Characteristic Transformation of the Fluctuation Spectrum of Radiation Passing Through a Resonant Medium

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The transformation of the fluctuation spectrum of radiation interacting with an absorbing or amplifying medium is studied theoretically and experimentally. It is shown that the initial white fluctuation spectrum acquires a set of extrema which characterize details of the medium-energy structure and, in particular, such details as are masked in ordinary spectroscopic investigations by large Doppler broadening. Transformation of the noise spectrum of the Xe 3.507 and 5.57 μ lines and the Ne 3.39 μ line is studied experimentally.

RECENT investigations[1,2] were devoted to the power spectrum of the discharge radiation in Xe gas at the 3.507 μ line under conditions of population inversion in the corresponding transition. When a magnetic field was applied to the discharge, narrow extrema were observed in the receiver photocurrent spectrum and were interpreted as the result of interference between the excited magnetic states of the atoms under conditions when there is no coherence in their ensemble. Further experimental and theoretical investigations have shown that the observed phenomenon is indeed connected with interference of states, but it turned out that the situation involves not correlation in the elementary act of the spontaneous emission, but changes in the power spectrum of the fluctuating light field as a result of stimulated interaction with the atoms. We present here theoretical and experimental results confirming this interpretation of the phenomenon.

2. THEORY

The task of the theory is to describe the following experimental situation. There is a radiation source with sufficiently broad spectrum (the width criterion will be spelled out more accurately later on). The radiation is made to pass through a cell with an atomic vapor capable of absorbing or amplifying this radiation. A magnetic field is applied to the cell and splits the energy levels of the atoms. Radiation passing through the cell strikes a photoreceiver, the signal from which is spectrally analyzed. The radiation used in all experiments has had such a high spectral density that the wave component of the fluctuations of its intensity (i.e., connected with the interference of different harmonics of the field) greatly exceeded the shot-noise component. This circumstance, in conjunction with the low level of the intrinsic noise of the receiver, has made it possible to identify the photocurrent power spectrum with the power spectrum of the recorded radiation. The task of the theory reduces thereby to a calculation of the power spectrum of the radiation as it leaves the cell.

Assume that the radiation flux incident on the cell has an intensity $I_\beta(t)$ which is a stationary random function of the time. The index β denotes the observed component of the radiation polarization. After passing through the cell, the intensity is changed by an amount $W_\beta(t)$ and becomes equal to $J_\beta(t)$:

$$J_\beta(t) = I_\beta(t) - W_\beta(t). \tag{1}$$

We assume that $W_\beta(t) \ll I_\beta(t)$. Then the correlation function $S_\beta(\tau)$ of the outgoing radiation is expressed as follows (the angle brackets denote averaging over the time):

$$S_\beta(\tau) = \langle J_\beta(t) J_\beta(t+\tau)\rangle \approx \langle I_\beta(t) I_\beta(t+\tau)\rangle - \Delta S(\tau).$$

The quantity $\Delta S(\tau)$ represents here the changes in the statistics of the initial radiation resulting from its interaction with the atoms in the cell:

$$\Delta S(\tau) = S'(\tau) + S'(-\tau), \tag{2}$$

where

$$S'(\tau) = \langle I_\beta(t) W_\beta(t+\tau)\rangle.$$

The change $W_\beta(t)$ of the radiation intensity as it passes through the system of atoms is given in first-order perturbation theory by the formula:

$$W_\beta(t) = AI_\beta(t)\left[\sum_{m\mu\mu'} d_{m\mu}{}^{\beta}\varphi_{\mu\mu'}(t)\, d^{\beta}_{\mu'm} - \sum_{mm'\mu} d_{\mu m}{}^{\beta} f_{mm'}(t)\, d^{\beta}_{m'\mu}\right]. \tag{3}$$

The following notation is used here: A is a proportionality coefficient, $d^{\beta}_{m\mu}$ is the matrix element of the projection of the dipole-moment operator on the polarization vector β, $f_{mm'}$ and $\varphi_{\mu\mu'}$ are the density-matrix components pertaining to the sublevels m of the upper state and the sublevels μ of the lower state participating in the optical transition. Expression (3) has been derived[1)] under the assumption that the width of the incident-radiation spectrum greatly exceeds the natural widths Γ_1 and Γ_0 of the upper and lower states and their magnetic splittings $\omega_{mm'}$ and $\omega_{\mu\mu'}$.

