Anomalous self-rotation of the polarization plane of a confined light beam in a gas located in a weak magnetic field

M. V. Krasheninnikov, V. S. Smirnov, M. V. Sultanov, A. M. Tumaĭkin, and V. I. Yudin

Siberian State Research Institute for Metrology

(Submitted 23 April 1987)
Zh. Eksp. Teor. Fiz. **91**, 203–209 (January 1988)

We consider self-rotation of the polarization plane of a spatially confined light beam in a gas of atoms initially in an unpolarized ground state, in the presence of an arbitrarily oriented constant magnetic field whose strength is substantially less than that of the weak fields that cause polarization rotation via the known Faraday, Hanle, and other effects. A linearly polarized field produces in the ground-state atoms a quadrupole moment whose orientation is changed by the magnetic field. In atoms having a ground-state electron angular momentum less than unity (e.g., alkali metals) the effect is possible only if account is taken of hyperfine splitting.

In the absence of a magnetic field, one of the mechanisms of self-rotation of the polarization in a nonlinear isotropic medium is formation of a magnetic moment (orientation) in a system of atoms via the circular component of the electromagnetic field.[1,2] This mechanism does not work in a linearly polarized field, and polarization rotation is possible either if the atom acquires a quadrupole moment by spontaneous emission[3,4] or as a result of parity nonconservation.[5]

Polarization-plane rotation in a magnetic field is due to the Faraday and Hanle effects. For Faraday rotation to occur in a gas, the Larmor frequency Ω must be of the order of the Doppler line width $k\bar{v}$ (rarefied gas, homogeneous broadening γ purely radiative, $\gamma \ll k\bar{v}$). We consider in the present paper interaction of linearly polarized resonant light with a gas of atoms that are initially in an unpolarized ground state, in the presence of an arbitrarily directed magnetic field. In the limiting case $\Omega \ll \gamma$, when the Faraday and Hanle effects can be neglected in the excited state, the magnetic field splits the Zeeman sublevels of only the ground state. For the Hanle effect to occur in the ground state of a rarefied gas it is necessary to make the Larmor frequency Ω comparable with the characteristic relaxation rate of the ground state, viz., with the rate $\gamma_\varkappa^0$ of the depolarizing collisions of the multipole moments of rank $\varkappa$, and with the reciprocal time $\bar{t}^{-1} = \bar{v}/r_0$ of interaction of the atom with the light beam (r_0 is the transverse beam dimension). Under real conditions the rate of depolarizing collisions in the ground state can be $\gamma_\varkappa^0 \ll \bar{t}^{-1}$. Indeed, in instantaneous quenching of the polarization moments by the cell wall the upper limit of $\gamma_\varkappa^0$ is determined by the reciprocal of the time to negotiate that transverse dimension R of the cell. If $R \gg r_0$ in that case, we have $\gamma_\varkappa^0 \bar{t} \ll t$. The increment added to $\gamma_\varkappa^0$ by the depolarizing collisions with the buffer gas is vanishingly small, $\gamma_\varkappa^0/\gamma \sim 10^{-8}$–$10^{-10}$ (Ref. 1). Moreover, collisions with the buffer gas suppress the depolarization on the walls, since the effective time of travel to the wall is determined in this case by the diffusion time $R^2/2D$ (D is the diffusion coefficient). We consider hereafter a one-component low-density gas, when all the collisions can be neglected. The characteristic scale of Ω is then $\bar{t}^{-1}$. Consequently, the region of strong and weak magnetic fields is now determined by the parameter $\Omega\bar{t}$ at $\gamma\bar{t} \gg 1$ and $\gamma_\varkappa^0 \ll \Omega \ll \gamma$. For alkali metals, for example, we have $\Omega\bar{t} \sim 1$ ($r_0 \sim 1$ cm) for fields $H \sim 10^{-2}$–10^{-3} G, i.e., the strong-field region begins with $H \sim 10^{-1}$–10^{-2} G and magnetic-field saturation is reached in fields traditionally regarded as weak.

