Orientation of hydrogen atoms by resonant radiation in a tenuous turbulent plasma

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A theoretical analysis is presented of the influence of an intense gyrofrequency plasma turbulence on the populations of the Zeeman sublevels of the $1S_{1/2}$ state of HI atoms in a tenuous turbulent plasma situated in the field of resonant radio emission ($\lambda = 21$ cm) or ultraviolet emission ($\lambda \sim 1216$–912 Å). The spin temperatures of the magnetic sublevels are determined and it is shown that in the plasma the criteria for producing invertered populations of the indicated sublevels differ substantially from those given by Varshalovich [Soviet Physics JETP 25, 157 (1967)]. The attenuation (amplification) coefficient of $\lambda = 21$ cm radio emission passing through the considered medium is obtained.

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The influence of plasma on the radiation of atoms, and the ensuing possibilities of plasma diagnostics, are the subject of a large number of studies.[1-9] This research developed predominantly in two directions. The first was connected with the study of the broadening of spectral lines in collisions of an atom with charged particles,[1,2] and the second with consideration of the influence of collective motions (plasma turbulence) on the radiation of atomic systems. In the second trend, one can separate two aspects. One is connected with research on the broadening of the spectral lines by the plasma turbulence,[3-5] whereas the other is connected with consideration of radiation plasmons by the atoms.[6-9]

We consider below hydrogen atoms situated in the field of $\lambda = 21$ cm radiation and of $\lambda \sim 1216-912$ Å ultraviolet radiation. A detailed analysis of orientation under these conditions was carried out by Varshalovich.[10] In the present paper we investigate the influence of the plasma turbulence, at frequencies corresponding to transitions between the Zeeman sublevels of the upper level of the hyperfine multiplet of the ground state $1S_{1/2}$ of the HI atoms, on the populations of the indicated sublevels.

We assume that all the atoms are on the levels of the $1S_{1/2}$ multiplet, and that their Zeeman splitting greatly exceeds the energy width, but is much less than the hyperfine splitting. This situation is typical, for example, of clouds of interstellar hydrogen.

1. The state of the atoms HI will be described by four quantities R_{FM}, which constitute the populations of the levels $F=0, M=0$ $F=1, M=0, \pm 1$ of the state $1S_{1/2}$, where $\mathbf{F}=\mathbf{J}+\mathbf{I}$ is the total angular momentum of the atom (the electron and the proton), and M is the projection of $\mathbf{F}$ on the quantization axis, assumed to be directed along the magnetic field.

We determine the spin temperatures $T_{s,+1}, T_{s0}, T_{s,-1}$:

$$R_{1M}/R_{00}=\exp(-\hbar\omega_0/\varkappa T_{sM}), \tag{1}$$

where $\hbar\omega_0$ is the hyperfine splitting of $1S_{1/2}$.[11] Since

$$T_*=\hbar\omega_0/\varkappa=0.0681\,\mathrm{K}\ll T_{sM},$$

we have

$$T_{sM}\approx T_*R_{00}/(R_{00}-R_{1M}). \tag{2}$$

The populations R_{FM} are determined from the system of stationarity equations

$$R_{FM}\sum_{F'M'}W_{FM\to F'M'}=\sum_{F'M'}R_{F'M'}W_{F'M'\to FM}, \tag{3}$$

where $W_{FM\text{-}F'M'}$ are the probabilities of the corresponding transitions. The general solution of (3) is of the form