The first sum in (3) corresponds to absorption of radiation by atoms going over from an aggregate of lower sublevels μ into the upper state. The second sum gives the contribution of the stimulated emission of the

[1)]Expression (3) is derived by starting from the formula $W(t) = \langle E(t) dP/dt\rangle$ (where P is the vector of the dipole moment of the system and E(t) is the electric field) by using a procedure analogous to that used in [3] to derive expression (14).

atoms under the influence of the same radiation $I_\beta(t)$.

The density-matrix components that depend on the intensity $I_\beta(t)$ can be determined from the following equations, derived from the Schrödinger equation in the first non-vanishing approximation in the perturbing light action[13]

$$\dot{f}_{mm'} + (\Gamma_1 + i\omega_{mm'}) f_{mm'} = -kN \sum_{\mu\alpha} d_{m\mu}{}^{\alpha} d_{\mu m'}^{\alpha} I_\alpha(t),$$
$$\dot{\varphi}_{\mu\mu'} + (\Gamma_0 + i\omega_{\mu\mu'}) \varphi_{\mu\mu'} = kN \sum_{m\alpha} d_{\mu m}{}^{\alpha} d_{m\mu'}^{\alpha} I_\alpha(t). \qquad (4)$$

Here k is a proportionality coefficient and N is the initial population difference between the system of the upper (m) and lower (μ) sublevels. For the sake of generality, summation was carried out over the polarization components of the radiation responsible for the occurrence of the components $f_{mm'}$ and $\varphi_{\mu\mu'}$ of the density matrix. We expand the intensities $I_{\alpha,\beta}(t)$ in a Fourier series in the finite time interval (0, T), with the aim of subsequently going over to a Fourier integral by means of the substitution $T^{-1} \rightarrow d\omega/2\pi$:

$$I_\alpha(t) = T^{-1} \sum_{\omega} I_\alpha(\omega) e^{-i\omega t}. \qquad (5)$$

Equations (4) then have a solution in the form of a Fourier series:

$$f_{mm'}(\omega) = -Nk \sum_{\mu'\alpha} d_{m\mu'} d_{\mu'm'}^{\alpha} I_\alpha(\omega) / [\Gamma_m + i(\omega_{mm'} - \omega)],$$
$$\varphi_{\mu\mu'}(\omega) = Nk \sum_{m'\alpha} d_{\mu m'}^{\alpha} d_{m'\mu'}^{\alpha} I_\alpha(\omega) / [\Gamma_\mu + i(\omega_{\mu\mu'} - \omega)]. \qquad (6)$$

Substituting (6) in (3) and using the expansion (5) for $I_\alpha(t)$ and $I_\beta(t)$ and analogous expressions for $f_{mm'}(t)$ and $\varphi_{\mu\mu'}(t)$, we have

$$W_\beta(t) = \frac{ANk}{T^2} \sum_{\substack{mm'\mu\mu' \\ \alpha\omega\omega'}} I_\alpha(\omega) I_\beta^*(\omega') \left[\frac{d_{m\mu}{}^{\beta} d_{\mu m'}{}^{\alpha} d_{m'\mu'}{}^{\alpha} d_{\mu' m}{}^{\beta}}{\Gamma_\mu + i(\omega_{\mu\mu'} - \omega)} + \frac{d_{\mu m}{}^{\beta} d_{m\mu'}{}^{\alpha} d_{\mu' m'}{}^{\alpha} d_{m'\mu}{}^{\beta}}{\Gamma_m + i(\omega_{mm'} - \omega)} \right] e^{-i(\omega-\omega')t}. \qquad (7)$$

We now proceed to calculate the correlation function $S'(\tau)$:

$$S'(\tau) = \frac{ANk}{T^3} \sum_{\substack{\omega\omega' \\ \omega''\alpha}} \langle I_\alpha(\omega) I_\beta^*(\omega') I_\beta(\omega'') \rangle e^{-i(\omega-\omega')\tau} e^{-i(\omega+\omega''-\omega')t} \sum_{\substack{mm' \\ \mu\mu'}} [\omega]. \qquad (8)$$

