Quasiresonant Stark broadening of optical spectra of quantum systems in a Gaussian noise field

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A theory is developed of the Stark broadening of an optical absorption line by a quantum system on excitation of a transition (adjoining the probed transition) by a quasiresonant noise electromagnetic field representing a complex Gaussian random process. A unified treatment is given of the evolution of the line profile as the spectral width of the field varies from small values (quasistatic limit) to asymptotically large values when the field is effectively "switched off."

1. Investigations of optical spectra of quantum systems in external fields with randomly varying parameters is one of the important tasks in optical spectroscopy. The interest in this task is stimulated by a large number of different physical problems including the dynamic Stark effect in the field of incoherent optical radiation,[1-8] resonant multiphoton ionization of atoms by radiation of high-power multimode lasers,[7-12] properties of magnetic resonance in spin systems excited by an rf noise field,[13,14] investigations of the profile of a cyclotron radiation line in precision experiments on the anomalous moment of an electron, Stark broadening of spectral lines in a turbulent plasma,[16] etc. In all these problems the response of a quantum system to a weak probe field manifests the full set of correlation functions of a noise field because a quantum system is a strongly nonlinear converter of a random process. This manifestation depends on the relationship between the width of the noise field spectrum, average depth of modulation of the natural frequency of the quantum system (average Stark shift), offset between resonances, and natural as well as field-induced widths of quantum transitions. Fairly obvious results are obtained in the following two opposite limiting cases. In the case of an infinitesimal narrow spectrum of a strong stochastic field the expression for the probability of a resonant optical transition under the action of probe radiation can be derived in the quasistatic approximation by averaging the relevant expressions for a strong monochromatic field (with fixed parameters) over the distribution of its intensity. In the case of a complex Gaussian random process, which we shall consider in the present paper, this is an exponential distribution of intensities (or a Rayleigh distribution of amplitudes). In the opposite limiting case of a very wide spectrum the response of a quantum system is governed by quantities averaged over fast fluctuations of the field. In this case the expression for the probability of a resonance transition is of the same form as for a monochromatic field. However, the Stark shifts and decay widths of the levels induced by the noise field should be replaced by the corresponding average values. Less evident and much more difficult to analyze is the intermediate case when neither the quasistatic approximation nor the approximation of a wide spectrum is valid.

The simplest model which makes it possible to study the relationships is a two-level system in the field of a strong nonresonant noise radiation and a weak probe radiation which is in resonance with a quantum transition. Then, the influence of the nonresonant noise field on a selected pair of levels reduces to increments in the phases of the wave functions of these levels dependent on the intensity of the field. Such changes in the phases can be obtained (allowing for the nonresonant nature of the noise field) by adiabatic inclusion of the contributions of the rest of the spectrum of the quantum system in equations describing the dynamics of this system as a whole. Such a description is essentially equivalent to inclusion of the dynamic shift of levels proportional to the instantaneous value of the noise field intensity.[2,5,7,11,12] The strongest perturbation of the spectrum of the quantum system occurs however under quasiresonant conditions when the central frequency of the noise field is close to the frequency of one of the quantum transitions adjoining the one which is probed (see Fig. 1). Then the above adiabatic approximation becomes invalid even in the case of a finite offset of the central frequency of the noise field from the frequency of the 2–3 transition because of resonant mixing of the levels 2 and 3 by the Fourier components of the field in the wing of the spectral distribution.

A fairly complete theoretical analysis of the quasiresonant Stark effect under these conditions has been made only for a somewhat artificial model (discontinuous Markov process) of a laser noise field.[3] In the case of a complex Gaussian process which describes well the radiation emitted by multimode lasers with modes which are not locked[10] it has been suggested that use can be made of iterative procedures convenient in numerical calculations.[6,17]

We shall develop an analytic theory of optical spectra of quantum systems subjected to the quasiresonant action of a noise field representing a complex Gaussian random process with an arbitrary spectral width. We shall provide a unified treatment of the evolution of the spectra of quantum systems for spectral widths of the field varying from infinitesimally small values when the quasistatic approximation is valid to asymptotically large ones when the field is effectively "switched off." We shall find the limits of validity of the results obtained in the adiabatic approximation and demonstrate the radical changes in the spectra of quantum systems when their parameters are outside these limits. The results obtained however are valid only if the intensity of the noise field is insufficient for the Autler–Townes splitting of the spectra. A comparison with the results of Ref. 3 is made within the limits of validity of the theory. We shall show that modeling of a noise field by a discontinuous Markov process describes correctly, in the qualitative sense, the behavior of the spectra in the limiting cases, but gives rise to quantitative

