EFFECT OF RESONANT-RADIATION CAPTURE ON THE CHARACTERISTICS OF A GAS LASER

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The nonlinear polarizability of an active gas medium is calculated by taking into account capture of resonant radiation. Radiation capture results in efficient mixing of the velocity distribution for atoms at the upper working level and also in leveling out of the population in the Zeeman sublevels of the atoms. It is shown that, along with formation of Bennett holes, generation leads to a general lowering of the amplification contour and the saturation effect acquires some properties which are characteristic of homogeneous broadening. It is shown that radiation capture alters the dependence of the generation intensity on resonator tuning (shape of Lamb dip). It is shown that the parameters characterizing the dip depend on the total moments of the operating levels and on polarization of the laser radiation.

INTRODUCTION

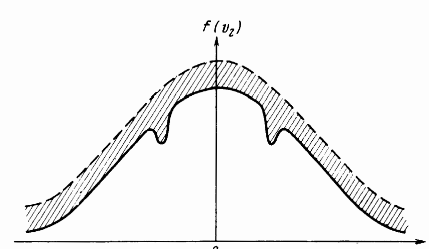
EFFECTS of pressure in gas lasers have been intensely investigated in recent times, both experimentally and theoretically (see the bibliographies in[1-3]). Usually the effects of pressure are ascribed to collisions of different types. The main interest is the influence of collisions on the form of the Lamb dip.

As to the dragging of resonant radiation, until recently it has not been investigated at all[1)]. Yet, if one of the working levels is connected with the ground state by an allowed optical transition, then the dragging at this transition is usually practically complete even at pressures when the collisions are not yet significant. Therefore, when considering collisions of atoms in a gas laser it is always necessary to take into account the fact that they occur against the background of almost complete dragging of the radiation. In the opposite case, the collision cross sections determined from the dependence of the shape of the hole on the pressure may turn out to be in error.

The role of dragging can be understood on the basis of the following qualitative considerations. We assume that the upper working level is coupled with the ground state. In the generation regime, the monochromatic field of the laser decreases the population of this level for those atoms for which the Doppler frequency shift coincides with the detuning of the resonator relative to the frequency of the working transition within the limits of the natural line width (the Bennet holes). The dragging of the resonant radiation leads to an effective mixing of the distribution with respect to the velocities for the atoms at the upper working level, and also equalizes the populations on the Zeeman sublevels of these atoms. As a result, the decrease of the population in the region of the Bennet holes brings about a decrease of the populations of the upper working level for atoms with all velocities (see the figure). Therefore, besides the formation of Bennet holes, generation causes a general lowering of the amplification contour. Thus, owing to the dragging, the effective saturation acquires singularities that are characteristic of homogeneous broadening.

It is easy to estimate the additional decrease of the number of atoms at the upper working level, as the result of dragging, i.e., the area S shown shaded in the figure.

We denote the total natural width of the upper level by γ_1; $\gamma_1 = \gamma_1' + \tilde{\gamma}_1$, where $\tilde{\gamma}_1$ is that part of the natural width which is connected with the transition to the ground state. Usually this transition makes the main contribution to γ_1, so that the quantity γ_1', which is connected with all the remaining transitions, is small: $\gamma_1' \ll \tilde{\gamma}_1$. If S_B is the area of the Bennet holes, then we can write the following balance equation: $\gamma_1' S = \tilde{\gamma}_1 S_B$. Indeed, owing to the dragging of the resonant radiation, the total number of atoms at the upper working level relaxes with a small decay constant γ_1', whereas the velocity mixing is determined by the large quantity $\tilde{\gamma}_1$. Consequently, $S \gg S_B$, but not all the excited atoms are



Distribution with respect to the velocities v_z of the atoms at the upper working level. Dashed line – in the absence of generation, solid line – in the presence of generation. Owing to the dragging of the resonant radiation, a mixing of the atoms by velocities takes place, and leads to a general lowering of the distribution function.