One more feature of the situation considered is optical orientation of the atoms in the ground state. It is known[3,4] that when dealing with laser orientation it is always possible to choose a coherent field in which there is no usual saturation [$G \ll 1$, where G is the saturation parameter, Eq. (8)]. The following situations can then be realized: $\gamma G\bar{t} \gg 1$, a case investigated in detail in Ref. 6, $G < \gamma_\varkappa^0/\gamma$ (Ref. 7), and finally the case

$$\gamma_\varkappa{}^0\bar{t} \ll \gamma G\bar{t} < 1. \tag{1}$$

The investigation of this intermediate regime is the subject of the present paper.

1. FORMULATION OF PROBLEM

Consider a gas of atoms in a cell of length l and of low optical density $\chi l < 1$ (χ is the absorption coefficient). Let a linearly polarized light beam with Gaussian profile of width r_0 be incident on the cell boundary $y = 0$. We neglect reflection from the cell boundaries and the spreading of the beam in the medium. The field inside the cell can then be regarded as a plane wave propagating along $\mathbf{e}_2$ (the y axis), having a specified transverse distribution, and usually elliptically polarized with the polarization ellipse rotated through an angle ψ relative to the initial coordinate frame:

$$E = \boldsymbol{\mu} E(r, y) \exp\{-i(\omega t - ky)\} + \text{c.c.}, \tag{2}$$

where $E(r,y)$ is the slowly varying amplitude and $\boldsymbol{\mu}$ the elliptic unit vector of the polarization:

$$\boldsymbol{\mu} = \{(\mathbf{e}_3 \cos\psi + \mathbf{e}_1 \sin\psi) + i \operatorname{tg}\alpha(\mathbf{e}_1 \cos\psi - \mathbf{e}_3 \sin\psi)\}(1+\operatorname{tg}^2\alpha)^{-1/2}. \tag{3}$$

The angle α is called the ellipticity angle ($-\pi/4 \leqslant \alpha \leqslant \pi/4$); $\tan\alpha$ is equal to the ratio of the ellipse axes, and the sign of α determines the direction of the rotation along the ellipse.

The problem is to find on the boundary $y = l$ the field polarization orthogonal to the initial one, i.e., to calculate the angles ψ and α. We assume the field at the entry $y = 0$ into the medium to be linearly polarized in the direction of $\mathbf{e}_3$ (the z axis). The orientation of the external magnetic field $\mathbf{h}$ is characterized by spherical angles (φ,θ) in a coordinate frame $(\mathbf{e}_1,\mathbf{e}_2,\mathbf{e}_3)$.

 0038-5646/88/010113-04$04.00

If the condition (1) is met, the excited level serves as the virtual level from which is emitted spontaneous radiation that redistributes the populations of the magnetic sublevels of the ground state. An excited level does not contribute to the polarization of the medium. Consequently, the complete set of equations for the density matrix, a set describing the resonant gas, reduces to a set of equations for the density matrix of the ground state in which we take the hyperfine splitting into account. In a magnetic field arbitrarily oriented relative to the beam it is more convenient to expand the ground-state density matrix in terms of the polarization operators $\rho_\varkappa^{(n)}$ (the $\varkappa$, q representation[8]), where the index numbers the sublevels of the hyperfine structure of the ground state:

$$\left(\frac{\partial}{\partial t}+\mathbf{v}\nabla\right)\rho_\varkappa^{(n)}-i\Omega g_n[\varkappa(\varkappa+1)]^{1/2}\{h_1\otimes\rho_\varkappa^{(n)}\}_\varkappa$$
$$=\gamma G\sum_{\varkappa_1\varkappa_2}(A_{\varkappa_1\varkappa_2}^{(1)nn_1}-A_{\varkappa_1\varkappa_2}^{(2)nn_1})(-1)^{\varkappa_1+\varkappa_2}3[(2\varkappa_1+1)(2\varkappa_2+1)]^{1/2}$$
$$\times\{\{\mu_1^*\otimes\mu_1\}_{\varkappa_1}\otimes\rho_{\varkappa_2}^{(n_1)}\}_\varkappa, \tag{4}$$