$$\frac{R_{1,\pm1}}{R_{00}}=\frac{1}{N}\left\{\left(W_{00\to1,\pm1}+\frac{W_{00\to10}W_{10\to1,\pm1}}{W_{10}}\right)\left(\frac{W_{1,\mp1\to10}W_{10\to1,\mp1}}{W_{10}}-W_{1,\mp1}\right)\right.$$
$$\left.-\left(W_{00\to1,\mp1}+\frac{W_{00\to10}W_{10\to1,\mp1}}{W_{10}}\right)\left(\frac{W_{1,\mp1\to10}W_{10\to1,\pm1}}{W_{10}}+W_{1,\mp1\to1,\pm1}\right)\right\},$$
$$\frac{R_{10}}{R_{00}}=\frac{1}{N}\left\{\frac{W_{00\to10}}{W_{10}}(W_{11\to1,-1}W_{1,-1\to11}-W_{11}W_{1,-1})\right.$$
$$-\left(W_{00\to11}+\frac{W_{00\to1,-1}W_{1,-1\to11}}{W_{1,-1}}\right)\frac{W_{1,-1}}{W_{10}}W_{11\to10}$$
$$\left.-\left(W_{00\to1,-1}+\frac{W_{00\to11}W_{11\to1,-1}}{W_{11}}\right)\frac{W_{11}}{W_{10}}W_{1,-1\to10}\right\}, \tag{4}$$

where

$$W_{FM}\equiv\sum_{F'M'}W_{FM\to F'M'}$$

is the total probability of knocking out the atoms from a given state, and the factor N is given by

$$N=\left(W_{11\to1,-1}+\frac{W_{11\to10}W_{10\to1,-1}}{W_{10}}\right)\left(\frac{W_{1,-1\to10}W_{10\to11}}{W_{10}}+W_{1,-1\to11}\right)$$
$$+\left(W_{11}-\frac{W_{11\to10}W_{10\to11}}{W_{10}}\right)\left(\frac{W_{1,-1\to10}W_{10\to1,-1}}{W_{10}}-W_{1,-1}\right). \tag{5}$$

The probabilities $W_{FM\to F'M'}$ are determined in our analysis by four processes. We start with the transitions $F=1\to F=0$, which are connected with the $\lambda = 21$ cm radiation. The sought probabilities $W^{(R)}_{FM\to F'M'}$ are given in the case of a directed flux of polarized radiation by[10]

$$W^{(R)}_{FM\to F'M'}=\gamma_0(C^{FM}_{F'M'jm})^2(1+\rho^{(R)}_m), \tag{6}$$

where $\gamma_0 = 2.85\cdot 10^{-15}$ sec^{-1} is the total probability of the spontaneous $FM\to F'M'$ transition, $C^{FM}_{F'M'jm}$ are Clebsch–Gordan coefficients, and the quantities $\rho^{(R)}_m$, which are the diagonal matrix elements of the radiation density in the angular-momentum representation, take the form

$$\rho^{(R)}_{\pm1}(\omega)=\frac{3}{2}N^{(R)}_\omega\frac{\Omega_R}{4\pi}\left[1\pm\eta\cos\theta-\frac{1}{2}\sin^2\theta(1+\xi\cos 2\alpha)\right],$$
$$\rho^{(R)}_0(\omega)=\frac{3}{2}N^{(R)}_\omega\frac{\Omega_R}{4\pi}\sin^2\theta(1+\xi\cos 2\alpha). \tag{7}$$

Here $N^{(R)}_\omega$ is the dimensionless number of quanta in the phase-space cell, Ω_R is the solid angle subtended by the radiation source ($\Omega_R/4\pi\ll 1$). Next, η and ξ are the degrees of circular and linear polarization, θ is the angle between the direction of the flux and the direction of the magnetic field **H**, and α is a position angle that characterizes the linear polarization.

We now stop to consider transitions between levels of the hyperfine structure of $1S_{1/2}$, which are due to resonant scattering of ultraviolet radiation $\lambda\sim 126-912$ Å (Lyman-series transitions). The expressions for the probabilities were obtained previously.[10] Without presenting the rather cumbersome concrete expressions, we note that these probabilities are proportional to the quantities $\rho^L_m(\omega)$ that describe the field of the ultraviolet radiation and are similar in form to (7).

Transitions are also possible between the levels of the $1S_{1/2}$ hyperfine structure as a result of collisions of the atoms with one another or their collisions under nonequilibrium conditions is an independent problem. Assuming the medium to be sufficiently tenuous, we neglect henceforth these transitions.