[1)]The significant consequences ensuing from dragging of resonant radiation in a laser were first discussed at the All-Union Symposium on Gas Laser Physics (Novosibirsk, July 1969) in papers by I. M. Beterov, Yu. A. Matyugin, S. G. Rautian, and V. P. Chebotaev, where convincing experimental data were also cited, and also in a paper by the present authors, where the present work was reported. The main conclusions of all these papers are in agreement.

important for generation, and only those whose velocities lie in the region of the hole. The additional decrease of the number of excited atoms in the region of the hole, due to the dragging, constitutes a small fraction of the value of S, on the order of γ_{10}/ku, where γ_{10} is the natural line width for the working transition, and ku is the Doppler width of this line. Thus, the role of the dragging of the radiation is characterized by the parameter $\tilde{\gamma}_1\gamma_{10}/\gamma_1' ku$.

The ratio γ_{10}/ku is small, and therefore the influence of the dragging on the generation is appreciable, provided only the transition to the ground state makes the overwhelming contribution to the natural width γ_1, i.e., if $\gamma_1' \ll \gamma_1$. This is a fairly common situation, and we shall henceforth assume it to obtain.

The role of dragging increases with increasing pressure, since the collisions cause the width of the Bennet hole to increase. The dragging of the radiation changes the dependence of the intensity of the generation on the tuning of the resonator (the shape of the Lamb dip). As a result of the overall settling of the amplification contour (which is proportional to the generation intensity), the Lamb dip becomes less pronounced.

In this paper we calculate the nonlinear polarizability of the active gas medium with allowance for dragging of the resonant radiation, and obtain an expression for the shape of the Lamb dip under these conditions. We show that the parameters characterizing the hole depend on the momenta j_1 and j_0 of the working levels and on the polarization of the laser radiation. These parameters are calculated for different cases.

ALLOWANCE FOR DRAGGING OF RADIATION IN THE EQUATIONS FOR THE DENSITY MATRIX

We shall assume that the pressure is so low that the collisions are still insignificant. Then the equations for the density matrix of the atoms, with allowance for the dragging of the resonant radiation, takes the form

$$\frac{df_{mm'}}{dt} + (\gamma_1' + \tilde{\gamma}_1) f_{mm'} - \tilde{\gamma}_1 (\hat{L}f)_{mm'}$$
$$= \frac{i}{\hbar} \sum_{\mu} [(\mathbf{E}\mathbf{d}_{m\mu})\psi_{\mu m'} - \psi_{m\mu}(\mathbf{E}\mathbf{d}_{\mu m'})] + \gamma_1' N_1 F(v)\delta_{mm'},$$
$$\frac{d\varphi_{\mu\mu'}}{dt} + \gamma_0\varphi_{\mu\mu'} = \frac{i}{\hbar}\sum_{m} [(\mathbf{E}\mathbf{d}_{\mu m})\psi_{m\mu'} - \psi_{\mu m}(\mathbf{E}\mathbf{d}_{m\mu'})]$$
$$+ \sum_{mm'} \Gamma_{\mu\mu'}^{mm'} f_{mm'} + \gamma_0 N_0 F(v)\delta_{\mu\mu'}, \tag{1}$$
$$\frac{d\psi_{\mu m}}{dt} + (\gamma_{10} - i\omega_0)\psi_{\mu m} = \frac{i}{\hbar}\Big[\sum_{m_1} (\mathbf{E}\mathbf{d}_{\mu m_1}) f_{m_1 m} - \sum_{\mu_1} \varphi_{\mu\mu_1}(\mathbf{E}\mathbf{d}_{\mu_1 m})\Big].$$

Here $d/dt = \partial/\partial t + \mathbf{v}\cdot\nabla$. The form of the equations and the notation is the same as in our earlier paper[4]. The term $-\tilde{\gamma}_1(\hat{L}f)_{mm'}$, added to the first equation of (1), takes into account the dragging of the resonant radiation,

$$(\hat{L}f)_{mm'} = \int d^3r' \int d^3v' \sum_{m_1 m_1'} K_{mm'}^{m_1 m_1'}(\mathbf{r}-\mathbf{r}', \mathbf{v}, \mathbf{v}') f_{m_1 m_1'}(\mathbf{r}', \mathbf{v}', t), \tag{2}$$

the kernel $K_{mm'}^{m_1 m_1'}$ was calculated by us in[5].

