Quantum features in the noise spectrum in radiation that has passed through a nonequilibrium gaseous medium

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The statistical properties of radiation passing through a nonequilibrium gaseous medium in the transparency band are considered. It is observed that the photocurrent noise spectrum is dispersive in the vicinity of the frequencies ω_0 and $2\omega_0$ (ω_0 is the Zeeman-splitting frequency), and the shot noise is partially suppressed. The dependence of sub-Poisson-statistics effects on the geometry of the experiment is investigated and the possibility of recording these effects is discussed.

1. INTRODUCTION

The fluctuation properties of radiation are known to be greatly altered by interaction with nonlinear optical media.[1,2] In particular, four-wave processes that affect the statistical properties of radiation can be observed in gaseous media. The reported experimental and theoretical research[3–9] in this branch of spectroscopy fluctuations deals mainly with media in which the internal states of the particles are in thermodynamic equilibrium. For magnetic fields with a Zeeman-splitting energy $\hbar\omega_0$ much lower than the temperature T, this corresponds, under ordinary laboratory conditions, to magnetic sublevels with equal atomic populations.

Yet in interactions between radiation and a gas medium with particles whose internal states are not in equilibrium, it is natural to expect the variation of the noise spectrum to be more substantial than under equilibrium conditions. A strongly nonuniform population of the Zeeman levels by atoms, which produces in a gas nonzero average polarization moments (PM) of the density matrix of the atomic ground state, can be achieved, for example in atoms with long polarization-relaxation times, by optical pumping. It was pointed out earlier[10] that intensity-fluctuation spectroscopy methods can be used in this case to record average PM of high rank. It was shown that the photocurrent correlator is determined not only by symmetric correlators (averaged anticommutator) of the PM fluctuations of the Wigner density matrix, but also by antisymmetric (averaged anticommutator) ones. The contribution of the antisymmetric correlator is most substantial in an interaction between radiation and a nonequilibrium gas in the transparency band. Since the noncommutativity of the density-matrix PM fluctuations is evidence that their evolution is basically quantum-mechanical, it is natural to expect that the transmitted radiation will acquire quantum-statistical properties, i.e., the photocurrent noise will become lower than the Poisson noise corresponding to the classical field in definite regions of the spectrum. In this paper we determine the quantum properties of the noise spectrum of radiation that probes a nonequilibrium gas medium in the transparency band.

2. CORRELATOR OF PHOTOCURRENTS PROBING THE TRANSPARENCY BAND

In the quantum theory of photodetection, the photocurrent-noise spectrum is defined in terms of the Fourier transform of the average anticommutator of the quantum operators of the current. The informative part of this anticommutator, which is due to fluctuations of the radiation parameters, is determined by the fourth-order correlation function of the electromagnetic field[1,2,11,12]:

$$\left\langle \frac{1}{2}[\delta i(t_1), \delta i(t_2)]_+ \right\rangle = \frac{e^2c^2}{(2\pi\hbar\omega)^2}\zeta^2 \int\int d^2r_1\, d^2r_2$$

$$\times \langle \tilde{T}(E^{(+)}_{\mu_2}(\mathbf{r}_2t_2)E^{(+)}_{\mu_1}(\mathbf{r}_1t_1))T(E^{(-)}_{\mu_1}(\mathbf{r}_1t_1)E^{(-)}_{\mu_2}(\mathbf{r}_2t_2))\rangle. \quad (1)$$

Here $E^{(\pm)}_{\mu_i}(\mathbf{r}_it_i)$ are the Heisenberg operators of the μ_ith positive- and negative-frequency components of the electromagnetic field at the space–time point $\mathbf{r}_it_i$, ($E^{(\pm)} \sim e^{\pm i\omega t}$); T and $\tilde{T}$ are the time ordering and anti-ordering operators; e is the electron charge; ω is the average frequency of the probing radiation; ζ is the quantum yield of the photodetector. Angle brackets denote averages over the nonequilibrium density operator ρ_0 of the system. The integration is over the photocathode surface.

