# On the theory of polarization phenomena in the spectroscopy of two-quantum transitions

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(Submitted November 8, 1973)
Zh. Eksp. Teor. Fiz. **66**, 1570–1577 (May 1974)

The effect is considered of depolarizing collisions on the shape of the absorption or of the amplification spectrum of a weak wave in the presence of a strong wave which is in resonance with an adjacent transition. Also, the dependence of the shape of the spectrum on the polarizations of the waves and on the values of the total angular momenta of the levels is investigated. It is shown that the determination of the parameters of spontaneous and collision relaxations can be significantly simplified by utilizing for this purpose variation of polarizations and the method of "magnetic scanning."

1. As is well known the construction of gas lasers has served as a stimulus for a sharp increase in interest in resonance phenomena in combination scattering spectra. One of most effective methods of nonlinear spectroscopy has become the method based on the study of the special features of resonance interaction of a weak and a strong[1)] travelling wave with two transitions having one common level (cf., the review[1]). An analysis of the absorption or of the amplification spectra of a weak signal enables one to draw conclusions concerning different physical processes in the gas (in particular, concerning the spontaneous and collision relaxations) and to determine important atomic or molecular characteristics.

In the process of collision an excited particle can alter not only its velocity but also the orientation of its angular momentum. But whereas the effect on combination scattering spectra of collisions altering the velocity has been investigated in sufficient detail[2], the question of the role played by depolarizing collisions (i.e., collisions which give rise to mixing with respect to Zeeman sublevels) remains completely open. Our paper is devoted to solving it. In addition the effect is studied of the polarization of the waves on the line shape of a weak signal, and also the possibilities are considered associated with placing a three-level system into a magnetic field. It should be noted that the latter questions have been partially investigated experimentally[3-4] and the theory presented below agrees with the results of these experiments.

2. We assume that the active medium is situated in an external electromagnetic field $\tilde{\mathbf{E}}$ represented by a strong and a weak travelling wave corresponding to resonance transitions $m \to n$, $m \to l$:

$$\tilde{\mathbf{E}} = \frac{1}{2}[\mathbf{E}e^{-i(\omega t - \mathbf{kr})} + \mathbf{E}_\mu e^{-i(\omega_\mu t - \mathbf{k}_\mu \mathbf{r})}]. \tag{1}$$

The subscript $\mu$ characterizes a weak field, $\omega = |\mathbf{k}|c$, $\omega_\mu = |\mathbf{k}_\mu|c$ are the frequencies of the waves. In the resonance approximation the presence of all the levels of the atom with the exception of $m$, $n$, $l$ can be neglected, and we can restrict ourselves in future to a three-level model of the medium.

The interaction of the medium with the waves is described by the matrix elements:

a) for a strong field

$$V_{mM,nM'} = \sum_{\alpha=0,\pm1} G_{-\alpha} \begin{pmatrix} J_m & 1 & J_n \\ -M & \alpha & M' \end{pmatrix} (-1)^{J_m - M + \alpha}, \quad G_\alpha = \frac{E_\alpha d_{mn}}{2\hbar}; \tag{2}$$

b) for a weak field

$$V_{mM,lM'} = \sum_{\alpha=0,\pm1} G_{-\alpha} \begin{pmatrix} J_m & 1 & J_l \\ -M & \alpha & M' \end{pmatrix} (-1)^{J_m - M + \alpha}. \tag{3}$$

Here $d_{mn}$ and $d_{ml}$ are the reduced matrix elements of the dipole moment, while the electric vectors of the fields are resolved into the components

$$E_0 = E_z, \quad E_{\pm1} = \mp\frac{E_x \pm iE_y}{\sqrt{2}}; \quad E_0^\mu = E_z^\mu, \quad E_{\pm1}^\mu = \mp\frac{E_x^\mu \pm iE_y^\mu}{\sqrt{2}}. \tag{4}$$