In Eqs. (1) and (2), $f_{mm'}$ denotes the elements of the density matrix pertaining to the upper working level (level 1), while $\varphi_{\mu\mu'}$ denotes the same for the lower working level (the zero level). The density matrix elements connecting the upper and lower states are denoted by $\psi_{\mu m}$ and $\psi_{m\mu}$; ω_0 is the frequency of the working transition, $\gamma_1 = \gamma_1' + \tilde{\gamma}_1$ and γ_0 are the reciprocal lifetimes of the upper and lower levels, $\gamma_{10} = (1/2)(\gamma_1 + \gamma_0)$, $\tilde{\gamma}_1$ is the part of the width of the upper working level due to the transition to the ground state.

The term containing $\Gamma_{\mu\mu'}^{mm'}$ describes the arrival of the atoms from the state 1 to the state 0 in spontaneous emission. In[4], the matrix Γ is expressed in terms of the probability γ of transitions from the level 1 to the level 0. The following relation holds:

$$\sum_{m} \Gamma_{\mu\mu'}^{mm} = \delta_{\mu\mu'} \frac{2j_1+1}{2j_0+1}\gamma.$$

We note that if spontaneous transitions from the level 1 are possible only to the level 0 and to the ground state, then γ and γ_1' coincide.

It is assumed that the pumping is homogeneous and isotropic, and the atoms at the levels 1 and 0 are produced with a Maxwellian velocity distribution. F(v) is the Maxwellian distribution normalized to unity; $\gamma_1' N_1$ is the number of atoms produced per unit time at each of the Zeeman sublevels of the state 1 as a result of the pumping, $\gamma_0 N_0$ is the same for the state 0; $\mathbf{E}$ is the field in the laser; $\mathbf{d}_{m\mu}$ are the matrix elements of the dipole moment.

We now assume that $\gamma_1' \ll \tilde{\gamma}_1$ and that the dragging is complete, i.e., the mean free path of the photon of the resonant radiation at the center of the line l_0 is much smaller than the radius of the tube:

$$l_0 = (2j_g+1)8\pi^{3/2}u / (2j_1+1) n_g \lambda^3 \tilde{\gamma}_1, \tag{3}$$

(here j_1 and j_g are the total momenta of the level 1 and of the ground level, n_g is the concentration of the atoms of the working medium in the ground state, $u = (2T/M)^{1/2}$, T is the temperature in energy units, and λ is the wavelength of the dragged radiation). For an He-Ne laser at a partial neon pressure of 0.1 Torr, for a resonant transition from the level $2s_2$, we have $l_0 = 10^{-2}$ cm, so that the dragging is indeed almost complete.

We shall assume the dragging to be so strong that the decay constant characterizing the change of the total number of excited atoms at the level 1 reduces to γ_1'.