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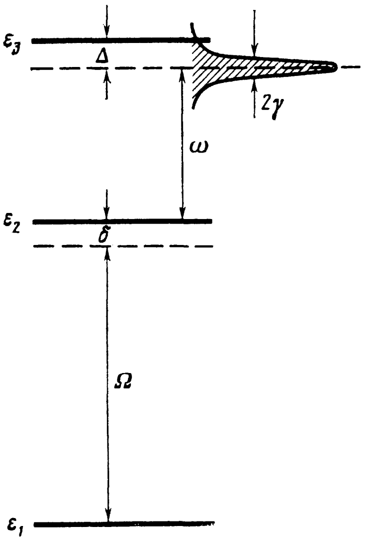

FIG. 1. Three-level quantum system in the field of quasiresonant noise (ω) and probe (Ω) radiations.

discrepancies compared with the model of a Gaussian process in the intermediate case.

2. We shall consider a three-level quantum system consisting of nondegenerate levels 1, 2, and 3 (with energies $\varepsilon_1 < \varepsilon_2 < \varepsilon_3$) and subjected to electromagnetic fields (Fig. 1). A strong stochastic field $\mathscr{F} = 2\mathrm{Re}[F(t)e^{-i\omega t}]$ has a carrier frequency ω close to a natural frequency of the transition $\omega_{32} = (\varepsilon_3 - \varepsilon_2)/\hbar$. The amplitude of this field $F(t)$ will be regarded as a stationary complex Gaussian random process with zero average and with a correlation function

$$\langle F(t)F^*(0)\rangle=\langle F^*(t)F(0)\rangle=B(t)={}^1/_2F_0{}^2e^{-\gamma|t|}. \qquad (1)$$

Here, F_0^2 is the average intensity of the noise field and γ^{-1} is the decay time of the correlation (1), governing the width of a Lorentzian profile of the radiation spectrum $\mathscr{F}(t)$.

We shall investigate the profile of an absorption line of a weak probe radiation $E = 2\mathrm{Re}(E_0 e^{-i\Omega t})$ which is in resonance with the 1–2 transition $[|(\omega_{21} - \Omega)/\omega_{21}| \ll 1]$. The imaginary part $\chi''(\Omega)$ of the polarizability $\chi(\Omega)$, governing the absorption coefficient of probe radiation, is proportional to the off-diagonal element σ_{12} of the density matrix of the system averaged over the ensemble of realizations of the random process $F(t)$. To lowest order in the interaction with weak radiation we can describe the vector $\boldsymbol{\sigma}$, composed of the matrix element σ_{12} of interest to us and the associated element σ_{13}, by the following equation

$$\frac{d}{dt}\boldsymbol{\sigma}+A(t)\boldsymbol{\sigma}=\mathbf{b}(t), \quad \boldsymbol{\sigma}=\begin{pmatrix}\sigma_{12}\\ \sigma_{13}\end{pmatrix}. \qquad (2)$$

Here,

$$A(t)=\begin{bmatrix}\Gamma_{12}, & i\hbar^{-1}d_{32}F(t)e^{i\Delta t}\\ i\hbar^{-1}d_{23}F^*(t)e^{-i\Delta t}, & \Gamma_{13}\end{bmatrix},$$

$$\mathbf{b}(t)=\begin{bmatrix}-i\hbar^{-1}d_{12}E_0{}^*e^{-i\delta t}\\ 0\end{bmatrix}, \qquad (3)$$

where $\Delta = \omega_{32} - \omega$; $\delta = \omega_{21} - \Omega$; d_{ij} is the dipole matrix element; Γ_{ij}^{-1} is the transverse relaxation time of the i–j transition. The solution of Eq. (2) with the initial condition $\boldsymbol{\sigma}(t_0) = \boldsymbol{\sigma}_0$ is

$$\boldsymbol{\sigma}(t)=U(t,t_0)\boldsymbol{\sigma}(t_0)+\int_{t_0}^{t} dt_1\,U(t,t_1)\mathbf{b}(t_1), \qquad (4)$$

where the evolution matrix $U(t, t_1)$ satisfies the homogeneous equation

$$\frac{d}{dt}U(t,t_1)+A(t)U(t,t_1)=0 \qquad (5)$$

and the condition $U(t_1, t_1) = 1$. Under steady-state absorption conditions ($t_0 \to -\infty$), we have

$$\sigma_{12}(t)=-\frac{id_{12}E_0{}^*}{\hbar}\int_{-\infty}^{t} dt_1\,U^{11}(t,t_1)e^{-i\delta t_1} \qquad (6)$$