The collision integrals for the elements of the density matrix $\rho_{iM,jM'}$ ($i, j = m, n, l$) have the form[5,6]

$$S_{iM,jM'} = -\sum_{\varkappa q M_1, M_1'} \gamma_{ij\varkappa}(2\varkappa+1)(-1)^{2J_i - M - M_1} \begin{pmatrix} J_i & \varkappa & J_j \\ -M & q & M' \end{pmatrix} \begin{pmatrix} J_i & \varkappa & J_j \\ -M_1 & q & M_1' \end{pmatrix} \rho_{iM_1,jM_1'} \tag{5}$$

and describe relaxation associated with mixing with respect to the magnetic sublevels (J is the total angular momentum of the level, M is the component of the angular momentum). The constants $\gamma_{ij}$ in (5) are expressed in the usual manner[5,6] in terms of the scattering matrices in the plane of the collisions. For given labels of the levels these constants depend only on the value of the angular momentum $\kappa$ ($|J_i - J_j| \le \kappa \le |J_i + J_j|$). Thus, for elastic collisions it follows from the law of conservation of the number of particles occupying a given level that $\gamma_{jj\kappa} = 0$ for $\kappa = 0$. If we neglect the fact that the remaining constants $\gamma_{jj\kappa}$ may be different, we shall be led to the strong collision model utilized in[7]. However, these constants may to a certain extent differ from one another[8-10], and therefore we shall here use the collision integral in the form (5).

From (2) and (5) one can see that in solving the equations for the density matrix which have a standard form it is convenient to go over to the representation of irreducible tensor operators in which the collision integrals become diagonalized. Following the method proposed by D'yakonov and Perel'[8,11], we obtain after simple but awkward calculations an expression for the work done by the weak field $P_\mu$ in the case when the waves are travelling in the same direction ($\mathbf{k} \uparrow\uparrow \mathbf{k}_\mu$) and the medium is placed in a magnetic field of intensity H:

$$P_\mu \sim \mathrm{Re}\Big[\Delta N_\mu \sum_\alpha |G_\alpha^\mu|^2 \exp\{-(\Omega_\mu + \alpha\Delta)^2/(k\bar{v})^2\}$$
$$-\Delta N \sum_{\substack{\varkappa q \alpha \\ \alpha_1 \alpha_2 \alpha_3}} \begin{pmatrix} 1 & 1 & \varkappa \\ \alpha_2 & -\alpha & q \end{pmatrix} \begin{pmatrix} 1 & 1 & \varkappa \\ \alpha_3 & -\alpha_1 & q \end{pmatrix} \frac{6(2\varkappa+1)(-1)^{\alpha_2+\alpha_3}}{\Gamma_{mn1} + \Gamma_{ml1} - i\varepsilon - i\Delta(\alpha - \alpha_3)}$$

  

$$\times\left(\frac{\beta_{nl\varkappa}G_\alpha^{\ \mu*}G_{\alpha_1}^{\ \mu}G_{\alpha_2}G_{\alpha_3}^{*}}{\Gamma_{nl\varkappa}-i\varepsilon-iq\Delta}+\frac{\beta_{m\varkappa}G_\alpha^{\ \mu*}G_{\alpha_1}G_{\alpha_2}^{\ \mu}G_{\alpha_3}^{*}}{\Gamma_{mm\varkappa}-iq\Delta}\right)\Bigg], \tag{6}$$

$$\Delta N=Q_m/\Gamma_{mm0}-Q_n/\Gamma_{nn0},\quad \Delta N_\mu=Q_m/\Gamma_{mm0}-Q_l/\Gamma_{ll0}.$$

Here $Q_j$ is the rate of excitation of the state j(j = m, n, $l$), while the coefficients $\beta_{nl\kappa}$ and $\beta_{m\kappa}$ are expressed in terms of Wigner's 6J-symbols:

$$\beta_{nl\varkappa}=\begin{Bmatrix}\varkappa & 1 & 1\\ J_m & J_n & J_l\end{Bmatrix}^2,$$
$$\beta_{m\varkappa}=(-1)^{J_l-J_n}\begin{Bmatrix}\varkappa & 1 & 1\\ J_n & J_m & J_m\end{Bmatrix}\begin{Bmatrix}\varkappa & 1 & 1\\ J_l & J_m & J_m\end{Bmatrix}. \tag{7}$$