We now dwell on transitions with participation of plasma waves. Transitions between Zeeman sublevels of one and the same level are possible in atoms situated in a plasma at frequencies $g\omega_{He}/2$ is the Landé factor and ω_{He} is the electron gyrofrequency) which proceed with emission (absorption) of waves in the plasma at these frequencies.[12] The values of the Lande factor for different terms vary most frequently from fractions to several times unity.[13] In the frequency interval corresponding to the indicated variations of the g factor, whistlers and longitudinal gyrofrequency waves propagate in the plasma. Despite the low frequency, transitions with participation of plasma waves can be quite effective. One must bear in mind here two circumstances. First, the refractive indices of the waves at these frequencies can greatly exceed unity; second, the wave energy density in a turbulent plasma is high. Expressions for the probabilities of magnetic-dipole transitions between Zeeman sublevels of a hyperfine structure level in a plasma can be easily obtained from the formulas of Ref. 14:

$$W_{FM\to F'M'}=W_F(C^{FM}_{FM'1m})^2\rho^{(p)}_m(\omega),\quad W_F=\frac{2}{3}\frac{\omega^3\mu_0^2}{\hbar c^3}g_F^2F(F+1). \tag{8}$$

Here μ_0 is the Bohr magneton, g_F is the Lande factor, $\omega=g_F\omega_{He}/2$ is the frequency at which the transition takes place. The quantities $\rho^{(p)}_m(\omega)$ are given by

$$\rho^{(p)}_{\pm1}(\omega)=\frac{3}{2}N^{(p)}_\omega\frac{\Omega_p}{4\pi}n^3A^2\cos^2\theta\left(1-\frac{n^2\sin^2\theta}{n^2\sin^2\theta-\varepsilon_2}\mp\frac{\varepsilon_2}{n^2-\varepsilon_1}\right)^2\times\left[1-\frac{c^2(\mathbf{kl}^*)(\mathbf{kl})}{n^2\omega^2}\right]^{-1},$$
$$\rho^{(p)}_0(\omega)=\frac{3}{2}N^{(p)}_\omega\frac{\Omega_p}{4\pi}n^3A^2\frac{\varepsilon_2^2\sin^2\theta}{(n^2-\varepsilon_1)}\left[1-\frac{c^2(\mathbf{kl}^*)(\mathbf{kl})}{n^2\omega^2}\right]^{-1}. \tag{9}$$

Here $N^{(p)}_\omega$ is a dimensionless number of plasma waves, **k** and **l** are the wave vector and polarization vector, while A is given by

$$A=\left[1+\frac{\varepsilon_2^2}{(n^2-\varepsilon_1)^2}+\frac{n^4\sin^2\theta\cos^2\theta}{(n^2\sin^2\theta-\varepsilon_2)^2}\right]^{-1/2}. \tag{10}$$

The expressions for the refractive index n and the components ε_1, ε_2, ε_3 of the dielectric tensor of the plasma are given in Ref. 15. In (9) and (10) we have left out, for simplicity, the indices that label the ordinary and extraordinary waves.

2. We consider now the case when the principal mechanism of excitation of the atoms is their interaction with the radio emission:

$$W^{(R)}_{FM\to F'M'}\gg W^{(L)}_{FM\to F'M'}. \tag{11}$$

The probabilities of the concrete transitions for the HI atom under the influence of the radio emission can be easily obtained from (6):

$$W^{(R)}_{00\to1M}=\gamma_0\rho^{(R)}_M,\quad W^{(R)}_{1M\to00}=\gamma_0(1+\rho^{(R)}_M). \tag{12}$$

In addition, we take into account the transitions between the Zeeman sublevels of the $F=1$ level [at the frequency $(1/2)\omega_{He}$] under the influence of a narrow intense beam of whistlers propagating along the magnetic field (we assume henceforth that the plasma frequency ω_{pe} greatly exceeds ω_{He}). The probabilities of these transitions can be easily obtained from (8)–(10):

$$W_{1M\to1M'}=\frac{2\omega_{pe}^3\mu_0^2}{\hbar c^3}(C^{1M}_{1M'11})^2N^{(p)}_\omega\frac{\Omega_p}{4\pi}. \tag{13}$$