(the matrix indices of U is shown as superscripts to distinguish them from the indices corresponding to atomic levels). Equation (5) readily yields

$$\frac{d}{dt}\bar U^{11}(t,t_1)=-\frac{|d_{23}|^2}{\hbar^2}F(t)\int_{t_1}^{t} dt_2 e^{i\tilde\Delta(t-t_2)}F^*(t_2)\bar U^{11}(t_2,t_1),$$

$$\bar U^{11}(t,t_1)=e^{\Gamma_{12}(t-t_1)}U^{11}(t,t_1), \quad \tilde\Delta=\Delta+i\Gamma, \quad \Gamma=\Gamma_{13}-\Gamma_{12}. \qquad (7)$$

We shall consider the specific case when $\Gamma > 0$.

We shall assume that the field $F(t)$ is not too high, i.e.,

$$\Omega_R{}^2/|\tilde\Delta|^2\ll1, \quad \Omega_R=|d_{23}|F_0/\hbar. \qquad (8)$$

In this case, integrating Eq. (7) by parts and retaining the main terms of the expansion in the above parameter, we obtain

$$\bar U(t,t')=\exp\left\{-\frac{|d_{23}|^2}{\hbar^2}\int_{t'}^{t} dt_1\int_{t'}^{t_1} dt_2F^*(t_2)F(t_1)e^{i\tilde\Delta(t_1-t_2)}\right\}. \qquad (9)$$

Consequently,

$$\chi''(\Omega)=\frac{|d_{12}|^2}{\hbar}\mathrm{Re}\int_0^\infty d\tau\,e^{i\delta\tau-\Gamma_{12}\tau}J(\tau), \qquad (10)$$

$$J(\tau)=\left\langle\exp\left\{-\frac{|d_{23}|^2}{2\hbar^2}\int_0^\tau\int_0^\tau dt_1\,dt_2F^*(t_2)F(t_1)e^{i\tilde\Delta|t_1-t_2|}\right\}\right\rangle \qquad (11)$$

[the angular brackets denote averaging over all possible realizations of the stationary random process $\mathscr{F}(t)$].

If the width of the radiation spectrum γ is considerably less than the quantity $|\tilde\Delta|$, then the integral with respect to t_2 in Eq. (9) is dominated by the contribution of the region where $t_2 \sim t_1$. Then, in the integral with respect to t_2 in Eq. (9) we can take out $F(t_2)$ corresponding to $t_2 = t_1$ and integrate, ignoring the rapidly oscillating term which makes a contribution small in terms of the parameter $\Omega_R/|\tilde\Delta|$. In this case we obtain

$$J(\tau)=\left\langle\exp\left\{-i\frac{|d_{23}|^2}{\hbar^2\tilde\Delta}\int_0^\tau dt\,|F(t)|^2\right\}\right\rangle. \qquad (12)$$

If $|\Delta| \gg \Gamma_{13} - \Gamma_{12}$ is obeyed, the quantity $|d_{23}|^2|F(t)|^2/\hbar^2\Delta$ represents an instantaneous value of the resonant part of the Stark shift of the natural frequency of the transition 1–2 (for simplicity, a small antiresonant part is ignored). The expression in the argument of the exponential function of Eq. (12) represents an adiabatic increase in the wave-function phase of the level 2 caused by mixing with the level 3 due to the interaction of the quantum system with the noise field. The profile of an optical transition line considered in the approximation of Eq. (12) had been discussed earlier in Refs. 2, 11, 12, and 15. Mathematically similar problems related to averaging of quantities of the type described by Eq. (12) have been encountered also[18] when dealing with the interaction of charged relativistic particle with a noise electromagnetic field, as well as in a study of photocount statistics[19] (see also Ref. 20). The average of Eq. (11) can be calculated conveniently by representing it in the form of a continuum integral over all possible realizations of the random process $F(t)$ (Ref. 21):

$$J(\tau)=\int D[F(t)]D[F^*(t)]P([F(t)],[F^*(t)])\cdot\exp\left\{-\int_0^\tau\int_0^\tau F^*(t_1)K(t_1,t_2)F(t_2)\,dt_1\,dt_2\right\}, \quad (13)$$

$$K(t_1,t_2)=(|d_{23}|^2/2\hbar^2)\exp\{i\tilde{\Delta}|t_1-t_2|\}. \quad (14)$$

Here, $P\{[F(t)],[F^*(t)]\}$ is a functional of the distribution of $F(t)$ and $F^*(t)$:

$$P\{[F(t)],[F^*(t)]\}\sim\exp\left\{-\int_0^\tau\int_0^\tau F^*(t_1)B^{-1}(t_1,t_2)F(t_2)\,dt_1\,dt_2\right\}. \quad (15)$$