$$A_{\varkappa_1\varkappa_2}^{(1)nn_1}=$$
$$\sum_m(-1)^{F_n+F_m}2(2J_1+1)$$
$$\times\frac{C_{mn}C_{mn_1}}{1+\Delta_{mn_1}^2}\begin{Bmatrix}\varkappa & F_n & F_n\\ 1 & F_m & F_m\end{Bmatrix}\begin{Bmatrix}1 & 1 & \varkappa_1\\ F_m & F_m & \varkappa\\ F_{n_1} & F_{n_1} & \varkappa_2\end{Bmatrix}, \tag{5}$$

$$A_{\varkappa_1\varkappa_2}^{(2)nn_1}=\delta_{nn_1}\sum_m(-1)^{F_n-F_m}C_{mn_1}\left[\frac{1}{1-i\Delta_{mn}}-\frac{(-1)^{\varkappa+\varkappa_1+\varkappa_2}}{1+i\Delta_{mn}}\right]$$
$$\times\begin{Bmatrix}F_m & F_n & 1\\ \varkappa_1 & 1 & F_n\end{Bmatrix}\begin{Bmatrix}F_n & \varkappa_2 & F_n\\ \varkappa_1 & F_n & \varkappa\end{Bmatrix}. \tag{6}$$

We use here the notation of Ref. 9 for the irreducible tensor products and the $3jn$ symbols. The summation in (5) and (6) is over the hyperfine structure of the excited state, labeled by the subscript m. The constants C_{mn} in (5) and (6) and g_n in (4) yield the dependences of the dipole moment d and of the ground-state g-factors on the total moments $F_{(n,m)}$ of the nuclear spin I and of the electron moments J_0 and J_1 (0 and 1 denote the lower and upper states):

$$C_{mn}=(2F_m+1)(2F_n+1)\begin{Bmatrix}I & J_0 & F_n\\ 1 & F_m & J_1\end{Bmatrix}^2,$$
$$g_n=(-1)^{F_n+J_0+I+1}g_0(2F_n+1)\begin{Bmatrix}I & J_0 & F_n\\ 1 & F_n & J_0\end{Bmatrix}. \tag{7}$$

The saturation parameter G and the spectral components Δ_{mn} are defined, with allowance for the hyperfine splitting and the Doppler shift, as follows:

$$G=|Ed/\hbar\gamma|^2,\quad \Delta_{mn}=(\omega-\omega_{mn}-\mathbf{kv})/\gamma. \tag{8}$$

The terms $A^{(1)}$ and $A^{(2)}$ in (4)–(6) describe respectively the arrival in the ground state via spontaneous emission and the departure from the ground state on account of the external field. We neglect the collisions in both the ground and the excited states.

The slowly varying field components E_j in the medium are described by simplified Maxwell equations in which the diffraction spreading of the initial packet can be neglected,

$$\partial E_j/\partial y=2\pi ikP_j, \tag{9}$$

with a polarization P_j that depends on the polarization moments of the ground state:

$$P_j=\frac{3i|d|N}{\hbar\gamma}\left\langle\sum_{\substack{mn\\ \varkappa}}\frac{(-1)^{F_n+F_m}C_{mn}}{1-i\Delta_{mn}}\begin{Bmatrix}F_m & F_n & 1\\ \varkappa & 1 & F_n\end{Bmatrix}\right.$$
$$\left.\times(\{\mathbf{e}_j\otimes\mu\}_\varkappa\rho_\varkappa^{(n)})\right\rangle_v E. \tag{10}$$

The symbol $\langle\ldots\rangle_v$ denotes averaging over the velocities with a Maxwellian distribution, and N denotes the atom density. We emphasize that the components P_j are expressed in invariant form, just as Eq. (4).