In a strongly turbulent plasma, this process can become more effective than interaction with the radio emission. The corresponding condition can be easily obtained from (12) and (13). For the case of an isotropic unpolarized radio-emission flux we have

$$\frac{T_p}{T_R}\gg\frac{\omega_{He}\omega_0^2}{3\omega_{pe}^3}\frac{4\pi}{\Omega_p}. \tag{14}$$

To obtain (14) we have introduced the effective temperatures T_p and T_R of the whistlers and of the radio emission, respectively:

$$T_p=\frac{\hbar\omega_{He}}{2\varkappa}N^{(p)}_\omega,\quad T_R=\frac{\hbar\omega_0}{\varkappa}N^{(R)}_{\omega_0}. \tag{15}$$

Using (4) and (5), we obtain the corresponding populations under the indicated assumptions, and relation (2) yields the spin temperatures.

Simple calculations give:

$$T_{S_1}\approx T_{S_0}\approx T_{S,-1}\approx T_R. \tag{16}$$

This expression is accurate up to terms $\sim T_*$. This result is close to that obtained in Ref. 10.

In the case of a directed beam of polarized radiation, the temperatures are different, but differ little from the corresponding ones in Ref. 10.

3. We consider now the case when the intensity of the $\lambda \sim 1216-912$ Å radiation is high enough to make the principal excitation mechanism of the hydrogen atoms their interaction with the ultraviolet radiation. We assume in addition that the following condition is satisfied:

$$W^{(p)}_{FM\to F'M'} \gg W^{(L)}_{FM\to F'M'}. \quad (17)$$

Using (2), (4), and (5) we obtain expressions for the spin temperatures, assuming that the atoms are excited by an unpolarized isotropic flux of ultraviolet radiation:

$$\begin{aligned} &T_{11}=C/B_1, \quad T_{1,-1}=C/B_2, \quad T_{10}=C/B_3; \\ &C=(3\rho^{(p)2}+3\rho^{(p)}+1)\,(\rho(\omega_1)+\rho(\omega_2))\,T_*, \\ &B_1=3\rho^{(p)2}[\rho(\omega_1)+\rho(\omega_2)-\rho(\bar\omega_1)-\rho(\bar\omega_2)] \\ &\quad+3\rho^{(p)}[\rho(\omega_1)+\rho(\omega_2)]+\rho(\omega_1)+\rho(\omega_2), \\ &B_2=3\rho^{(p)2}[\rho(\omega_1)+\rho(\omega_2)-\rho(\bar\omega_1)-\rho(\bar\omega_2)] \\ &\quad-3\rho^{(p)}[\rho(\bar\omega_1)+\rho(\bar\omega_2)] \\ &\quad+3\rho^{(p)}[\rho(\omega_1)+\rho(\omega_2)-\rho(\bar\omega_1)-\rho(\bar\omega_2)]-3[\rho(\bar\omega_1)+\rho(\bar\omega_2)], \\ &B_3=3\rho^{(p)2}[\rho(\omega_1)+\rho(\omega_2)-\rho(\bar\omega_1)-\rho(\bar\omega_2)] \\ &\quad+3\rho^{(p)}[\rho(\omega_1)+\rho(\omega_2)-\rho(\bar\omega_1)-\rho(\bar\omega_2)]+\rho(\omega_1)+\rho(\omega_2). \end{aligned} \quad (18)$$

Here ω_1, $\bar\omega_1$, ω_2, and $\bar\omega_2$ are the respective frequencies of the transitions

$$\begin{aligned} &2P_{1/2}(F=1) \leftrightarrow 1S_{1/2}(F'=1), \\ &2P_{1/2}(F=1) \leftrightarrow 1S_{1/2}(F'=0); \\ &2P_{3/2}(F=1) \leftrightarrow 1S_{1/2}(F'=1), \\ &2P_{3/2}(F=1) \leftrightarrow 1S_{1/2}(F'=0). \end{aligned}$$

In (18) and subsequently we rewrite the index L of the quantities ρ describing the ultraviolet radiation.