The quantity $B^{-1}(t_1,t_2)$ is the kernel of an integral operator, which is the inverse of Eq. (1):

$$\int_0^\tau dt'\,B^{-1}(t_1,t')B(t',t_2)=\delta(t_1-t_2). \quad (16)$$

The Gaussian functional integral of Eq. (13) can be calculated directly:

$$J(\tau)=[\det(\hat{O})]^{-1}. \quad (17)$$

The kernel of the integral operator $\hat{O}$ is

$$O(t_1,t_2)=\delta(t_1-t_2)+\int_0^\tau dt'\,B(t_1,t')K(t',t_2). \quad (18)$$

The determinant of $\hat{O}$ is an infinite product of all eigenvalues of this operator, which are found using the equation

$$(\lambda-1)f(t)=\int_0^\tau\int_0^\tau dt'\,dt''\,B(t,t')K(t',t'')f(t'') \quad (19)$$

[$f(t)$ is the eigenfunction of the operator $\hat{O}$ corresponding to the value of λ].

We can easily demonstrate that the quantities $B(t_1,t_2)$ and $K(t_1,t_2)$ satisfy the equations

$$\left(\frac{d^2}{dt_1^2}-\gamma^2\right)B(t_1,t_2)=\gamma F_0^2\delta(t_1-t_2), \quad (20)$$

$$\left(\frac{d^2}{dt_1^2}+\tilde{\Delta}^2\right)K(t_1,t_2)=i\frac{|d_{23}|^2}{\hbar^2}\tilde{\Delta}\delta(t_1-t_2). \quad (21)$$

Using Eqs. (20) and (21) we can show that the solutions of the integral equation (19) obey also a differential equation

$$\frac{d^4f}{dt^4}-(\gamma^2-\tilde{\Delta}^2)\frac{d^2f}{dt^2}+\left(i\frac{\gamma\tilde{\Delta}\Omega_R^2}{\lambda-1}-\gamma^2\tilde{\Delta}^2\right)f=0. \quad (22)$$

The general solution of Eq. (22) is

$$f(t)=\sum_{i=1}^4 C_i e^{L_i t}, \quad (23)$$

where L_i are the roots of the characteristic equation

$$L^4-(\gamma^2-\tilde{\Delta}^2)L^2+\left(i\frac{\gamma\tilde{\Delta}\Omega_R^2}{\lambda-1}-\gamma^2\tilde{\Delta}^2\right)=0, \quad (24)$$

i.e.,

$$L_{1,2}=\pm P,\qquad L_{3,4}=\pm Q,$$
$$P=\frac{1}{2^{1/2}}\left\{\gamma^2-\tilde{\Delta}^2+\left[(\gamma^2+\tilde{\Delta}^2)^2-i\frac{4\gamma\tilde{\Delta}\Omega_R^2}{\lambda-1}\right]^{1/2}\right\}^{1/2}, \quad (25)$$
$$Q=\frac{1}{2^{1/2}}\left\{\gamma^2-\tilde{\Delta}^2-\left[(\gamma^2+\tilde{\Delta}^2)^2-i\frac{4\gamma\tilde{\Delta}\Omega_R^2}{\lambda-1}\right]^{1/2}\right\}^{1/2}.$$

Substituting Eq. (23) into Eq. (19) and using Eq. (24), we obtain a system of four homogeneous linear equations for the quantities C_i. Equating the determinant of this system to zero, we obtain an equation which is satisfied by the eigenvalues λ:

$$\Phi(\lambda,\tau)=\Phi_+(\lambda,\tau)\Phi_-(\lambda,\tau)=0, \quad (26)$$

$$\Phi_\pm(\lambda,\tau)=\exp\left(\frac{P+Q}{2}\tau\right)\left[\frac{P+\gamma}{Q+i\tilde{\Delta}}-\frac{Q+\gamma}{P+i\tilde{\Delta}}\right]+\exp\left(-\frac{P+Q}{2}\tau\right)\left[\frac{P-\gamma}{Q-i\tilde{\Delta}}-\frac{Q-\gamma}{P-i\tilde{\Delta}}\right]$$
$$\pm\exp\left(\frac{P-Q}{2}\tau\right)\left[\frac{P+\gamma}{Q-i\tilde{\Delta}}-\frac{Q-\gamma}{P+i\tilde{\Delta}}\right]\pm\exp\left(\frac{Q-P}{2}\tau\right)\left[\frac{P-\gamma}{Q+i\tilde{\Delta}}-\frac{Q+\gamma}{P-i\tilde{\Delta}}\right]. \quad (27)$$