The set of equations that describes the interaction of the field with the resonant gas is now closed. The boundary conditions for the field were formulated above. We specify the initial conditions for $\rho_\varkappa^{(n)}$ by starting from an adiabatic start of the interaction:

$$\rho_\varkappa^{(n)}=\delta_{\varkappa,0}N^{(n)}(2F_n+1)^{-1/2},\quad \sum_n N^{(n)}=1, \tag{11}$$

where $N^{(n)}$ are the equilibrium occupation numbers of the hyperfine structure of the ground state. Bearing in mind that $l\gg r_0$, the boundary effects for $\rho_\varkappa^{(n)}$ can be disregarded.

2. POLARIZATION MOMENTS OF THE GROUND STATE

To set apart in explicit form the parameter for the expansion of the solution of Eq. (4), we transform this equation into an integral one. We denote the entire right-hand side of (4) by $\gamma\Phi_\varkappa f(r)$, where the dimensionless function $f(r)$ specifies the transverse distribution of the field intensity. We rotate the coordinate system in (4) so that the quantization axis coincides with the direction of the magnetic field $\mathbf{h}$. Integrating next (4) with the specified right-hand side and rotating back the coordinate frame, we obtain the set of integral equations of interest to us:

$$\rho_\varkappa^{(n)}=\delta_{\varkappa,0}N^{(n)}(2F_n+1)^{-1/2}+\gamma\int_0^\infty dt f(\mathbf{r}-\mathbf{v}t)$$
$$\times\sum_L\{\Phi_\varkappa\otimes n_L\}_\varkappa\sum_q(-1)^{\varkappa-q}C_{\varkappa-q,\varkappa q}^{L0}\exp\{-it\Omega qg_n\}. \tag{12}$$

Here $C_{\varkappa_1q_1,\varkappa_2q_2}^{L\lambda}$ are the Clebsch-Gordan coefficients resulting from the two rotations of the coordinate frame, and $n_L(\theta,\varphi)$ are unit spherical harmonics. The tensor $\Phi_\varkappa=\Phi_\varkappa(\mathbf{r}-\mathbf{v}t)$ depends on the coordinates via $\rho_{\varkappa_2}^{(n)}$ from the right-hand side of (4).

A perturbation theory suitable for the considered situation can be constructed by iteration in the system (12). We confine ourselves to the first iteration. To this end we must substitute in (4) [or in (12)] the equilibrium distribution (11), as a result of which $\Phi_\varkappa$ ceases to depend on the coordinates, and the "interaction time" of the atom with the light beam is set apart:

$$t_{\varkappa,L}^{(n)}(\mathbf{r},\mathbf{v})=\sum_{q}(-1)^{\varkappa-q}C_{\varkappa-q,\varkappa q}^{L0}\int_{0}^{\infty}dt f(\mathbf{r}-\mathbf{v}t)\exp\{-it\Omega q g_n\}, \tag{13}$$

$t_{\varkappa,L}^{(n)}(\mathbf{r},\mathbf{v})=\delta_{L,0}(2\varkappa+1)^{1/2}t(\mathbf{r},\mathbf{v})$ determines the time of interaction of the atom with a bounded light beam as a function of its coordinate and velocity in the absence of a magnetic field.

As seen from (12) and (4), the multipole moments induced by the field at the atoms in the ground state are proportional to tensor products of the type $\{\{\boldsymbol{\mu}^*\otimes\boldsymbol{\mu}\}_\varkappa\otimes n_L\}_\varkappa$, where the rank $\varkappa$ is set by the field polarization $\{\boldsymbol{\mu}^*\otimes\boldsymbol{\mu}\}_\varkappa$ and the direction n_L by the magnetic field. There is no magnetic moment in a linearly polarized optical field $\{\boldsymbol{\mu}\otimes\boldsymbol{\mu}\}_{\varkappa=1}=0$.

