. (A1). They are expressed in number of photons emitted per steradian, per unit of volume, per unit of time and per unit of interval of frequency.

This equation has to be applied with $c_{ul} = 1/2$ (Li -like resonance line $2p^2P_{3/2} \rightarrow 2s^2S_{1/2}$) and $c_{ul} = 0$ for the other component of the doublet ($2p^2P_{1/2} \rightarrow 2s^2S_{1/2}$).

Without limb-brightening,

$$f\left(\alpha\right) = 1$$

and

$$\sin\alpha_{\max} = \frac{R_\odot}{R_\odot(1+h)} \ .$$

After integration over the scattered line-profile, one obtains the global Stokes parameters of the scattered line I, Q, U (in number of photons emitted per steradian, per unit of volume, and per unit of time). They are given in Appendix A: Eq. (A8).

These formulae show that the three first Stokes parameters (intensity and linear polarization parameters) are sensitive to the three components of the vector velocity of the drift. The order of magnitude of the sensitivity to these components will be the object of the next paper (cf. also Sahal-Bréchot et al. 1992b and Sahal-Bréchot & Choucq-Bruston 1994 for preliminary results).

3.1. Asymptotic limit: perfectly directive case

At great distances from the limb (*i.e.*, $h > 5$ in units of solar radius), the results corresponding to a perfectly directive incident radiation (Hyder & Lites 1970; Beckers & Chipman 1974), are recovered. By taking the limit $\alpha \rightarrow 0$, the integration of Eq. (A1) over α and β becomes analytic. The scattered line has also a gaussian profile, and the frequency dependence is the same for all the Stokes parameters: they are thus characterized by 5 parameters $w_\infty, d_\infty, p_\infty, s_\infty, \phi_\infty$ which are defined in the following.

Writing the intensity of the purely directive incident radiation as

$$I_{\rm i}\left(\nu\right) = I_0 \exp\left[-\left(\frac{\nu-\nu_0}{\Delta\nu_{D_{\rm i}}}\right)^2\right] ,$$

with

$$I_0 = I_{\rm C}\frac{1}{\sqrt{\pi}\sqrt{\Delta\nu_{D_{\rm i}}}} ,$$

we obtain a scattered intensity in the AZ direction

$$I_{\rm s}\left(\nu\right) = I \exp\left[-\left(\frac{\nu-\nu_0-d_\infty}{w_\infty}\right)^2\right] ,$$

with a width

$$w_\infty = \Delta\nu_{D_{\rm s}}\left(1-\frac{\Delta\nu_{D_{\rm s}}^2}{\Delta\nu_{D_{\rm i}}^2+\Delta\nu_{D_{\rm s}}^2}\cos^2\theta\right) , \quad (51)$$

where

$$\Delta\nu_{D_{\rm s}} = \frac{\nu_0}{c}\sqrt{\frac{2kT}{m_A}} ,$$

a shift

$$d_\infty = \frac{\nu_0}{c}\left(V_Z-\frac{\Delta\nu_{D_{\rm s}}^2}{\Delta\nu_{D_{\rm i}}^2+\Delta\nu_{D_{\rm s}}^2}V_z\cos\theta\right) , \quad (52)$$

a dimming

$$\frac{I}{I_0} = \exp\left[-s_\infty\right] = \exp\left[-\left(\frac{\frac{\nu_0}{c}V_z}{\sqrt{\Delta\nu_{D_{\rm i}}^2+\Delta\nu_{D_{\rm s}}^2}}\right)^2\right] , \quad (53)$$

which characterizes the dimming of the scattered line intensity with respect to the scattered intensity in the absence of velocity field, a degree of polarization

$$p_\infty = \frac{|Q|}{I} = \frac{3c_{ul}\sin^2\theta}{4-c_{ul}+3c_{ul}\cos^2\theta} \quad (54)$$

and

$$\phi_\infty = 0$$

which corresponds to a direction of polarization perpendicular to the scattering plane ($U=0$ and $Q>0$).