We shall seek stationary solutions of the system of equations (1) by the method of successive approximations. The field $\mathbf{E}$ will be written in the form

$$\mathbf{E} = (\mathbf{E}_0 e^{i\omega t} + \mathbf{E}_0^* e^{-i\omega t}) \sin kz.$$

In the zeroth approximation in the field $\mathbf{E}$, we have

$$f_{mm'}^{(0)} = N_1 F(v)\delta_{mm'}, \qquad \varphi_{\mu\mu'}^{(0)} = \tilde{N}_0 F(v)\delta_{\mu\mu'}, \tag{4}$$

$$\tilde{N}_0 = N_0 + \frac{\gamma}{\gamma_0}\frac{2j_1+1}{2j_0+1}N_1. \tag{5}$$

We have used here the following property of the kernel K:

$$\int d^3r' \int d^3v' \sum_{m_1 m_1'} K_{mm'}^{m_1 m_1'}(\mathbf{r}-\mathbf{r}', \mathbf{v}, \mathbf{v}')\delta_{m_1 m_1'} F(v') = \delta_{mm'} F(v), \tag{6}$$

where F(v) is the Maxwellian distribution.

In second order in the field $\mathbf{E}$, we have an equation for the spatially-homogeneous part of the density matrix of the upper working level $f_{mm'}^{(2)}$:

$$(\gamma_1' + \tilde{\gamma}_1) f^{(2)}_{mm'} - \tilde{\gamma}_1 (\hat{L} f^{(2)})_{mm'} = R_{mm'},$$

where

$$R_{mm'} = -\frac{NF(v)}{2\hbar^2} \sum_{\mu} (\mathbf{E}_0^* \mathbf{d}_{m\mu})(\mathbf{E}_0 \mathbf{d}_{\mu m'}) \left[\frac{\gamma_{10}}{\gamma_{10}^2 + (\delta + kv)^2} + \frac{\gamma_{10}}{\gamma_{10}^2 + (\delta - kv)^2} \right]. \quad (7)$$

Here $N = N_1 - \tilde{N}_0$ and $\delta = \omega - \omega_0$. The quantity $R_{mm'}$, due to the induced transitions, is a sharp function of the projection of the velocity of the atom on the direction of the wave vector **k**.

It is convenient to represent $f^{(2)}_{mm'}$ as the sum of a sharp function and a smooth function of the velocity:

$$f^{(2)}_{mm'} = R_{mm'}/\gamma_1 + f^{s}_{mm'}. \quad (8)$$

For the smooth function $f^{s}_{mm'}$ we have the equation

$$(\gamma_1' + \tilde{\gamma}_1) f^{s}_{mm'} - \tilde{\gamma}_1 (\hat{L} f^{s})_{mm'} = \frac{\tilde{\gamma}_1}{\nu} (\hat{L} R)_{mm'}. \quad (9)$$

Equation (9) can be solved in the case $\gamma_1' \ll \gamma_1$ under consideration. We seek the solution in the form

$$f^{s}_{mm'} = A \delta_{mm'} F(v) + \eta_{mm'}, \quad (10)$$

where A is a certain constant and $\eta_{mm'}$ is a small correction. Then, using the property (6) and neglecting the term $\gamma_1' \eta_{mm'}$, we obtain an equation for $\eta_{mm'}$:

$$\tilde{\gamma}_1 \eta_{mm'} - \tilde{\gamma}_1 (\hat{L} \eta)_{mm'} = -\gamma_1' \delta_{mm'} A F(v) + (LR)_{mm'}. \quad (11)$$

The condition for the solvability of this equation determines the constant A. Indeed,

$$\int d^3 v \sum_{m} (\hat{L} \eta)_{mm} = \int d^3 v \sum_{m} \eta_{mm} \quad (12)$$

for any function $\eta_{mm'}$, since under the conditions of total dragging the spontaneous transition to the ground state does not change the total number of excited atoms at the level 1. From (11) and (12) we obtain the quantity A:

$$A = \frac{1}{(2j_1 + 1)\gamma_1'} \int d^3 v \sum_{m} R_{mm}(\mathbf{v}). \quad (13)$$