The expansion parameter in (12) is the quantity $\gamma\bar{G}\bar{t}$. Accurate to the first power of this parameter, it is necessary to substitute in (12) the polarization of the initial field, i.e., $\mu=\mathbf{e}_3$ and, if $\chi l<1$, also the initial value of the field amplitude $E(r,0)$. In this limit there is induced in the ground state only a quadrupole moment ($\varkappa=2$) with components proportional to $n_{L\lambda}(\theta,\varphi)$ ($|\lambda|\leqslant 2$). This can be easily verified by orienting the quantization axis along the vector $\mathbf{e}_3$: $\rho_{\varkappa q}^{(n)}\sim\delta_{q\lambda}n_{L\lambda}(\theta,\varphi)$. It is precisely by virtue of the fact that the quadrupole moment of the medium has now components oriented by the magnetic field ($\lambda\neq 0$), the polarization (10) of the medium acquires a component along the axis $\mathbf{e}_1$ and perpendicular to the initial $\mathbf{e}_2$. In other words, the distinctive feature of the considered problem is that the polarization-plane rotation is now due to the quadrupole and not the magnetic moment.

3. ROTATION ANGLE AND DEGREE OF ELLIPTICITY

Integrating Eq. (9) along the $\mathbf{e}_3$ axis and taking (2) and (3) into account, we obtain the rotation angle ψ and the ellipticity angle α at the exit from the cell:

$$\psi=-\chi l\,\mathrm{Re}\,\xi_{13},\quad \alpha=-\chi l\,\mathrm{Im}\,\xi_{13}. \tag{14}$$

Here ξ_{12} are the off-diagonal components of the dielectric susceptibility tensor:

$$\xi_{13}=\gamma G\bar{t}\sum_{L,n}[6L(L+1)]^{-1/2}C_{21,20}^{L1}P_L^{(1)}(\theta)$$
$$\times\{[1+(-1)^L]\tau_{L,c}^{(n)}\cos\varphi$$
$$-[1-(-1)^L]\tau_{L,s}^{(n)}\sin\varphi\}B_n, \tag{15}$$

$$B_n=\sum_{\substack{n_1\\ m,m_1}}\frac{3(2J_0+1)N^{(n_1)}}{2F_{n_1}+1}\frac{C_{mn}C_{m_1n_1}\begin{Bmatrix}F_m & 1 & F_n\\ 2 & F_n & 1\end{Bmatrix}}{1-i(\omega_{nn_1}-\omega_{mm_1})/2\gamma}$$
$$\times\Big[(-1)^{-F_m+F_{m_1}}\delta_{nn_1}\begin{Bmatrix}F_{m_1} & 1 & F_n\\ 2 & F_n & 1\end{Bmatrix}+(-1)^{F_m+F_{n_1}}(2J_1+1)C_{m_1n}$$
$$\times\begin{Bmatrix}F_{n_1} & 1 & F_{m_1}\\ 2 & F_{m_1} & 1\end{Bmatrix}\begin{Bmatrix}1 & F_{m_1} & F_n\\ 2 & F_n & F_{m_1}\end{Bmatrix}\Big], \tag{16}$$

and ω_{nn_1} and ω_{mm_1} are the hyperfine splitting frequencies of the ground and excited states. It was assumed for simplicity in the derivation of (16) that the hyperfine splitting is overlapped by the Doppler broadening, $|\omega_{ll_1}|\ll k\bar{v}$. In this case $N^{(n)}=1/\min[(2J_0+1);(2L+1)]$.