We now analyze the obtained expressions (18). At a constant spectrum

$$\rho(\omega_1)=\rho(\bar\omega_1), \quad \rho(\omega_2)=\rho(\bar\omega_2),$$

we have

$$T_{11}\approx T_*\rho^{(p)}, \quad T_{1,-1}\approx -T_*\rho^{(p)}, \quad T_{10}\approx 3T_*\rho^{(p)2}. \quad (19)$$

It follows from (19) that $T_{1,-1}<0$, i.e., circularly polarized radiation (with right-hand polarization) at $\lambda=21$ cm can build up. It is important to note that if no account were taken of the plasma there would be no inverted population in this case.[10]

The presence of intense ($N^{(p)}_\omega \gg 1$) plasma waves with circular polarization singles out a preferred direction of the transitions

$$M=1\to M=0\to M=-1$$

between the sublevels of the level $F=1$ against the background of their uniform population by an isotropic flux of unpolarized ultraviolet radiation with a constant spectrum. The probabilities of these transitions exceed those for the competing processes in which ultraviolet radiation takes part [see the condition (17)]. These circumstances lead to a repopulation of the sublevel $F=1$, $M=-1$.

The inverted population can occur also in the case of isotropic unpolarized radiation with a rapidly growing spectrum in the region of $\lambda \sim 1216$ Å (see Ref. 10). This takes place also in our case. Thus, T_{10} can become negative under the condition

$$\frac{d\rho}{d\omega} > \frac{\rho(\omega_1)+\rho(\omega_2)}{6\rho^{(p)2}\omega_0}. \quad (20)$$

To obtain negative T_{11} it is necessary to satisfy the much more stringent condition

$$\frac{d\rho}{d\omega} > \frac{\rho(\omega_1)+\rho(\omega_2)}{3\rho^{(p)}\omega_0}. \quad (21)$$

The criteria (20) and (21) obtained from (18) differ substantially from the analogous ones in Ref. 10, and the first of them can be much more effective than those indicated in the case of developed whistler turbulence.

We note also that there is no inverted population in the case when the atoms are excited by unpolarized radiation directed along the quantization axis, with a constant spectrum.

4. The anisotropic properties of a medium containing oriented hydrogen atoms are characterized by a gain (or attenuation) τ of the passing radiation. The attenuation coefficient (or gain) of $\lambda=21$ cm circularly polarized radiation is given by (the plus and minus signs pertain respectively to right- and left-hand polarization)[10]

$$\tau_\pm = -\frac{3}{8\pi}\gamma_0 S(\omega) N R_{00}\frac{T_*}{T_{S\pm}}\left[1+\frac{\sin^2\theta_H}{2}\left(\frac{T_{S\pm}-T_{S0}}{T_{S0}}\right) \pm\cos\theta_H\left(\frac{T_{S1}-T_{S,-1}}{T_{S1}+T_{S,-1}}\right)\right], \quad (22)$$

where

$$S(\omega)=\frac{1}{\Delta\omega_D\sqrt{\pi}}\exp\left[-\left(\frac{\omega-\omega_0}{\Delta\omega_D}\right)^2\right]$$

is the Doppler contour of the spectral line

$$\Delta\omega_D=\omega_0(\varkappa T/mc^2)^{1/2}, \quad N=\lambda^2\int n_{at}(r)\,dr$$

is the number of hydrogen atoms in a column of cross section λ^2[cm^2] along the line of sight, and the angle θ_H characterizes the observation direction. In (22) and below we normalize the population to unity, and

$$T_{S\pm}=2T_{S1}T_{S,-1}/(T_{S1}+T_{S,-1}).$$

The attenuation (or gain) of linearly polarized $\lambda=21$ cm radiation is given by[10]

$$\tau_0 = -\frac{3}{8\pi}\gamma_0 S(\omega) N R_{00}\frac{T_*}{T_{S\pm}}\left[1+\frac{\sin^2\theta_H}{2}\left(\frac{T_{S\pm}-T_{S0}}{T_{S0}}\right)(1+\cos 2\alpha)\right]. \quad (23)$$