Integrating in (13) with respect to the velocities (with allowance for the fact that $\gamma_{10} \ll ku$) and neglecting the correction $\eta_{mm'}$, we obtain with the aid of (8) and (10) a final expression for the density matrix of the upper working level 1 in second order in the field:

$$f^{(2)}_{mm'} = -\frac{NF(v)}{2\hbar^2} \left\{ \frac{1}{\gamma_1} \sum_{\mu} (\mathbf{E}_0^* \mathbf{d}_{m\mu})(\mathbf{E}_0 \mathbf{d}_{\mu m'}) \left[\frac{\gamma_{10}}{\gamma_{10}^2 + (\delta + kv)^2} + \frac{\gamma_{10}}{\gamma_{10}^2 + (\delta - kv)^2} \right] + \frac{\delta_{mm'}}{\gamma_1'} |\mathbf{E}_0|^2 \frac{2\sqrt{\pi} d^2}{3(2j_1 + 1)} \frac{1}{ku} \exp\left(-\frac{\delta^2}{k^2 u^2} \right) \right\},$$

where

$$d^2 = \sum_{m\mu} |\mathbf{d}_{m\mu}|^2. \quad (14)$$

The first term in the curly brackets in (14) describes the formation of the Bennet holes. The second term is a reflection of the settling of the entire distribution of the atoms at the level 1 with respect to the velocities, due to the dragging of the resonant radiation. In the region of the holes, the ratio of the second term to the first is determined by the parameter $\gamma_1 \gamma_{10} / \gamma_1' ku$.

The appearance in formula (14) of an addition, smoothly dependent on the velocity, to the density matrix causes the change of the spontaneous emission from the upper working level, due to the laser field, to occur not only in the region of the holes (narrow lines[6]), but also in the entire Doppler contour (broad line).

This singularity appears also for transitions beginning at the lower working level. The reason for the appearance of the broad line is in this case the spontaneous transition 1 → 0 (the term containing the matrix Γ in Eqs. (1)). We shall not present here expressions for the spontaneous emission from the working levels of the laser. We note only that the integral intensity of the broad line is larger than the intensity of the narrow lines in a ratio γ_1/γ_1'. The broad line is in the main not polarized. This is connected with the fact that in each act of absorption of a resonant photon the polarization returns only partially, as a result of which the relaxation of the orientation and the alignment of the atoms occur not with a small decay constant γ_1', but with a decay constant on the order of the value of γ_1 itself. The polarized part of the broad line is due to the function $\eta_{mm'}$, which we have discarded. The integral intensity of the polarized part of the broad line is of the same order as the intensity of the narrow lines.

We note that formula (14) and the results that follow concerning the shape of the Lamb dip can be obtained by replacing the integral term $(\hat{L}f)_{mm'}$ in Eqs. (1) by the expression

$$(\hat{L} f)_{mm'} \to \frac{\delta_{mm'}}{2j_1 + 1} F(v) \int d^3 v' \sum_{m_1} f_{m_1 m_1}(\mathbf{r}, \mathbf{v}', t).$$

Neglect of $\eta_{mm'}$ in (10) is in essence equivalent to such a replacement.

It must be borne in mind, however, that this replacement is valid only for that part of the density matrix $f_{mm'}$, which changes sufficiently slowly in space. A characteristic scale is the photon mean free path l_0 (formula (3)), which determines the spatial dispersion of the kernel K in formula (2). In the equation for the spatially modulated part $f_{mm'}$ (which is significant in the higher approximations with respect to the field or with respect to the ratio γ_{10}/ku), the integral term must be discarded, since usually $k l_0 \gg 1$.

SHAPE OF THE LAMB DIP

We now proceed to determine the dependence of the generation intensity on the detuning δ. Substituting (14) (and the corresponding expression for $\varphi_{mm'}$) in the third equation of (1), we can find $\psi_{\mu m}$ in third order in the field, and calculate the dipole moment of the gas in the mode under consideration, using the formula

$$\mathbf{P} = \frac{2}{l} \int_0^l dz \sin kz \int d^3 v \sum_{m\mu} \mathbf{d}_{m\mu} \psi_{\mu m}.$$