In particular, for a $J_u = 3/2 \rightarrow J_l = 1/2$ line and for $\theta = \pm\pi/2$, we obtain the well-known result

$$p_\infty = 0.428\,. \quad (55)$$

Therefore, at great distances from the surface of the sun the analysis of the scattered line intensity gives only two components of the velocity field: its projection on the direction of the incident radiation V_z and its projection on the line-of-sight V_Z through the interpretation of the dimming and of the shift. The other Stokes parameters do not bring any additional information. The third component of the velocity field V_X – or V_x –, *i.e.*, the perpendicular to the scattering plane, cannot be determined. Consequently the vectorial diagnostic cannot be performed with this method for a perfectly directive incident radiation.

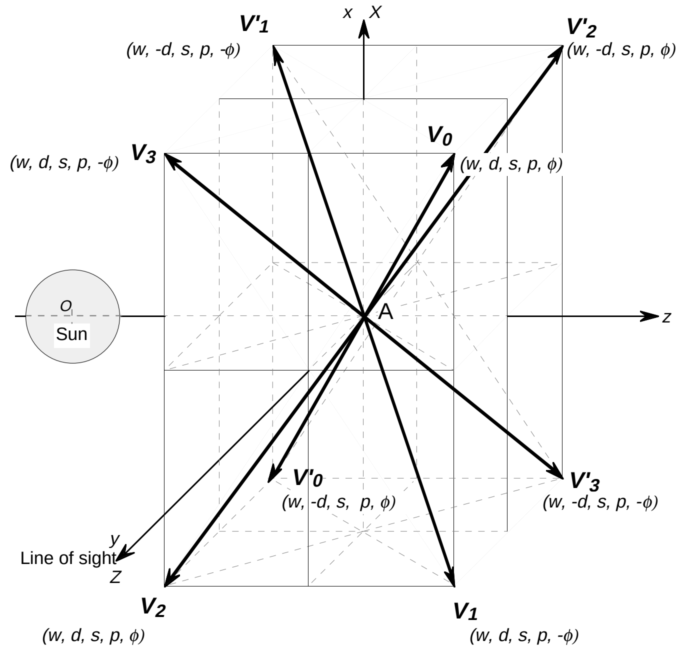

Fig. 4. The symmetries and degeneracies of the solution for the diagnostic of the velocity field vector. Oz is the preferred direction of incident radiation

3.2. *Symmetries of the problem*

The following discussion is restricted to a scattering atom in the plane of the sky ($\theta = \pm\pi/2$). We show that the results for $\pi/2 < \alpha_V < \pi$, and for $\pi/2 < \beta_V < 2\pi$ can be deduced from those of the first quadrant ($0 < \alpha_V < \pi/2$ and $0 < \beta_V < \pi/2$). The symmetries of the problem are illustrated on Fig. 4.

3.2.1. Symmetry due to the change of the direction of propagation of the incident light into the opposite sense

For the scattering atom, this is equivalent to

$$\begin{cases}\alpha_V \rightarrow \pi-\alpha_V\\ \beta_V \rightarrow \pi-\beta_V\end{cases} ,$$

and the line-of-sight is conserved. For $\theta = \pm\pi/2$, (the present figure is drawn for $\theta = -\pi/2$), this correspond to a symmetry with respect to the line-of-sight AZ.

As the recoil of the atomic motion due to absorption or emission of light is negligible, the atom is not sensitive to the sense of propagation of the light and two opposite directions of propagation have the same effect. Consequently the Stokes parameters of the scattered line will be the same for two velocity field vectors which are symmetrical with respect to the line-of-sight. This is the fundamental degeneracy, which exists in every polarimetric diagnostic, and which has been abundantly discussed in Hanle effect studies (Bommier & Sahal-Bréchot 1978 and further papers); $\boldsymbol{V}_0 \rightarrow \boldsymbol{V}_2$ on Fig. 4. In fact, if the scattering atom was not in the plane of the sky ($\theta \neq \pm\pi/2$), the two corresponding velocity field vectors would not be symmetrical with respect to the line-of-sight.

3.2.2. Symmetry with respect to the preferred direction of the incident radiation

$$\begin{cases} \alpha_V \to \alpha_V \\ \beta_V \to \beta_V + \pi \end{cases} .$$

The characteristics of the scattered radiation are the same for two opposite lines of sight, the only changes are the sign of the shift d of the scattered line

$$d \to -d$$

and that of the polarization rotation: I (apart from the sign of the shift) and Q are unchanged and the sign of U is changed

$$\phi \to -\phi \; ;$$

this symmetry exists also in Hanle effect studies. $\boldsymbol{V}_0 \to \boldsymbol{V}'_3$ on Fig. 4.