The entire dependence on the magnetic-field orientation is set apart in (15) in explicit form [$P_L^{(1)}(\theta)$ is the associated Legendre polynomial, $L\leqslant 4$]. The dependence on the magnetic field strength is contained in the mean dimensionless times τ of the interaction between the atom and the field ($\Omega_n=\Omega_{gn}$) [see Eq. (13)]:

$$\tau_{L,c}^{(n)}=\sum_{q}(-1)^q C_{2-q,2q}^{L0}\int_0^\infty\frac{dt}{\bar{t}}\cos q\Omega_n t\langle f(\mathbf{r}-\mathbf{v}t)\rangle_{\mathbf{v}},$$

$$\tau_{L,s}^{(n)}=\sum_{q}(-1)^q C_{2-q,2q}^{L0}\int_0^\infty\frac{dt}{\bar{t}}\sin q\,\Omega_n t\langle f(\mathbf{r}-\mathbf{v}t)\rangle_{\mathbf{v}}. \tag{17}$$

The typical values of the magnetic fields are easiest to discern for a Gaussian beam $f(r)=\exp(-\rho^2)$, $\rho=r/r_0$. In this case the problem has only one parameter $\varepsilon_n=\Omega_n t$.

Let us consider the limiting cases. We denote the integrals in (17) by τ_c^q and τ_s^q. In a weak magnetic field $|\varepsilon_n|\ll 1$ and

$$\tau_s^q=-q\varepsilon_n\left(\ln|q\varepsilon_n|+C-1+\frac{1}{2}\int_0^{\rho^2}\frac{1-e^{-x}}{x}dx\right),$$

$$\tau_c^q=-\frac{\pi}{2}|q\varepsilon_n|+\tau_0,\quad \tau_0=\frac{\pi}{2}e^{-\rho^2/2}I_0(\rho^2/2), \tag{18}$$

$I_0(x)$ is a modified Bessel function and C is the Euler constant. It can be seen that in the limit considered there is no ε_n section linear in the magnetic field. It appears in much weaker fields $\Omega<\chi\bar{v}$ ($\chi l<1$) for example, at $H<10^{-4}$ G for $\chi^{-1}\sim l\sim 10^2$ cm, and $\bar{v}\sim 10^5$ cm/s.

It is seen from (15) that there is no rotation if the magnetic field is oriented along the electric-field vector $\mathbf{e}_3$ or along the vector $\mathbf{e}_2$. The dependences of the time $\tau(\rho)$ on the transverse coordinates is a consequence of the polarization of the atoms in the ground state (infinite lifetime).

In the opposite limit $|\varepsilon_n|\gg 1$ the oscillations in (17) lead to rapid saturation, and as a result

$$\tau_c^q=\delta_{q0}\tau_0,\quad \tau_s^q=(q\varepsilon_n)^{-1}e^{-\rho^2}. \tag{19}$$

The principal correction to (15) does not depend on the magnetic field, its dependence on the field direction remains and has a specific form: L is even, by virtue of the selection rules, and the angular dependence is

$$\xi_{13}(\theta,\varphi)\propto\cos\varphi\sin 2\theta(1-3\cos^2\theta).$$

It must be emphasized that under ordinary conditions $|\varepsilon_n|\gg 1$ for the earth's magnetic field. Therefore, if the effect of the earth's field cannot be controlled, a circular component can always appear in an initially linearly polarized field. As a result, the medium itself can acquire a magnetic moment in the next order of the light intensity.

The ellipticity angle α, in contrast to the rotation angle ψ, is determined by the imaginary part of the spectral constant B_n. If the hyperfine splitting in the ground (ω_{nn_1}) and excited (ω_{mm_1}) states is large compared with the radiative broadening, it follows from (16) that $\alpha/\psi\sim\gamma/\omega_{ll_1}\ll 1$, i.e., the degree of ellipticity is low compared with the rotation angle. In the opposite limiting case $\omega_{ll_1}\ll\gamma$ it is possible to

carry out the summation in (16), as a result of which the constant B_n is proportional to the $3j$ symbol

$$\begin{Bmatrix} J_1 & J_0 & 1 \\ 2 & 1 & J_0 \end{Bmatrix},$$

from which it follows that, in view of the selection rules, rotation is possible only if $J_0 \geqslant 1$. For alkali metals ($J_0 = 1/2$), for example, the hyperfine splitting must be taken into account. Moreover, as seen from (16), neglect of the hyperfine splitting only in the excited state ($\omega_{mm_1} \ll \gamma$) also leads to the absence of rotation at $J_0 < 1$.