Here l is the resonator length. We then obtain

$$P_1 = (a - b_{11}|E_1|^2 - b_{1,-1}|E_{-1}|^2) E_1,$$
$$P_{-1} = (a - b_{-1,1}|E_1|^2 - b_{-1,-1}|E_{-1}|^2) E_{-1}. \quad (15)$$

Here $P_{\pm 1}$ and $E_{\pm 1}$ are the circular components respectively of the dipole moment and of the field, and a is the linear polarizability:

$$a = i \frac{d^2 N}{3\hbar} \frac{\sqrt{\pi}}{ku} W\left(\frac{\delta}{ku} \right), \quad W(z) = e^{-z^2} \left(1 - \frac{2i}{\sqrt{\pi}} \int_0^z e^{t^2} dt \right). \quad (16)$$

The coefficients b, which characterize the nonlinearity, were calculated by us in[4] without allowance for the dragging. When dragging is taken into account we have

$b_{qq'} = b^0_{qq'} + b^s$, where $b^0_{qq'}$ are quantities calculated in[4]:

$$b_{11}^0 = b^0_{-1,-1} = igA_1^0\left(\frac{1}{\gamma_1}+\frac{1}{\gamma_0}\right)\left(1+\frac{\gamma_{10}}{\gamma_{10}+i\delta}\right)\mathcal{E},$$
$$b^0_{1,-1} = b^0_{-1,1} = ig(A_2^0 + A_3^0)\left(\frac{1}{\gamma_1}+\frac{1}{\gamma_0}\right)\left(1+\frac{\gamma_{10}}{\gamma_{10}+i\delta}\right)\mathcal{E}, \quad (17)$$
$$g = \sqrt{\pi}d^4N(4\hbar^3ku\gamma_{10})^{-1}, \quad \mathcal{E} = \exp\{-\delta^2/(ku)^2\}.$$

The numbers A^0_n, which depend only on the momenta j_1 and j_0, have been calculated in[4]. For $j_1 = 1$ and $j_0 = 2$ we have $A^0_1 = 23/450$, $A^0_2 = 1/900$, and $A^0_3 = 7/300$. The addition b^s resulting from the dragging is given by

$$b^s = ig\frac{4\sqrt{\pi}}{9(2j_1+1)}\frac{\gamma_{10}}{\gamma_1' ku}\mathcal{E}W\left(\frac{\delta}{ku}\right). \quad (18)$$

In formulas (17) and (18) we have neglected the influence of the spontaneous transition 1 → 0 (the term with $\Gamma^{mm'}_{\mu\mu'}$ in Eqs. (1)) in view of the smallness of the ratio γ/γ_1.

With the aid of expressions (15) we can determine the intensity of the generation $I = |E_0|^2$ in the stationary regime (see, e.g.,[4]). If the laser radiation is circularly polarized, then

$$I_1 = (a'' - (4\pi Q)^{-1})\,/\,b_{11}''. \quad (19)$$

Here Q is the quality factor of the resonator, and a'' and b''_{11} are the imaginary parts of the coefficients a and b_{11}.

In the case of plane polarization, the intensity I_x is given by another formula:

$$I_x = (a'' - (4\pi Q)^{-1})\,/\,b_{xx}, \quad (20)$$

where $b_{xx} = (1/4)(b_{11} + b_{1,-1} + b_{-1,1} + b_{-1,-1})$.