In deriving expression (6) we chose the z axis of quantization along the magnetic field vector **H** and restricted our consideration to weak saturation. Moreover, we utilized the assumption usual for gas lasers that the constants

$$\Gamma_{ij\varkappa}=\Gamma_{ij}+\gamma_{ij\varkappa} \tag{8}$$

are much narrower than the Doppler width of the line $k\bar{v}$. From (6) it may be seen that it is useful to regard $P_\mu$ as a function of the frequency $\epsilon = \Omega_\mu - \Omega$ which is equal to the difference in the detuning of the weak and the strong fields with respect to the frequency of its own transition ($\Omega_\mu = \omega_\mu - \omega_{ml}$, $\Omega = \omega - \omega_{mn}$). The function $P_\mu(\epsilon)$ has the form of broad Doppler contours and of narrow dips (peaks) of Lorentz shape situated against their background. The widths of the dips contain information concerning spontaneous damping, which is taken into account by the quantities $\Gamma_{ij} = (\Gamma_{ii} + \Gamma_{jj})/2$, and concerning collision relaxation characterized by the constants $\gamma_{ij\kappa}$. The positions of the dips along the frequency scale can be altered depending on the value of the Zeeman splitting $\Delta = \mu_0 gH/\hbar$ (the Lande factors g of the m, n and $l$ levels are assumed to be equal) and on the direction of the magnetic field **H**.

Expression (6) has been obtained for values of k and $k_\mu$ close to each other ($|k - k_\mu|\bar{v} \ll \Gamma_{ij}$). However, it can be easily generalized to the case when k and $k_\mu$ are significantly different. For this it is sufficient to carry out in formula (6) the following replacements:

$$\varepsilon\to\frac{k}{k_\mu}\Omega_\mu-\Omega,\quad \Delta(\alpha-\alpha_3)\to\Delta\left(\frac{k}{k_\mu}\alpha-\alpha_3\right), \tag{9}$$
$$\Gamma_{ml1}\to\frac{k}{k_\mu}\Gamma_{ml1},$$
$$\Gamma_{nl\varkappa}\to\begin{cases}\Gamma_{nl\varkappa}+\Gamma_{mn1}(k_\mu/k-1) & \text{if}\quad k_\mu>k,\\ \Gamma_{nl\varkappa}+\Gamma_{mn1}(k/k_\mu-1) & \text{if}\quad k_\mu<k.\end{cases}$$

With the aid of the replacement

$$\varepsilon\to\frac{k}{k_\mu}\Omega_\mu+\Omega,\quad \Delta(\alpha-\alpha_3)\to-\Delta\left(\frac{k}{k_\mu}\alpha+\alpha_1\right),$$
$$\Gamma_{ml1}\to\frac{k}{k_\mu}\Gamma_{ml1},\quad \beta_{nl\varkappa}\to 0 \tag{10}$$

one can also easily obtain from (6) an expression for the work done by the weak field corresponding to waves travelling in opposite directions ($k \downarrow\uparrow k_\mu$).

In future we shall discuss in greater detail only the simple situation when the propagation vectors have values close to one another ($k \approx k_\mu$), having in mind that the analysis of the case $k \neq k_\mu$ can be carried out in an analogous manner by utilizing the substitutions (9) and (10).

3. The description of the amplification of a weak wave is simplest in the spontaneous relaxation approximation which is valid at low pressures. As can be deduced from (6), in this case for waves travelling in the same direction and for zero magnetic field (the axis of quantization coincides with k) we have

$$P_\mu\sim\sum_{\alpha=\pm1}\mathrm{Re}\left[\Delta N_\mu\exp\{-\Omega_\mu^2/(k\bar{v})^2\}-\frac{C\Delta N}{\Gamma_{mm}(\Gamma_{nl}-i\varepsilon)}\right], \tag{11}$$

$$C=A_0|G_\alpha^{\ \mu}|^2|G_\alpha|^2+A_1|G_\alpha^{\ \mu}|^2|G_{-\alpha}|^2+A_2G_\alpha G_{-\alpha}^{*}G_{-\alpha}^{\ \mu}G_\alpha^{\ \mu*},$$
$$A_{0,2}=2\beta_{nl0}\pm\beta_{nl1}+3\beta_{nl2},\quad A_1=6\beta_{nl2}. \tag{12}$$