3.2.3. Symmetry with respect to the center of scattering A

$$\begin{cases} \alpha_V \to \pi - \alpha_V \\ \beta_V \to \beta_V + \pi \end{cases} ,$$

The only change is the sign of the shift

$$d \to -d$$

This symmetry does not exists in Hanle effect studies, because $\boldsymbol{V}$ is a real (polar)-vector, whereas the magnetic field $\boldsymbol{B}$ is an axial-vector; $\boldsymbol{V}_0 \to \boldsymbol{V}'_0$ on Fig. 4.

3.2.4. Symmetry with respect to the scattering plane

$$\begin{cases} \alpha_V \to \alpha_V \\ \beta_V \to \pi - \beta_V \end{cases} ;$$

this symmetry is in fact the product of the three preceding symmetries: the result is that I and Q are unchanged, and that the sign of U is changed: the rotation of polarization is changed in its opposite and the shift is unchanged; $\boldsymbol{V}_0 \to \boldsymbol{V}_1$ on Fig. 4.

4. Conclusion

We have rederived the basic theory of resonance scattering for an incident radiation partially directive and frequency-dependent, the moving scattering atoms having an anisotropic distribution of velocities. We have shown that the three first Stokes parameters of the reemitted line are sensitive to the three components of the macroscopic velocity field. Therefore the interpretation of these Stokes parameters may offer a method for deriving the complete velocity field vector. The formalism is general and be applied to a variety of astrophysical problems where anisotropies of resonance scattering occur.

Then we have applied the formalism to the two-level atom and we have given the expression of the Stokes parameters of the O VI 103.2 nm line of the solar corona formed in the solar wind acceleration region. This line which has a Doppler profile is formed by resonance scattering of the same line originating from the transition region and which has also a Doppler profile. The first numerical results of this formalism have been given by Sahal-Bréchot et al. (1992b) and Sahal-Bréchot & Choucq-Bruston (1994). They will be analysed and discussed in the next paper, in view of future interpretation of solar data obtained by SUMER-SOHO.

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Appendix A: stokes parameters of the scattered radiation

The Stokes parameters of the radiation scattered by a two-level atom having an ensemble velocity $\boldsymbol{V}$, as derived from the calculation described in the present paper, are given by

– ν-dependent Stokes parameters

$$\begin{aligned}\begin{pmatrix} I(\nu) \\ Q(\nu) \\ U(\nu) \end{pmatrix} \mathrm{d}\nu = \frac{\mathrm{d}\nu}{4\pi} N_{\mathrm{A}} \frac{1}{\sqrt{\pi}\Delta\nu_{\mathrm{D_s}}} \exp\left[-\left(\frac{\nu-\nu_0-\Delta\nu_Z}{\Delta\nu_{\mathrm{D_s}}}\right)^2\right] N_{\mathrm{e}}\,\alpha(J_l \to J_u, T_{\mathrm{e}}) \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix} \\ + \frac{\mathrm{d}\nu}{4\pi} N_{\mathrm{A}} I_{\mathrm{C}} B(J_l \to J_u) \frac{1}{\sqrt{\pi}\sqrt{\Delta\nu_{\mathrm{D_i}}^2 + \Delta\nu_{\mathrm{D_s}}^2}} \frac{1}{4\pi} \int_0^{\alpha_{\max}} f(\alpha) \sin\alpha \,\mathrm{d}\alpha \int_0^{2\pi} \mathrm{d}\beta \\ \times \exp\left[-s(\alpha,\beta)\right] \frac{1}{\sqrt{\pi}} \frac{1}{w(\alpha,\beta,\theta)} \exp\left[-\left(\frac{\nu-\nu_0-d(\alpha,\beta,\theta)}{w(\alpha,\beta,\theta)}\right)^2\right] \\ \times \left\{ \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix} + c_{ul} \begin{pmatrix} \frac{1}{8}\left(3\cos^2\theta - 1\right)\left(3\cos^2\alpha - 1\right) - \frac{3}{8}\sin 2\theta \sin 2\alpha \sin\beta - \frac{3}{8}\sin^2\theta \sin^2\alpha \cos 2\beta \\ \frac{3}{8}\sin^2\theta\left(3\cos^2\alpha - 1\right) + \frac{3}{8}\sin 2\theta \sin 2\alpha \sin\beta - \frac{3}{8}\left(1+\cos^2\theta\right)\sin^2\alpha\cos 2\beta \\ -\frac{3}{4}\sin\theta\sin 2\alpha\cos\beta - \frac{3}{4}\cos\theta\sin^2\alpha\sin 2\beta \end{pmatrix} \right\}\end{aligned}, \quad \text{(A1)}$$