It is convenient to calculate the determinant of Eq. (17) by the following procedure (see, for example, Ref. 11). We can readily see that

$$\frac{d}{d\tau}\ln\det(\hat{O})=\sum_n\frac{d\lambda_n/dt}{\lambda_n}=-\sum_n\frac{\partial\Phi(\lambda,\tau)/\partial\tau}{\lambda\,\partial\Phi(\lambda,\tau)/\partial\lambda}\Bigg|_{\lambda=\lambda_n}. \quad (28)$$

The last sum in Eq. (28) can be represented by an integral in the complex λ plane along a closed contour of sufficiently large radius going round the point $\lambda=0$ in the positive direction:

$$\frac{d}{d\tau}\ln\det(\hat{O})=\frac{\partial\Phi(0,\tau)/\partial\tau}{\Phi(0,\tau)}-\frac{1}{2\pi i}\oint\frac{\partial\Phi(\lambda,\tau)/\partial\tau}{\lambda\Phi(\lambda,\tau)}d\lambda. \quad (29)$$

It follows from Eqs. (25) and (27) that

$$\frac{1}{\Phi_\pm(\lambda,\tau)}\frac{\partial\Phi_\pm(\lambda,\tau)}{\partial\tau}\Bigg|_{|\lambda|\to\infty}\to\frac{1}{2}(\gamma-i\tilde{\Delta}). \quad (30)$$

Increasing the radius of the contour in Eq. (29) to infinity, we finally obtain

$$J(\tau)=[\det(\hat{O})]^{-1}=\frac{\Phi_+(0,0)\Phi_-(0,0)}{\Phi_+(0,\tau)\Phi_-(0,\tau)}e^{(\gamma-i\tilde{\Delta})\tau}. \quad (31)$$

Equation (31) is derived bearing in mind that $\det(\hat{O}) = 1$ when $\tau = 0$.

3. We shall now consider the limiting cases of these formulas. We shall assume that

$$\Omega_R^2\gamma|\tilde{\Delta}|/|\gamma^2+\tilde{\Delta}^2|^2 \ll 1. \quad (32)$$

Within the limits of validity of Eq. (11) [see Eq. (8)] the inequality of Eq. (32) is disobeyed only for $\Delta \ll \Gamma$ and $\gamma \sim \Gamma$ When Eq. (32) is obeyed, the expressions in the system (25) simplify greatly ($\lambda = 0$):

$$P \approx q = [\gamma^2 + i\gamma\tilde{\Delta}\Omega_R^2/(\gamma^2+\tilde{\Delta}^2)]^{1/2},$$
$$Q \approx -i\tilde{\Delta} + \psi, \quad \psi = \gamma\Omega_R^2/2(\gamma^2+\tilde{\Delta}^2). \quad (33)$$

We thus find that

$$J(\tau) = 4\gamma q[(q+\gamma)^2 e^{q\tau} - (q-\gamma)^2 e^{-q\tau}]^{-1} e^{(\gamma-\psi)\tau}. \quad (34)$$

Far from a resonance ($\Delta \gg \Gamma$) and for $\gamma \ll \Delta$, Eq. (34) reduces to an expression derived in Ref. 11 using the adiabatic approximation to calculate the Stark shift. An expression analogous to Eq. (34) was obtained first in Ref. 20 by a different method. We shall first consider the case when

$$\gamma \ll \frac{\Omega_R^2}{|\tilde{\Delta}|} = \frac{|d_{23}|^2 F_0^2}{\hbar^2(\Delta^2+\Gamma^2)^{1/2}}. \quad (35)$$

We then have

$$\chi''(\Omega) = \frac{|d_{12}|^2}{\hbar}\,\mathrm{Re}\int_0^\infty \frac{\exp(i\delta\tau - \Gamma_{12}\tau)}{1+i\Omega_R^2\tau/2(\tilde{\Delta}+i\Gamma)}\,d\tau. \quad (36)$$