The frequency dependence $P_\mu(\epsilon)$ here contains against a Doppler background one line of width $\Gamma_{nl}$ and with its center at $\epsilon = 0$. Further, for the sake of definiteness we assume that $\Delta N$ and $\Delta N_\mu$ have the same sign and we speak of a dip. From (11) and (7) it can be seen that the depth of the dip depends in an essential manner on the relationship of the polarizations of the strong and the weak waves and on the values of the total angular momenta of the working levels $J_m$, $J_n$ and $J_l$. For certain values of these parameters the dip can disappear completely. Such a situation occurs, for example, for waves with opposite circular polarizations, when $C = A_1|G_\mu|^2|G|^2$, and for a three-level system with total angular momenta $J_n = 0$, $J_m = J_l = 1$. Figure 1 corresponds to this situation. The solid lines in it correspond to a strong field and the wavy ones to a weak field. Since the weak wave acts on transitions between sublevels unperturbed by the strong field (i.e., on transitions m, M = 0 → $l$, M = – 1 and m, M = 1 → $l$, M = 0), naturally the dip is absent. An exact calculation utilizing formula (11) also yields C = 0, since (cf.,[12])

$$A_1=6\begin{Bmatrix}2 & 1 & 1\\ 1 & 0 & 1\end{Bmatrix}^2=0.$$

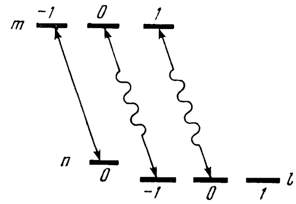


FIG. 1. Scheme of transitions between Zeeman sublevels of levels m, $J_m = 1$; n, $J_n = 0$; $l$, $J_l = l$. The strong and the weak waves are circularly polarized in opposite directions.

The polarization properties of the line shape for combination scattering were investigated experimentally in[3-4]. In these papers measurements were reported of the absorption spectrum for a weak field interacting with the transitions $2s_2 - 2p_4$ ($J_{2s_2} = 1$, $J_{2p_4} = 2$) or $2s_2 - 2p_2$)$J_{2p_2} = 1$) of the $Ne^{20}$ isotope under the action of a strong field on the adjacent transition $2s_2 - 2p_1$ ($J_{2p_1} = 0$). For the system of levels $2p_1 - 2s_2 - 2p_2$ there was observed, in particular, the effect described here of the disappearance of the narrow structure. Moreover, for the system $2p_1 - 2s_2 - 2p_4$ a sixfold increase in the depth of the dip was recorded in going over from waves circularly polarized in the same sense to waves with orthogonal circular polarizations. The same result can also be easily obtained with the aid of the theory developed above. Indeed, from (11) it follows directly that the ratio of the depths of the dips measured for waves with the same and with opposite circular polarizations is equal to $A_0 : A_1$; with this ratio having the value $A_0 : A_1 = 1/6$ for transitions with total angular momenta $J_n = 0$, $J_m = 1$, $J_l = 2$ and, in general, for all transitions of the type $J_m = J_n + 1 = J_l - 1$.

4. New possibilities for spectroscopic investigations

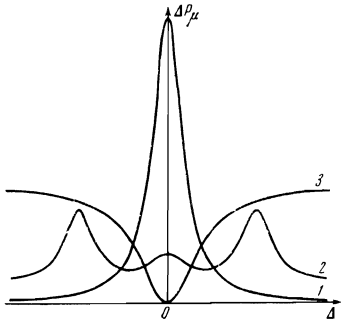


FIG. 2. Dependence of the nonlinear addition to the absorption coefficient for the weak field $\Delta P_\mu$ on the magnitude of the Zeeman splitting Δ; curve 1-waves circularly polarized in opposite directions acting on the transition $J_n = 0, J_m = 1, J_l = 2$ ($\epsilon = 0$); 2—the waves have the same linear polarizations and act on the same transition ($\epsilon \neq 0$); 3—waves having the same polarizations act on the transition $J_n = 0, J_m = J_l = 1$) ($\epsilon = 0$).

arise when the gas system is placed in a longitudinal magnetic field (H ↑↑ k). We consider them here on the basis of an analysis of the structure of an addition to $P_\mu$ nonlinear in energy, having in mind that it is relatively easy to separate out this addition by experimental means. For example, in order to separate it out it is possible to use the method utilized in [13], where polarization effects were investigated in the context of resonance fluorescence.