Here $[\omega]$ denotes the expression in the square brackets in (7). Owing to the independence of the different harmonics of the intensities $I_\alpha(t)$ and $I_\beta(t)$, the averaging leads to the vanishing of all the terms in the sum (8) except those corresponding to the following sets of frequency indices: 1) $\omega = \omega' = \omega'' = 0$; 2) $\omega = 0$, $\omega' = \omega''$; 3) $\omega'' = 0$, $\omega = \omega'$; 4) $\omega' = 0$, $\omega = -\omega''$. We are interested only in the last set. We write out the corresponding component of the correlation function $S_4'(\tau)$, changing over to a Fourier integral:

$$S_4'(\tau) = \frac{1}{2\pi} \int_{-\infty}^{\infty} S_4'(\omega) e^{i\omega\tau} d\omega, \qquad (9)$$

$$S_4'(\omega) = AkNI_\beta G_{I_\beta}(\omega) \sum_{\substack{mm' \\ \mu\mu'}} \left[\frac{d_{\mu m'}{}^{\beta} d_{m'\mu'}{}^{\beta} d_{\mu' m}{}^{\beta} d_{m\mu}{}^{\beta}}{\Gamma_0 + i(\omega_{\mu\mu'} - \omega)} + \frac{d_{m\mu'}{}^{\beta} d_{\mu' m'}{}^{\beta} d_{m'\mu}{}^{\beta} d_{\mu m}{}^{\beta}}{\Gamma_1 + i(\omega_{mm'} - \omega)} \right];$$

$$I_\beta = \lim_{T\to\infty} \frac{\langle I_\beta(0) \rangle}{T}, \qquad G_{I_\beta}(\omega) = \lim_{T\to\infty} \frac{\langle I_\beta^2(\omega) \rangle}{T}. \qquad (10)$$

Here I_β is the dc component of the intensity and $G_{I_\beta}(\omega)$ is the spectral density of the noise in the initial radiation.

The spectral density of the noise in the radiation passing through the cell can be written in the form

$$G_{J_\beta}(\omega) = G_{I_\beta}(\omega) - \Delta G(\omega),$$

and, according to formula (2), the component $\Delta G_4(\omega)$ of interest to us in the total increment $\Delta G(\omega)$ is given by

$$\Delta G_4(\omega) = 2\mathrm{Re}\, S_4'(\omega). \qquad (11)$$

In deriving (10) we have assumed that there is no correlation of the fluctuations in the orthogonal polarizations of the light. As a result, the intensity fluctuations of the observed component with polarization β do not depend on whether light of another polarization passes through the cell.

The remaining components of the increment $\Delta G(\omega)$, which have not been used here, either differ from zero only at zero frequency (the first and third set of frequencies), or are similar to $G_{I_\beta}(\omega)$ (second set of frequencies). We note that $\Delta G_2(\omega)$, corresponding to the second set of frequencies, duplicates the spectrum of the incident radiation, but its value depends on the magnetic field in the region of weak magnetic fields $\omega_{mm'} < \Gamma_1$ and $\omega_{\mu\mu'} < \Gamma_0$. This phenomenon is analogous to the well-known Hanle effect (magnetic depolarization of radiation).

On passing through the investigated volume, the radiation-power spectrum-increment component $\Delta G_4(\omega)$ of interest to us contains, in accord with formulas (10) and (11), two terms pertaining to the structures of the upper and lower levels. Each term has an extremum in the vicinity of the zero frequencies, with a width on the order of the natural width of the corresponding state, and a system of extrema of the same width in the vicinity of the resonances $\omega = \omega_{mm'}$ and $\omega = \omega_{\mu\mu'}$. The type of the extremum is determined by the sign of N; it is a maximum in an absorbing system and a minimum in an amplifying system.

Thus, the theory predicts that the initial broad "white" fluctuation spectrum of the radiation flux passing through a resonant medium is transformed and acquires extrema whose widths and positions characterize the energy structure of the resonating atomic transition. The Doppler broadening does not limit the accuracy with which this structure is determined by such a method.