Using Eqs. (18) and (22) in the case of a constant ultraviolet spectrum, we obtain an expression for the gain of circularly polarized light with right-hand polarization:

$$\tau_+ = \frac{3}{8\pi}\gamma_0 S(\omega) N R_{00}\frac{\cos\theta_H}{\rho^{(p)}}. \quad (24)$$

For the case of ultraviolet radiation with a rapidly growing spectum (20) we obtain, using (18) and (23), an expression for the gain of linearly polarized radiation:

$$\tau_0 = \frac{3}{8\pi}\gamma_0 S(\omega) N R_{00}\sin^2\theta_H(1+\cos 2\alpha)\frac{\omega_0 d\rho/d\omega}{\rho(\omega_1)+\rho(\omega_2)}. \quad (25)$$

It is interesting to note that under these conditions the gain is independent of the characteristics of the plasma turbulence.

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Diamagnetism of cyclotron waves in plasma

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The nonlinear diamagnetic correction to the magnetic field, needed to account for the high-frequency pressure of a wave packet on the plasma, is obtained for potential cyclotron waves (Bernstein modes) in the case of arbitrary spatial dispersion, with account taken of the nonlinearity of the cyclotron resonance. With allowance for this correction, simplified equations are obtained for the evolution of the cyclotron waves. In the case of ion-cyclotron oscillations on the first harmonic, numerical methods are used to obtain the stationary solutions of these equations. In the one-dimensional case of the solution is a truncated periodic wave, while the cylindrical-symmetry case of the solution is a cyclotron soliton stretched along the external magnetic field. The formation of a magnetic well leads to self-localization, as a result of which the high-frequency pressure of the cyclotron waves can become comparable with the plasma pressure even in the case of weak instability and when the region of existence of the instability is small. This leads to anomalous resistance and to an increase of the particle energy in a direction perpendicular to the magnetic field. Similar effects might be observed in a magnetized semiconductor plasma.

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1. INTRODUCTION

A phenomenon intensively investigated during the last decade in plasma theory is the abrupt increase in the interaction between the Fourier components of the oscillations when the energy density of the latter reaches a certain threshold—the limit of applicability of the weak-turbulence approximation. As indicated in Refs. 1–4 and elsewhere, in the case of Langmuir waves this amplification leads to strongly nonlinear effects, namely to wave collapse or to formation of a set of solutions.

It is of interest to ascertain whether similar effects occur for cyclotron waves, which play very frequently an important role in plasma behavior. This is evidenced, in particular, by the experimental observation of these waves in tokamaks[5,6] and in the auroral region of the earth's magnetosphere,[7] from which intense radiation of electromagnetic waves was observed in the electron cyclotron frequency band (the kilometer radiation of the earth).[8,9] In addition, it was shown theore tically that in an isothermal plasma, in which no ion sound can build up, the principal mechanism of the anomalous resistance can be the interaction of the electrons with the ions via buildup of ion-cyclotron waves.[10]

The first step in the study of strongly nonlinear cyclotron wave is to obtain simplified equations that take into account only the principal linear and nonlinear effects. In Langmuir waves, owing to the decisive dependence of the Langmuir frequency on the plasma density, the main mechanism that leads to the aforementioned strongly nonlinear effects is the formation of density wells in the localization region of the wave packet. The frequency of the cyclotron waves depends mainly on the external magnetic field. Therefore, as shown in Ref. 11 and as will be shown more accurately and in greater detail in the the present paper, the main nonlinear mechanism that determines the behavior of the cyclotron waves is the formation of wells of a constant magnetic field in the region of localization of the cyclotron waves.

The decrease of the magnetic field in the region of localization of the wave packets, called high-frequency (HF) diamagnetism,[11] was discussed in Refs. 12–14. A self-focusing influence of HF diamagnetism on Alfven waves was observed in Ref. 15.

The HF diamagnetism is caused by the diamagnetic current produced in the plasma in a direction perpendicular to the constant magnetic field in the region of