In accordance with (6) in the spontaneous approximation the expression for the nonlinear addition has the form

$$\Delta P_\mu \sim \mathrm{Re} \sum_{\alpha=\pm1} \left[ \frac{A_0 |G_\alpha{}^\mu|^2 |G_\alpha|^2}{\Gamma_{mm}(\Gamma_{nl} - i\varepsilon)} + \frac{A_1 |G_\alpha{}^\mu|^2 |G_{-\alpha}|^2}{\Gamma_{mm}(\Gamma_{nl} - i\varepsilon - i2\alpha\Delta)} + \frac{A_2 G_\alpha G_{-\alpha}{}^* G_{-\alpha}{}^\mu G_\alpha{}^{\mu *}}{(\Gamma_{mm} - i2\alpha\Delta)(\Gamma_{nl} - i\varepsilon)} \right]. \tag{13}$$

From a practical point of view it is convenient to investigate the dependence of $\Delta P_\mu$ on the magnitude of the Zeeman splitting Δ. This is associated with the fact that the method based on a variation of H (the "magnetic scanning" method) because of its simplicity and a considerably smaller number of errors possesses a number of indubitable advantages compared to other nonlinear-spectroscopic methods.

The nonlinear addition $\Delta P_\mu$ as a function of 2Δ is a superposition of Lorentz contours of widths $\Gamma_{nl}$ and $\Gamma_{mm}$. The position of these contours along the Δ scale depends on the magnitude of the detuning $\epsilon = \Omega_\mu - \Omega$, while their amplitudes depend on the polarizations of the waves and on the type of transition. Thus, if the waves are circularly polarized in opposite directions, then we have one Lorentzian $L(\Gamma_{nl})$ of width $\Gamma_{nl}$ and centered at $\Delta = -\epsilon/2\alpha$ (curve 1 in Fig. 2). In the case of identical linear polarizations there appear three split contours: two $L(\Gamma_{nl})$ centered at $\Delta = \pm\epsilon/2$ and one $L(\Gamma_{mm})$ centered on zero magnetic field (curve 2, Fig. 2). This splitting makes possible an individual investigation of the contours with the aim of determining their widths. In some cases, for example, for transitions of the type $J_n = 0$, $J_m = J_l = 1$ and identical linear polarizations of the strong and the weak waves, only the contour $L(\Gamma_{mm})$ remains (curve 3 in Fig. 3). In this unique situation the resonance characteristic features of the absorption spectrum of the weak field are manifested only in the form typical for the Hanle effect [14-15], and it can be utilized for a very precise measurement of the width of the common level m. If the waves are travelling in opposite directions (k ↓↑ k), then the nonlinear addition

$$\Delta P_\mu \sim \mathrm{Re} \sum_{\alpha=\pm1} \left[ -\frac{A_0 |G_\alpha{}^\mu|^2 |G_\alpha|^2 / \Gamma_{mm}}{\Gamma_{mn} + \Gamma_{ml} - i(\Omega + \Omega_\mu) - i2\alpha\Delta} + \frac{A_1 |G_\alpha{}^\mu|^2 |G_{-\alpha}|^2 / \Gamma_{mm}}{\Gamma_{mn} + \Gamma_{ml} - i(\Omega + \Omega_\mu)} + \frac{A_2 G_\alpha G_{-\alpha}{}^* G_{-\alpha}{}^\mu G_\alpha{}^{\mu *}}{\Gamma_{mn} + \Gamma_{ml} - i(\Omega + \Omega_\mu) - i2\alpha\Delta} \frac{1}{\Gamma_{mm} - i2\alpha\Delta} \right] \tag{14}$$

is a combination of the Lorentzians $L(\Gamma_{mm})$ and $L(\Gamma_{ml} + \Gamma_{mn})$, whose relative contribution to the total contour can also vary depending on the polarization of the waves. When $\Omega + \Omega_\mu = 0$ these Lorentzians have the same center $\Delta = 0$. Since always $\Gamma_{mm} < \Gamma_{mn} + \Gamma_{ml}$, then the contours in this case can be separated according to their widths: $L(\Gamma_{mm})$ will stand out against the background of the broader contour $L(\Gamma_{mm} + \Gamma_{ml})$.