Equation (36) is readily reduced to

$$\chi''(\Omega) = \frac{|d_{12}|^2}{\hbar}\int_0^\infty dx\,\frac{(\Gamma_{12}+xs'')e^{-x}}{(\delta - xs')^2 + (\Gamma_{12}+xs'')^2}, \quad (37)$$

where

$$s' = \frac{\Omega_R^2\Delta}{2(\Delta^2+\Gamma^2)}, \quad s'' = \frac{\Omega_R^2\Gamma}{2(\Delta^2+\Gamma^2)}. \quad (38)$$

The quantities xs' and xs'' represent the shift and induced broadening of the level 2 by a field of intensity xF_0^2. As expected, in the quasistatic limit defined by Eq. (35) the line profile can be obtained from Eq. (37) by averaging over the exponential distribution of the intensity of the noise field of a modified Lorentzian profile. If $\Delta \gg \Gamma$, $s' \gg s''$, Γ_{12}, Eq. (37) describes a strongly asymmetric quasistatic profile due to the Stark broadening.[2–5]

In the limit opposite to that defined by Eq. (35), it follows from Eq. (34) that

$$\chi''(\Omega) = \frac{|d_{12}|^2}{\hbar}\frac{\Gamma_{eff}}{(\delta - s_{eff})^2 + \Gamma_{eff}^2}, \quad (39)$$

where

$$s_{eff} = \frac{\Omega_R^2}{2}\frac{\Delta}{\Delta^2+(\Gamma+\gamma)^2},$$
$$\Gamma_{eff} = \Gamma_{12} + \frac{\Omega_R^2}{2}\frac{(\Gamma+\gamma)}{\Delta^2+(\Gamma+\gamma)^2}. \quad (40)$$

Equation (39) describes a Lorentzian absorption profile shifted by an amount $s_{\rm eff}$ and including additional broadening proportional to the noise field intensity.

In contrast to the results of Refs. 2, 11, and 15, the expression obtained by us is valid under quasiresonant conditions and in the limit $\gamma \to \infty$ it gives an unshifted line with a natural width Γ_{12}. Figures 2a and 2b demonstrate evolution of the profile of an optical absorption line of frequency Ω due to the 1–2 transition when the spectral width γ of the noise field is varied within wide limits. For low values of γ (curve 1) the line profile has a typical asymmetric form characteristic of the quasistatic case. When γ increases to $\sim \Omega_R^2/2|\tilde{\Delta}|$ (i.e., when it becomes of the order of the average Stark shift), the line maximum is shifted and it becomes centered around $\delta = \delta_{st} = \Omega_R^2/2|\tilde{\Delta}|$ (curves 2 and 3). This stage of the line profile evolution can be described qualitatively also using the adiabatic approximation[2,11,15]; however, there are considerable quantitative differences compared with Refs. 2, 11, and 15 because of broadening of the Lorentzian line (see

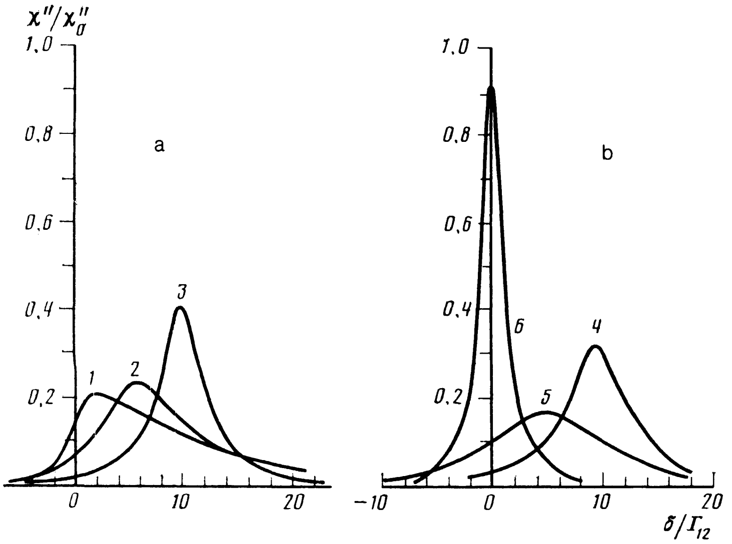

FIG. 2. Evolution of the profile of an absorption line when the spectral width of the noise field is varied ($\chi_0'' = |d_{12}|^2/\hbar\Gamma_{12}$, $\Delta/\Gamma_{12} = 10^3$, $\Gamma \ll \Delta$, $\delta_{st}/\Gamma_{12} = \Omega_R^2/2\Delta\Gamma_{12} = 10$): 1) $\gamma/\delta_{st} = 10^{-3}$; 2) 0.3; 3) 10; 4) 20: 5) 10^2; 6) 10^4.