The foregoing analysis reveals a kinship between the described phenomenon and the interference beats produced in luminescence or in absorption by a system of atoms in the case of modulated excitation.[3-6] They differ in the fact that in our case the exciting-light intensity spectrum contains not a single harmonic but a continuous set. The analogy is quite far-reaching in the case of [6], in which atoms were optically oriented by modulated light transversely to a magnetic field. When the modulation frequency coincided with the sublevel splitting frequency of the lower state, a decrease was observed in the integral absorption of the orienting light. We point out also the later experiments,[7,8] analogous to those described in [6], in which the excitation was modulated by laser intermode beats.

3. EXPERIMENTAL RESULTS AND DISCUSSION

We investigated experimentally the conversion of the noise spectra of two Xe lines, namely, 3.507μ ($5d[7/2]_3$–$6p[5/2]_2$ transition) and 5.57μ ($5d[7/2]_4$–$6p[5/2]_3$ transition), and also the 3.39μ Ne line ($5s'[1/2]_1$–$4p'[3/2]_2$ transition). The source of the strongly fluctuating radiation was a gas discharge in thin tubes in Xe vapor or in a Ne and He mixture. Under certain conditions, the discharge in these gases produced a strongly pronounced population inversion for the indicated transitions, as a result of which the intrinsic spontaneous radiation was amplified in the discharge. The investigated volume was either a section of the source tube or an independent discharge tube. The experimental conditions were varied over a wide range without noticeably changing the results. Most completely investigated was the 5.57-μ Xe line, the noise of which was investigated with apparatus that ensured the cleanest experimental conditions (Fig. 1). In this apparatus, the tubes of the source and of the working cell, with inside diameter 5 mm and lengths 100 and 15 cm respectively, were separated, had independent electric supplies, and could be filled with gases independently. The tube windows were made of fluorite. To avoid parasitic feedback, the window planes were inclined several degrees away from the normal to the tube axis. An additional unidirectional Faraday cell was placed between the source and the investigated cell, thereby ensuring elimination or reduction of the feedback produced in the system by scattering of the radiation from the windows and from the photoreceiver crystal. The active medium of the Faraday cell was also a discharge in Xe. An axial magnetic field of intensity on the order of 100 Oe ensured 45° rotation of the polarization plane of the radiation over an approximate length of 50 cm. In conjunction with the input polarizer and the output analyzer, which were turned through 45°, the system ensured a difference of not less than one order of magnitude between the transmissions in both directions.

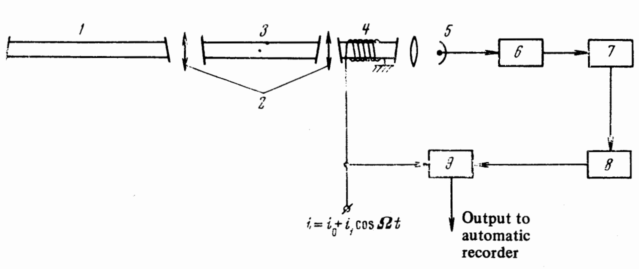

FIG. 1. Block diagram of experimental setup. 1—radiation source, 2—polarizers crossed at an angle 45°, 3—Faraday cell, 4—investigated cell, 5—infrared receiver, 6—tuned amplifier, 7—detector, 8—low-frequency amplifier, 9—synchronous detector.

A dc discharge was excited in all tubes. In the source and in the Faraday cell, the Xe vapor pressure corresponded to saturation at 77°K (2×10^{-3} Torr). The pressure in the cell was varied in the range from 2×10^{-3} to 10^{-1} Torr, making it possible to investigate conditions corresponding to both amplification and absorption of light.[9]

The investigations have shown that the source produced a collimated radiation beam with a power in the range 10^{-5}–10^{-4} W. In this range of powers, the output noise of the photoreceiver (InSb) greatly exceeded the intrinsic noise of the receiver and the noise of the following amplification channel. Direct experiments have shown that the noise was definitely not of the shot type, since equivalent illumination from a thermal source produced no optical noise at all at our sensitivity level. By splitting the beam with a light-splitting plate and feeding the two resultant beams to two different receivers, we were able to prove, by subtracting the signals, that the noises of the two receivers are almost fully correlated, thus confirming the initial assumption that the fluctuations of the photocurrent are similar to the fluctuations of the radiation intensity.