5. We consider the effect of depolarizing collisions on the shape of the absorption spectrum for a weak field only for waves travelling in the same direction and for a longitudinal magnetic field when

$$\Delta P_\mu \sim \mathrm{Re} \sum_{\substack{\alpha=\pm1 \\ \varkappa=0,1,2}} \Delta N \left[ \frac{|G_\alpha{}^\mu|^2 |G_\alpha|^2}{\Gamma_{mn1} + \Gamma_{ml1} - i\varepsilon} p_\varkappa \left( \frac{\beta_{nl\varkappa}}{\Gamma_{nl\varkappa} - i\varepsilon} + \frac{\beta_{m\varkappa}}{\Gamma_{mm\varkappa}} \right) + \frac{|G_\alpha{}^\mu|^2 |G_{-\alpha}|^2}{\Gamma_{mn1} + \Gamma_{ml1} - i\varepsilon - i2\alpha\Delta} \left( \frac{6\beta_{nl2}}{\Gamma_{nl2} - i\varepsilon - i2\alpha\Delta} + \frac{p_\varkappa' \beta_{m\varkappa}}{\Gamma_{mm\varkappa}} \right) + \frac{G_\alpha G_{-\alpha}{}^* G_{-\alpha}{}^\mu G_\alpha{}^{\mu *}}{\Gamma_{mn1} + \Gamma_{ml1} - i\varepsilon - i2\alpha\Delta} \left( \frac{6\beta_{m2}}{\Gamma_{mm2} - i2\alpha\Delta} + \frac{p_\varkappa' \beta_{nl\varkappa}}{\Gamma_{nl\varkappa} - i\varepsilon} \right) \right]; \tag{15}$$

$$p_0 = p_0' = 2, \quad p_1 = -p_1' = 3, \quad p_2 = p_2' = 1.$$

Moreover, in a detailed analysis of spectral characteristic features we restrict ourselves to a transition of the type $J_n = 0$, $J_m = J_l = 1$. This simple example will enable us to elucidate certain specific properties of depolarizing collisions without recourse to awkward calculations.

It was shown above that in the spontaneous approximation the nonlinear addition $\Delta P_\mu$ for this transition was equal to zero if the waves were circularly polarized in opposite directions and were travelling in the same direction. This was explained by the fact that the weak wave interacted with transitions between Zeeman sublevels unperturbed by the strong field. Depolarizing collisions mix atomic states corresponding to different sublevels of the same level and therefore lead to a transfer of the perturbation produced by the strong field. A narrow dip of purely collisional origin appears on the Doppler contour of the absorption line of the weak wave. In this case the nonlinear addition

$$\Delta P_\mu \sim \frac{2}{9} \Delta N \tau_{2m} \frac{|G_\mu|^2 |G|^2 (\Gamma_{mn1} + \Gamma_{ml1})}{(\Gamma_{ml1} + \Gamma_{mn1})^2 + \varepsilon^2}, \tag{16}$$

responsible for the appearance of the dip has a form different from zero. The time

$$\tau_{2m} = \frac{1}{\Gamma_{mm}} \left( \frac{3}{4} \frac{\gamma_{mm1}}{\Gamma_{mm} + \gamma_{mm1}} + \frac{1}{4} \frac{\gamma_{mm2}}{\Gamma_{mm} + \gamma_{mm2}} \right), \tag{17}$$