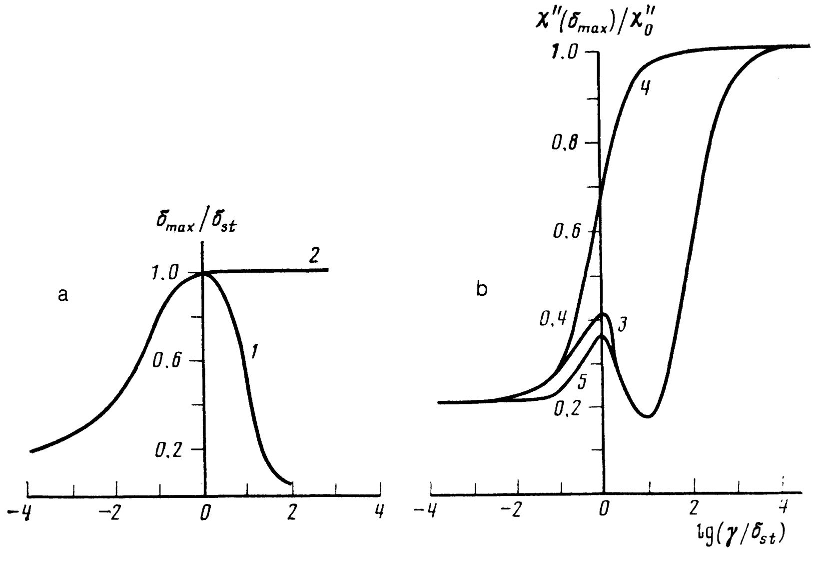

FIG. 3. Dependences of the position of the maximum δ_{max} of the absorption coefficient (a) and of the maximum value of the coefficient on the spectral width of the noise field (b) ($\Delta/\Gamma_{12} = 10^3$, $\Gamma \ll \Delta, \delta_{st}/\Gamma_{12} = 10$). Curves 1 and 3 were obtained in the present study, curves 2 and 4 were obtained in the adiabatic approximation, and curve 3 was obtained in the model of a discontinuous Markov process (see Ref. 3).

curve 3) due to transitions between the levels 2 and 3 induced by the noise field. Figure 2b shows the subsequent stage of the evolution of the absorption line profile on further increase in γ. In the limit $\gamma \to \infty$ the maximum of the line returns to its unshifted position and its width decreases right down to the natural width Γ_{12}. Figures 3a and 3b compare the positions of the maximum of the absorption coefficient and its value at this point obtained from Eq. (34) (curves 1 and 3) with those deduced using the adiabatic approximation (curves 2 and 4). Figure 3b includes also the result obtained using a model of a discontinuous Markov process for the field[3] with the same spectral width and the same average intensity (curve 5). It follows from Fig. 3 that if $\gamma/\Delta \sim 1$, the deviations from the adiabatic theory become very large. If $\Omega_R^2/2\Delta \lesssim \gamma \lesssim \Delta$, the models of complex Gaussian noise and of a discontinuous Markov process[3] give quantitatively different results (see Figs. 3b and 4). In the limits $\gamma \to 0$ and $\gamma \to \infty$ the results of both models are naturally identical.

Figure 5 shows the dependence of the shift of the line maximum on the average intensity of the noise field plotted for different values of the spectral width γ. At low intensities, when the average Stark shift is much less than γ ($\Omega_R^2/2|\tilde{\Delta}| \ll \gamma$), the shift of the line maximum depends linearly on the average field intensity. When the intensity exceeds this value, the dependence weakens considerably. If $\gamma \ll \Delta$ (curves 1, 2, and 3 in Fig. 5) the slope of the initial part of the dependence is not influenced by γ. The dependences plotted in Fig. 5 are in qualitative agreement with the experimental data.[2] It should be noted that if $\gamma \gtrsim \Delta$, the slope of the initial part decreases with γ (curve 4 in Fig. 5).