The signal from the photoreceiver was amplified by a tunable resonant amplifier with bandwidth 80 kHz, and was then detected, amplified by a low-frequency amplifier, and then again detected by a synchronous detector. The reference voltage from the latter was the magnetic-field modulation voltage applied axially to the cell. The mean value of the magnetic field was varied slowly with the aid of an automatic device. The magnetic field was swept in the range 0–12 Oe, and the amplitude modulation depth did not exceed 0.05 Oe. Thus, a fixed frequency interval was separated in the apparatus and the dependence of the noise power in this interval on the magnetic field intensity was investigated.

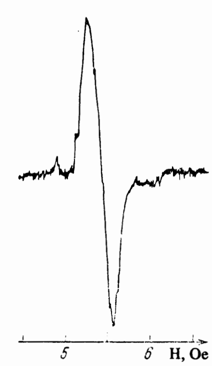

FIG. 2. Derivative of the noise power at 21 MHz vs. the magnetic field intensity (Xe^{136}, λ = 5.57μ).

After preliminary experiments performed with a natural Xe isotope mixture, and then with Xe^{129}, which led to very complicated signals and to complicated dependences of these signals on the experimental conditions,[2)] we settled on the simplest object, namely the even isotope Xe^{136}.

A typical plot of the noise power as a function of the magnetic field intensity is shown in Fig. 2. The receiving section was tuned to 21 MHz, and the recording time constant was 2 sec. An analogous resonance having the same width was observed for a somewhat different field value in the analysis of the 3.507-μ line of Xe. A similar signal was found in the magnetic power spectrum of the 3.39-μ line of Ne. The last signal was much broader

[2)]Results of these experiments for the 3.507-μ line were reported in[1,2]. The interpretation presented there for the signals and their evolution remains in force also in light of the latest results.

than the preceding ones, making it necessary to retune the receiving system to the higher frequency 42 MHz. The signal for Ne was much weaker and was reliably registered only at an accumulation time of 10 sec.

Thus, in all the investigated cases we observed one magnetic resonance of the noise power with a near-Lorentzian contour (the plot represents the derivative of the contour with respect to the magnetic field). To interpret these results, it is useful to elaborate somewhat on formulas (10) and (11), which describe the conversion of the noise spectrum. At a fixed observation frequency, variation of the magnetic field should lead in the general case to observation of two series of resonances, $\omega_{mm'} = \omega$ and $\omega_{\mu\mu'} = \omega$, which are connected with separate interferences of the upper and lower sublevels. Splitting in fields of moderate intensity H is linear:

$$\omega_{mm'} = \Delta m g_1 \mu_B H, \qquad \omega_{\mu\mu'} = \Delta\mu g_0 \mu_B H, \tag{12}$$

where g_1 and g_0 are the corresponding Lande factors, μ_B is the Bohr magneton, and Δm and $\Delta\mu$ are the differences of the angular-momentum projections of the interfering sublevels. For electric dipole transitions, by virtue of the selection rules $m - \mu = \pm 1$ and 0, the projection differences Δm and $\Delta\mu$ assume values 0, ± 1, and ± 2, as can be verified by considering the sets of matrix elements in (10). From this we see, taking (12) into account, that each of the two aforementioned series contains only two[3)] degenerate resonances in the vicinity of the fields $H_1 = \omega/g_1\mu_B$ and $2H_1$ (the series $\omega_{mm'} = \omega$) and analogously $H_0 = \omega/g_0\mu_B$ and $2H_0$ (the series $\omega_{\mu\mu'} = \omega$). In our experiments we realized the particular case of observation along the magnetic field. Optical transitions of the type $m - \mu = \pm 1$ were registered, so that the observed interference was between levels with alternating angular-momentum projections (Δm, $\Delta\mu = \pm 2$), which should lead to only two resonances in the fields $2H_1$ and $2H_0$. In the experiment, however, we observed only one resonance for each spectral line. To explain the absence of the second resonance and to identify the obtained one, it is necessary to calculate the relative amplitudes of the expected resonances. The degeneracy of resonances with equal difference in the indices Δm or $\Delta\mu$ makes it possible to simplify expressions (10) and (11) by summing the coefficients of like frequencies $\omega_{mm'}$ and $\omega_{\mu\mu'}$. The result is the expression