with

$$\Delta\nu_{\mathrm{D_s}} = \frac{\nu_0}{c}\sqrt{\frac{2kT}{m_{\mathrm{A}}}}\ , \quad \text{(A2)}$$

$$\Delta\nu_Z = \frac{\nu_0}{c} V_Z = \frac{\nu_0}{c}\left(-V_y \sin\theta + V_z\cos\theta\right)\ . \quad \text{(A3)}$$

$$\Delta\nu_{\boldsymbol{n}} = \frac{\nu_0}{c}\left(\boldsymbol{n}\cdot\boldsymbol{V}\right) = \frac{\nu_0}{c}\left(V_x\sin\alpha\cos\beta + V_y\sin\alpha\sin\beta + V_z\cos\alpha\right)\ , \quad \text{(A4)}$$

$$s(\alpha,\beta) = \frac{\Delta\nu_{\boldsymbol{n}}^2}{\Delta\nu_{\mathrm{D_i}}^2 + \Delta\nu_{\mathrm{D_s}}^2}\ , \quad \text{(A5)}$$

$$w(\alpha,\beta,\theta) = \Delta\nu_{\mathrm{D_s}}\left[1 - \frac{\Delta\nu_{\mathrm{D_s}}^2}{\Delta\nu_{\mathrm{D_i}}^2 + \Delta\nu_{\mathrm{D_s}}^2}\left(-\sin\alpha\sin\beta\sin\theta + \cos\alpha\cos\theta\right)^2\right]^{1/2}\ , \quad \text{(A6)}$$

$$d(\alpha,\beta,\theta) = \Delta\nu_Z - \frac{\Delta\nu_{\mathrm{D_s}}^2}{\Delta\nu_{\mathrm{D_i}}^2 + \Delta\nu_{\mathrm{D_s}}^2}\left(-\sin\alpha\sin\beta\sin\theta + \cos\alpha\cos\theta\right)\Delta\nu_{\boldsymbol{n}}\ , \quad \text{(A7)}$$

– ν-integrated Stokes parameters

$$\begin{aligned}\begin{pmatrix} I \\ Q \\ U \end{pmatrix} = \frac{1}{4\pi} N_{\mathrm{A}} N_{\mathrm{e}}\,\alpha(J_l \to J_u, T_{\mathrm{e}}) \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix} \\ + \frac{1}{4\pi} N_{\mathrm{A}} I_{\mathrm{C}} B(J_l \to J_u) \frac{1}{\sqrt{\pi}\sqrt{\Delta\nu_{\mathrm{D_i}}^2 + \Delta\nu_{\mathrm{D_s}}^2}} \frac{1}{4\pi} \int_0^{\alpha_{\max}} f(\alpha)\sin\alpha\,\mathrm{d}\alpha \int_0^{2\pi} \mathrm{d}\beta\ \exp\left[-s(\alpha,\beta)\right] \\ \times \left\{ \begin{pmatrix} 1 \\ 0 \\ 0 \end{pmatrix} + c_{ul} \begin{pmatrix} \frac{1}{8}\left(3\cos^2\theta - 1\right)\left(3\cos^2\alpha - 1\right) - \frac{3}{8}\sin 2\theta \sin 2\alpha \sin\beta - \frac{3}{8}\sin^2\theta \sin^2\alpha \cos 2\beta \\ \frac{3}{8}\sin^2\theta\left(3\cos^2\alpha - 1\right) + \frac{3}{8}\sin 2\theta \sin 2\alpha \sin\beta - \frac{3}{8}\left(1+\cos^2\theta\right)\sin^2\alpha\cos 2\beta \\ -\frac{3}{4}\sin\theta\sin 2\alpha\cos\beta - \frac{3}{4}\cos\theta\sin^2\alpha\sin 2\beta \end{pmatrix} \right\}\end{aligned}. \quad \text{(A8)}$$

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