The type of polarization can be determined both at the moments of the working levels[4] and by the anisotropy of the resonator. Using formulas (16)—(18), we can represent both expressions (19) and (20) in the form of Szoke and Javan[1]:

$$I = \text{const}\cdot\left[1-\frac{N^{(0)}}{N\mathcal{E}}\right]\left(1+\frac{\gamma_{10}\bar{\gamma}_{10}}{\gamma_{10}^2+\delta^2}\right)^{-1}, \quad (21)$$

where $N^{(0)}$ is the threshold value of N at $\delta = 0$. For $\bar{\gamma}_{10}$ we obtain the following expression:

$$\bar{\gamma}_{10} = \gamma_{10}\left[1+\frac{4\sqrt{\pi}}{3}\beta\frac{\gamma_{10}}{\gamma_1' ku}\left(\frac{1}{\gamma_1}+\frac{1}{\gamma_0}\right)^{-1}\right]^{-1}. \quad (22)$$

We recall that γ_1 and γ_0 are the reciprocal lifetimes of the upper and lower levels, γ_{10} is the natural line width, γ_1' is the probability of the spontaneous radiative transition from the level 1 to all the states except to the ground state.

The numerical coefficient β in formula (22) depends on the momenta of the working levels j_1 and j_0 and is different for planar and circular polarization. This coefficient can be expressed in terms of 6j-symbols. For plane polarization

$$\beta = \frac{1}{2j_1+1}\left[\begin{Bmatrix}1 & 1 & 0\\ j_0 & j_0 & j_1\end{Bmatrix}^2 + 2\begin{Bmatrix}1 & 1 & 2\\ j_0 & j_0 & j_1\end{Bmatrix}^2\right]^{-1}.$$

For circular polarization

$$\beta = \frac{1}{2j_1+1}\left[\begin{Bmatrix}1 & 1 & 0\\ j_0 & j_0 & j_1\end{Bmatrix}^2 + \frac{3}{2}\begin{Bmatrix}1 & 1 & 1\\ j_0 & j_0 & j_1\end{Bmatrix}^2 + \frac{1}{2}\begin{Bmatrix}1 & 1 & 2\\ j_0 & j_0 & j_1\end{Bmatrix}^2\right]^{-1}$$

The values of the coefficients β for the case when $j_1 = 1$ are listed in the table.

Formula (22) leads to the important conclusion that $\bar{\gamma}_{10}$ is smaller than γ_{10} even in the absence of collisions. For the neon line $\lambda = 1.15\ \mu$, an estimate gives $\gamma_{10}/\bar{\gamma}_{10} \approx 1.5$ in the case of plane polarization.

So far we have assumed for simplicity that the time τ of the emergence of the resonant photon from the volume is large compared with the time $1/\gamma_1'$. If this is not the case, then it is necessary to replace in formula (22) the quantity γ_1' by $\gamma_1' + \tau^{-1}$. An estimate of the time τ is given in[5]. In the case of strong dragging $\bar{\gamma}_1\tau \gg 1$ it is possible to use Holstein's result of variational calculation[7] for an infinite cylinder:

$$\bar{\gamma}_1\tau = \frac{5}{8}\frac{L}{l_0}\left(\pi\ln\frac{L}{l_0}\right)^{1/2},$$

where l_0 is the photon mean free path in the center of the line (formula (3)), and L is the radius of the laser beam. Formulas (18) and (20) can be generalized in obvious fashion to the case when dragging of the spontaneous radiation from the lower working level (to the metastable level) is also significant.

We note in conclusion that radiation dragging can influence appreciably not only the form of the Lamb dip, but also other characteristics of the laser. As already mentioned above, the spectrum and the polarization of the fluorescence from the working levels change. In particular, the degree of polarization drops sharply. Contributions to the polarized component of the fluorescence are made not only by the narrow lines, but also by the broad part of the spectrum.

Allowance for dragging is very important also when mode competition is considered. Because of the mixing of the excited atoms with respect to velocities and with respect to the Zeeman sublevels, an effective interaction takes place between modes that are separated in frequency or that have different polarizations.

The appreciable settling of the entire amplification contour in generation in a He-Ne laser was observed by Spiller[8], who also observed other features characteristic of a homogeneously broadened line. It can be assumed that these effects are the consequences of the dragging of the resonance radiation.

Values of the coefficients β at $j_1 = 1$

j_0	Circular polarization	Plane polarization
0	1	1
1	2	2
2	50/23	50/17

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