$$\Delta G_4(\omega) = D\left[\frac{A_0^{(1)}\Gamma_1^2}{\Gamma_1^2+\omega^2} + \frac{A_2^{(1)}\Gamma_1^2}{\Gamma_1^2+(\omega_1-\omega)^2} + \frac{A_{-2}^{(1)}\Gamma_1^2}{\Gamma_1^2+(\omega_1+\omega)^2}\right. \tag{13}$$

$$\left. + \frac{A_0^{(0)}\Gamma_0^2}{\Gamma_0^2+\omega^2} + \frac{A_2^{(0)}\Gamma_0^2}{\Gamma_0^2+(\omega_0-\omega)^2} + \frac{A_{-2}^{(0)}\Gamma_0^2}{\Gamma_0^2+(\omega_0+\omega)^2}\right],$$

where

$$A_0^{(1)} = \frac{1}{4\Gamma_1}\begin{Bmatrix} j_1 & 1 & j_0 \\ 1 & j_1 & 2 \end{Bmatrix}^2,$$

$$A_2^{(1)} = A_{-2}^{(1)} = \frac{1}{\Gamma_1}\left[\frac{1}{9(2j_1+1)} + \frac{1}{6}\begin{Bmatrix} j_1 & 1 & j_0 \\ 1 & j_1 & 2 \end{Bmatrix}^2\right]$$

$A_0^{(0)}$ and $A_2^{(0)}$ are obtained from $A_0^{(1)}$ and $A_2^{(1)}$ by making the substitutions $j_1 \rightleftarrows j_0$ and $\Gamma_1 \rightleftarrows \Gamma_0$. j_1 and j_0 denote the total angular momenta of the upper and lower states. The curly brackets contain Wigner's 6j-symbols; D is a proportionality coefficient, $\omega_1 = 2g_1\mu_B H$ and $\omega_0 = 2g_0\mu_B H$.

From the foregoing expression we can readily obtain the ratio of the resonance amplitudes of the sublevels m and μ. If the total angular momenta j_1 and j_0 differ by unity, i.e., $j_1 = j_0 \pm 1$, as occurred in the investigated cases, then it follows from (13) that

$$\frac{A_2^j}{A_2^{j-1}} = \frac{\Gamma_{j-1}}{\Gamma_j}\frac{(j+1)(2j+3)}{(j-1)(2j-3)}. \tag{14}$$

In order to account for both variants $j_1 = j_0 \pm 1$ in a single fashion, we have introduced here new symbols, namely, j is the larger angular momentum of the two states, Γ_j and Γ_{j-1} are the corresponding widths, and A_2^j and A_2^{j-1} are the amplitudes of the resonances due to the interference between the sublevels of the states with angular momenta j and $j - 1$.

For a numerical estimate of the ratio $A_2^{(0)}/A_2^{(1)}$ we use the published data on the widths of the states under conditions close to those prevailing in the experiments (we have in mind the influence of different broadening actions). According to the data of [10], the ratio Γ_1/Γ_0 for both investigated transitions in Xe is of the order of 0.03–0.09, i.e., the lower state $6p[^7\!/_2]_{2,3}$ is much broader than the upper one $5d[^7\!/_2]_{3,4}$. This gives for the 5.57- and 3.507-μ lines estimates of $A_2^{(0)}/A_2^{(1)}$ in the ranges 0.01–0.025 and 0.005–0.017, from which it follows that in both cases the smaller resonance is not observed at the obtained sensitivity level, so that the observed signal should be related to the upper states $5d[^7\!/_2]_{3,4}$. In the case of Ne, the estimate of Γ_1/Γ_0 is even more uncertain. This quantity, however, is apparently no smaller than 0.3,[10-12] leading with the aid of (14) to the conclusion that the ratio $A_2^{(0)}/A_2^{(1)}$ is no smaller than 6, i.e., in the case of Ne the signal should be related to the lower state $4p[^3\!/_2]_2$.