to which it is proportional may, evidently, be interpreted as the lifetime of an atom in the level m after the first collision [7]. In obtaining (17) we consider for the sake of simplicity that the collisions are elastic so that $\gamma_{mm0} = 0$ and the total lifetime of the atom in this level $\tau_m = 1/\Gamma_{mm}$ is determined by the spontaneous decay. The collision broadening of the level m is characterized by the constants $\gamma_{mm\kappa}$ ($\kappa = 1, 2$) appearing in relation (17) and in the general formula (15): $\gamma_{mm1}$ is responsible for the decay of the magnetic dipole moment of the level m, while $\gamma_{mm2}$ is responsible for the decay of the electric quadrupole moment [14]. From (17) it follows that the probabilities of a change in the magnetic dipole moment and in the electric quadrupole moment in the process of collision are different. The ratio of these probabilities is

$$\frac{W_{\mathrm{mm}}}{W_{\mathrm{eqm}}} = \frac{3\gamma_{mm1}(\Gamma_{mm} + \gamma_{mm2})}{(\Gamma_{mm} + \gamma_{mm1})\gamma_{mm2}}. \tag{18}$$

As can be seen from (7), for $J_n = 0$, $J_m = J_l = 1$ the coefficients $\beta_{nl0} = \beta_{nl2} = 0$, while $\beta_{nl1} \neq 0$. Therefore in this case the collision broadening of the forbidden transition $n \to l$ is determined only by the constant $\gamma_{nl1}$. In the general case this is not so. The situation can change depending on the values of the total angular momenta of the levels. Thus, in transitions of the type $J_n - J_l = \pm 2$ Lorentz contours of widths $\Gamma_{nl0}$ and $\Gamma_{nl1}$ ($\beta_{nl0} = \beta_{nl1} = 0$) are completely absent in the nonlinear addition independently of the kind of polarization of the waves, while in the case of transitions $J_n - J_l = \pm 1$ ($J_m > 1$) the Lorentzian $L(\Gamma_{nl0})$ is absent since $\beta_{nl0} = 0$. In transitions of the type $J_n - J_l = 0$ all three Lorentz contours $L(\Gamma_{nl\kappa})$ are present whose collision widths $\gamma_{nl\kappa}$ ($\kappa = 0, 1, 2$), as was shown in [6], can be different.

In Sec. 4 it was pointed out that for waves with the same linear polarizations acting on transitions $J_n = 0$, $J_m = J_l = 1$ the dependence of the nonlinear addition $\Delta P_\mu$ on the intensity of the longitudinal magnetic field had the form characteristic of the Hanle effect. At sufficiently low pressures ($|\Gamma_{mn1} + \Gamma_{ml1} - \Gamma_{nl1} - \Gamma_{mm1,2}| \ll \Gamma_{mn1} + \Gamma_{ml1}$) this assertion remains valid also in the situation when the collision broadening is comparable to the spontaneous broadening. In this case the nonlinear addition

$$\Delta P_\mu(\Delta, \varepsilon=0) \sim \frac{\Delta N}{6}|G_\mu|^2|G|^2\left[\frac{1}{\Gamma_{mm1}(\Gamma_{mn1}+\Gamma_{ml1}-\Gamma_{mm1})} - \frac{\Gamma_{mm2}}{\Gamma_{mm2}^2+4\Delta^2}\frac{1}{\Gamma_{mn1}+\Gamma_{ml1}-\Gamma_{mm2}}\right] \quad (19)$$

as a function of $2\Delta$ contains the Lorentzian $L(\Gamma_{mm2})$ the collision broadening for which is $\gamma_{mm2}$. The latter circumstance also serves as an indication of the general nature of the Hanle effect and of the resonance feature of (18). At higher pressures and for other transitions when $\Delta P_\mu(\Delta)$ contains several Lorentz contours $L(\Gamma_{mm2})$ is distinguished by the fact that of all the contours only this one is centered on zero magnetic field if the detuning is given by $\epsilon \neq 0$.

In conclusion it should be noted that although the results of this paper are formulated by us only for resonance scattering through the upper level, by proceeding in analogy to Rautian and Feoktistov [16] these results can be easily applied to the study of polarization and collision effects in the case of scattering through the lower level, and also in the case of two-quanta luminescence and two-quanta absorption.

[1)]The waves differ in intensity. The weak wave in contrast to the strong one does not produce saturation effects.


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Translated by G. Volkoff
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