We shall now consider the case of an exact resonance $\Delta = 0$ (within the limits of the validity of the theory we must then have $\Omega_R^2/\Gamma^2 \ll 1$). Since in this case when γ is varied continuously, there is a range of values ($\gamma \sim \Gamma$) where the expansion of Eqs. (32) and (33) fails, we shall construct the line profile using the general formulas of Eqs. (27) and (31). Figure 6 demonstrates the corresponding evolution of the

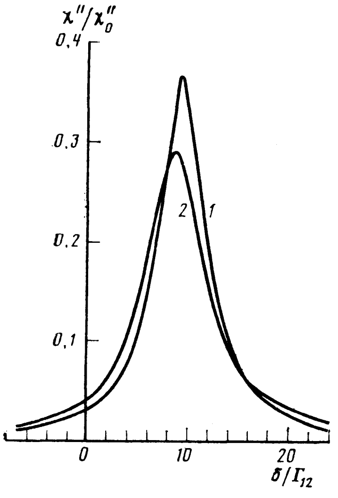

FIG. 4. Profile of an absorption line in the case when $\Delta/\Gamma_{12} = 10^3$, $\Gamma \ll \Delta$, $\delta_{st}/\Gamma_{12} = 10$, and $\gamma/\Gamma_{12} = 30$. Curve 1 is the result obtained in the present paper and curve 2 was obtained in the model of a discontinuous Markov process (see Ref. 3).

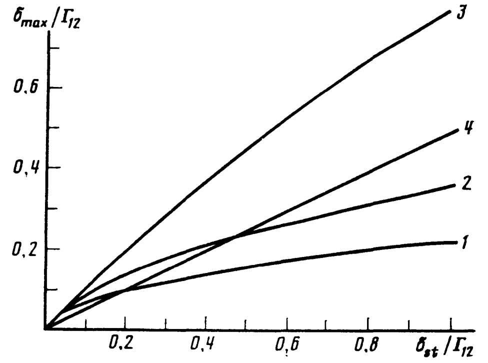

FIG. 5. Dependence of the shift of the maximum of the absorption line on the average intensity of the noise field obtained for different values of the spectral width of this field: 1) $\gamma/\Gamma_{12} = 3$; 2) 10; 3) 10^2; 4) 10^3.

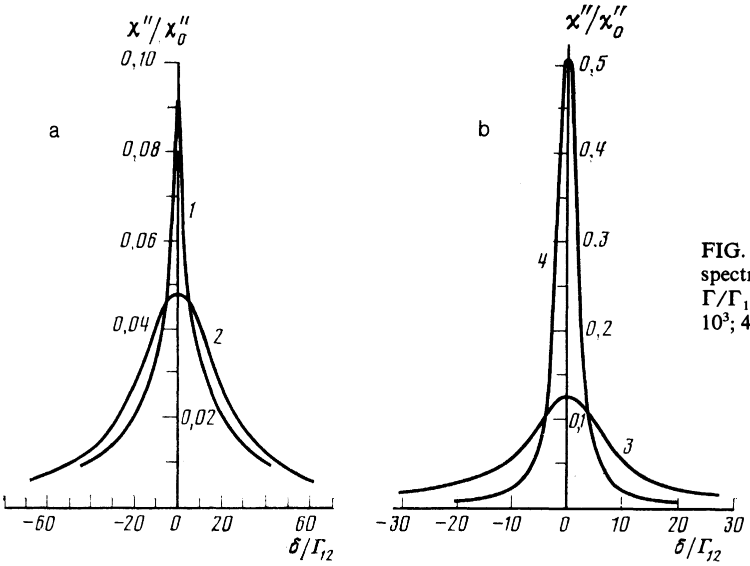

FIG. 6. Evolution of the absorption line profile due to variation of the spectral width of the noise field in the case of an exact resonance ($\Delta = 0$, $\Gamma/\Gamma_{12} = 3\times 10^2$, $\Omega_R^2/\Gamma\Gamma_{12} = 0.66\times 10^2$); 1) $\gamma/\Gamma_{12} = 10^{-2}$; 2) 10^2; 3) 10^3; 4) 10^4.

absorption line profile of probe radiation. If $\gamma \ll \Omega_R^2/2\Gamma$ the symmetric line profile is greatly extended at the center and it has elongated wings due to statistical peaks of the noise field (curve 1 in Fig. 6a). The central part of the band becomes broader on increase in γ and, consequently, its maximum amplitude decreases. The line profile approaches the Lorentzian form with a width $\Omega_R^2/2\Gamma$ (curve 2 in Fig. 6a). A further increase of γ to values of the order of Γ or more reduces the line width but its profile is retained. In the limit $\gamma \gg \Gamma$ the noise field is effectively "switched off," i.e., the line profile is a Lorentzian with a natural width Γ_{12}.

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