The extrema of the noise power were observed at different frequencies of the receiving section. In all cases, a linear connection between the intensity of the resonant magnetic field and the receiver frequency was confirmed, i.e., linearity of the level splitting in a magnetic field. A special study was made of the sign of the extremum. Direct observation of the detector current during the course of the magnetic-field scanning has shown that, as expected, the signal corresponds to the maximum in the power spectrum under conditions of light amplification in the cell. The change of the signal phase on going from the absorption regime to the amplification regime is demonstrated in Fig. 3.

A singularity in the noise spectrum was observed only in linearly polarized light, regardless of whether the polarizer was placed before or after the cell (the position of the plane of polarization is, of course, likewise immaterial). The vanishing of the magnetic singularity in the noise when unpolarized light is used for the observation follows from an analysis analogous to that given in Sec. 2, in which it is necessary to replace the correlation function $S_\beta(\tau) = \langle J_\beta(t) J_\beta(t+\tau)\rangle$ by the function $S(\tau) = \langle J(t) J(t+\tau)\rangle$, where $J(t) = \Sigma_\beta J_\beta(t)$. Account

[3)]An additional singularity of the spectrum in the vicinity of zero frequencies, corresponding to interference of each sublevel with itself within the limits of its width, Δm, $\Delta\mu = 0$, does not depend on the magnetic field and is not registered with the aid of the described magnetic-scanning procedure; see below.

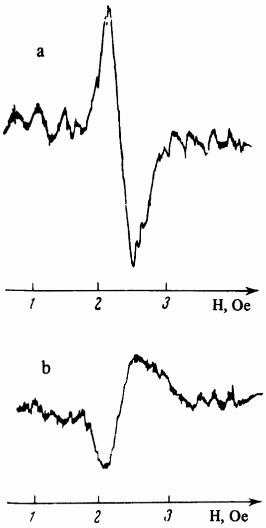

FIG. 3. Inversion of signal on going from an amplifying medium (a) to an absorbing one (b) (Xe^{136}, $\lambda = 5.57\mu$, f = 9.1 MHz).

must be taken here of the fact that the magnetic field is parallel to the beam axis.

As already mentioned, besides the high-frequency extremum of the spectral noise density, which depends on the magnetic field, there should also be observed a low-frequency extremum (which is always a minimum under our conditions), centered at the zero frequency. As shown by a calculation of the amplitudes $A_0^{(1)}$ and $A_0^{(0)}$ in accordance with formula (13), this minimum should be much deeper than the magnetic one. In addition, since its position cannot be controlled, the cell for this extremum is made up of the entire source-plus-cell system. As a result, this dip in the noise spectrum at low frequencies turns out to be very strong and can be observed with the aid of standard apparatus. The impossibility of scanning the radiation spectrum makes the observed photocurrent spectrum a product of the sought spectrum by the spectral characteristic of the photoreceiver, which is not known reliably. At low frequencies, however, the inertia of the receiver can be neglected.

Figure 4 shows curves taken from the screen of the spectrum analyzer for two gas-discharge regimes in Xe. Figure 4a corresponds to maximum amplification and maximum light intensity, and Fig. 4b corresponds to the minimum radiation intensity at which the noise can still be analyzed. Both spectrograms show a deep noise-power dip in the vicinity of zero. This is the sought-for minimum. It is strongly broadened in comparison with the magnetic minimum as a result of saturation. The difference in the degrees of saturation explains the different widths of the dip in cases a and b. The decrease of the spectral noise density in the vicinity of high frequencies is connected mainly with the descending frequency characteristic of the receiver.

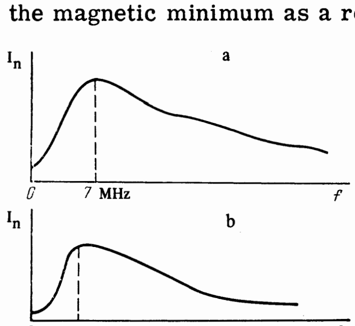

FIG. 4. Noise power spectrum of discharge in Xe^{136}: a—maximum light intensity; b—minimum intensity registered by the receiver.