The physical causes of the anomalous rotation of the polarization plane are quite obvious. A linearly polarized field produces a quadrupole moment in atoms in the ground state, and a strong magnetic field ($\Omega\bar{t} \sim 1$) reorients the field in such a way that it acquires components (1,3) or (x,z). The latter leads to elliptic birefringence, which is in fact described by Eqs. (14)–(16). Alkali metals have $J_0 = 1/2$, therefore their quadrupole moment is determined entirely by the hyperfine splitting and vanishes if $\omega_{II_1} \ll \gamma$. In atoms with $J_0 \geqslant 1$ the effect is preserved also without hyperfine splitting.

In conclusion, we estimate the rotation angle ψ for laser radiation frequency resonant to the transition $F_n = 2 \to F_m = 1.2$ in the D_1 line of sodium. Since the Doppler broadening kv is small compared with the hyperfine splitting of the ground state (1740 MHz), the contribution of the transitions $F_n \to F_m = 1.2$ to B_n can be disregarded. Summing in (15) and (16) we have then for strong magnetic field (19) at the center of the beam ($\rho = 0$)

$$\psi \approx 0{,}01 \chi l \gamma G \bar{t} \cos\varphi \sin 2\theta (1 - 3\cos^2\theta)$$

($I = 3/2$, $J_0 = J_1 = 1/2$, and the hyperfine splitting in the excited state is 190 MHz). The parameter $\gamma G\bar{t}$ can be made of the order of unity for the following values: saturation intensity in sodium D_1 line $I_{sat} \approx 64$ Mw/mm^2, $\gamma \approx 6 \cdot 10^7$ mW/mm^2, $r_0 \approx 1$ cm, and $v \approx 5 \cdot 10^4$ cm/s. The value of ψ at the optical density $\chi l \sim 1$ can reach then an order of 1°.

An experimental check on the predicted effect is of considerable interest. Note that the anomalous polarization-plane rotation recently observed in Ref. 10 was produced in strong magnetic fields ($\Omega \gtrsim \gamma$) outside the limits of the approximation $\gamma G\bar{t} < 1$ for an atom beam; the experimental conditions were therefore likewise outside the scope of the model considered in the present paper.

[1]W. Happer, Rev. Mod. Phys. **44**, 169 (1972).
[2]V. S. Smirnov and A. M. Tumaikin, Opt. Spektr. **45**, 442 (1978).
[3]A. P. Kazantsev, V. S. Smirnov, A. M. Tuaikin, and I. A. Yagofarov, *ibid.* **58**, 500 (1985).
[4]A. P. Kazantsev, V. S. Smirnov, G. I. Surdutovich, *et al.* J. Opt. Soc. Am. **B2**, 1731 (1985).
[5]I. B. Khriplovich, Parity Nonconservation in Atomic Phenomena [in Russian], Nauka, 1981, Chap. 1.
[6]A. P. Kazantsev, V. S. Smirnov, A. M. Tumaĭkin, and I. A. Yagoforov, Preprint No. 5, Inst. Atmosph. Optics, Siberian Div. USSR Acad. Sci., 1982.
[7]A. I. Okunevich, Opt. Spektrosk. **50**, 443 (1981).
[8]M. I. D'yakonov and V. I. Perel', Zh. Eksp. Teor. Fiz. **47**, 1483 (1964) [Sov. Phys. JETP **20**, 997 (1965)].
[9]D. A. Varshalovich, A. N. Moskalev, and V. K. Khersonskiĭ, Quantum Theory of Angular Momentum. Nauka, 1975, Chap. 5.
[10]O. E. Beuvich, I. S. Grigor'ev, A. F. Semerok, and V. A. Firsov, Pis'ma Zh. Eksp. Teor. Fiz. **45**, 322 (1987) [JETP Lett. **45**, 407 (1987)].

Translated by J. G. Adashko