When comparing the experimental results with the theory, one can speak only of qualitative agreement. In the experiment it was impossible not to violate the initial condition for the validity of the theory, namely that the change of the light intensity on passing through the investigated volume be small. In the experiments, the intensity of the light was changed on passage through the cell by not less than a factor of two. In the experiments with the zero extremum, one cannot even speak of satisfaction of this condition. This is precisely why a quantitative treatment of the results concerning the level widths is difficult. One can speak only of an estimate of the homogeneous width of the sublevels whose interference is observed. The experimentally measured width of the signal for Xe (from extremum to extremum on Fig. 2) is 0.30 Oe or 1.1 MHz. Assuming formula (10) to be valid, this yields $\Gamma = 6 \times 10^6$ sec^{-1}. With increasing light intensity, the broadening due to saturation becomes noticeable.[4)] The presented value corresponds to rather weak light at which the dependence of the signal width on the intensity is not yet observed.

The signal widths for the two xenon lines turned out to be the same, which is natural, since the initial states differ only in the total angular momentum. The width of the analogous extremum in the 3.39μ emission of Ne turned out to be 15 times larger. These results agree qualitatively with the available data on the level widths.[11-12]

Much more definite quantitative results were obtained with respect to the Lande factors of the magnetic splitting of the $5d[^7/_2]_{3,4}$ states of xenon. The table gives the measured values, including data on the odd isotope Xe^{129} for the hyperfine states with total angular momenta $F = ^5/_2$, $^7/_2$, and $^9/_2$, as compared with the calculated values obtained from the measured value of g_j for the even isotope.

State	Xe^{136}	Xe^{129}					
	g_j	$F = ^5/_2$		$F = ^7/_2$		$F = ^9/_2$	
		measurement	calculation	measurement	calculation	measurement	calculation
$5d\ [^7/_2]_3$	1.13	1,27	1.29	0.96	0.97	—	—
$5d\ [^7/_2]_4$	1.36	—	—	1.52	1.51	1,22	1.21

One of the difficulties in the comparison of the experiments with the theory was the spectrum width of the radiation incident on the cell. It is assumed in the theory that the spectrum width is much larger than the investigated splittings in the magnetic field. This condition is fulfilled with difficulty. The Doppler line width, from which the spectrum width at 5.57μ is estimated, is only 60 MHz, with splittings up to 30 MHz investigated. The true width of the spectrum can be either larger or smaller than this value, owing to saturation and regenerative narrowing. In addition, there is no

4)In addition, broadening of the signals was observed when Ne was added to the discharge. At an Ne pressure of about 0.1 Torr, the noise-spectrum line broadened by 1.5 times.

guarantee that the spectrum is perfectly smooth. We attribute to this uncertainty in the emission spectrum a certain nonreproducibility in the shape of the magnetic-extremum line. This shape was sometimes noticeably asymmetrical, with the maximum asymmetry observed at the highest frequencies, in the vicinity of 30 MHz, and at the lowest frequencies, in the region of several MHz. These observations correlate qualitatively with direct observations of the photocurrent spectrum; see Fig. 4. The asymmetry increased with increasing cell length. An example of an asymmetrical signal obtained with a cell 60 cm long is shown in Fig. 5. We note that at such lengths, even the inhomogeneity of the laboratory magnetic field becomes noticeable.

FIG. 5. Signal obtained with a cell 60 cm long (Xe^{136}, $\lambda = 5.57\mu$, f = 26.4 MHz). The registration time constant is 10 sec.

4. CONCLUSION

On the basis of the foregoing we can regard as proved the existence of a nonlinear process of conversion of the fluctuation spectrum of radiation passing through an absorbing or an amplifying medium. The resultant changes of the spectrum make it possible to assess the fine details of the energy structure of the atoms, including those inaccessible to direct spectroscopic investigation, owing to the much larger Doppler broadening. A side result of the investigation is an experimental confirmation of the conclusion that a coherent light amplifier possesses frequency distortion in the vicinity of zero frequencies, with a characteristic width on the order of the homogeneous broadening of the working transition. As applied to the problem of amplification of modulated radiation, this phenomenon was considered theoretically in [13].

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