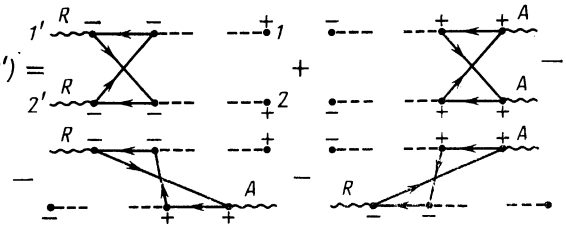
It is convenient to express the positive- and negative-frequency component correlator in expression (1) in terms of a two-photon Green's function,

$$\mathcal{D}(12;2'1') = \langle \tilde{T}(A(2)A(1))T(A(1')A(2'))\rangle, \quad (2)$$

for which there exists a standard diagram in the Keldysh technique[13,10]; $A(i) \equiv A_{\mu\cdot\infty}(\mathbf{r}_it_i)$ are the Heisenberg operators of the vector-potential components. In the quasimonochromatic case the field correlator contained in (1) is obtained from the two-photon Green's function (2) by separating the corresponding frequency components and multiplying by $(\omega/c)^4$.

The two-photon Green's function (2) in the case of Gaussian radiation interacting with a nonequilibrium gas medium was calculated in Ref. 10 in the approximation of an optically thin layer. In the case of a classical electromagnetic field, with account taken of the finite optical thickness of the medium, a similar calculation leads to the following diagrammatic expression:

$\mathcal{D}(12;1'2') =$



(3)

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The straight and wavy lines on these diagrams represent exact atomic and photonic advanced (A) and retarded (R) Green's functions. The dashed lines represent exact classical fields at points in the medium and on the photocathode surface. It is assumed in the derivation of (3) that there are no collisions in the interior of the beam. The probing light is assumed to be weak and not to affect the state of the atom. Equation (2) for the two-photon Green's function was therefore obtained in the lowest order in the field amplitude. No account was taken of the spontaneous emission, which we assume to be small because of the small gathering angle of this radiation. The classical fields contained in (3), obtained at various space-time points by the Huygens principle, are expressed in terms of an incident-wave field specified on some wave surface. By using Eq. (3), we can relate the electromagnetic-field correlator in (1) to the analogous correlator of the incident radiation.

The analytical interpretation of (3) depends substantially on the space group of the average PM of the density matrix of the atoms of the medium. The main difficulty is that the polarization of probing radiation propagating in an optically dense layer can vary, and in the general case the retarded Green's function of a photon in a polarized gas cannot be expressed in simple functional manner in terms of the atomic density matrix.[14] We confine ourselves hereafter to the case when the radiation polarization remains constant and the influence of the finite optical thickness reduces essentially to attenuation of the radiation intensity. This assumption undoubtedly limits the number of conceivable experimental situations. We shall show below, however, that even this approximation can yield important physical results that are absent when an equilibrium gas is probed. This approximation can be realized in experiment, for example, by placing the gas in a magnetic field $\mathbf{H}_0$ perpendicular to the radiation propagation direction. All the PM of the atoms of the medium are then directed along $\mathbf{H}_0$. Assuming also that the atom density matrix has no alignment, a situation attainable by a special optical pumping method, the radiation polarization remains unchanged (see Ref. 14).

The resultant final expression for the photocurrent correlator coincides in fact with the analogous expression for incident radiation having Gaussian statistics, a case considered earlier in Ref. 10. The informative correlator of the transmitted-radiation photocurrents has for incident radiation with Gaussian statistics the same form as for radiation with Poisson statistics because the gas cell influences the noise spectrum in the region of low frequencies Ω: $\Omega \sim \omega_0 \ll \Gamma$, where Γ is the width of the Gaussian-distribution spectrum and determines also the width of the Gaussian increment to the photocurrent noise spectrum. In the calculation that follows we shall assume that both classical and Gaussian statistics are possible for the probing radiation.

We represent the photocurrent correlator in the form

$$\langle {}^1/_2[\delta i(t_1), \delta i(t_2)]_+\rangle = \langle {}^1/_2[\delta i(t_1), \delta i(t_2)]_+\rangle_s + \langle {}^1/_2[\delta i(t_1), \delta i(t_2)]_+\rangle. \quad (4)$$

The first and second terms in the right-hand side of (4) correspond to contributions of the symmetric and antisymmetric correlators of the PM density-matrix fluctuations:

$$K^{\varkappa_1 q_1;\varkappa_2 q_2}_{(s,a)j;j}(\mathbf{P}_1\mathbf{R}_1 t_1;\, \mathbf{P}_2\mathbf{R}_2 t_2) = \left\langle \frac{1}{2}[\delta\hat{\rho}_j^{\varkappa_1 q_1}(\mathbf{P}_1\mathbf{R}_1 t_1), \delta\hat{\rho}_j^{\varkappa_2 q_2}(\mathbf{P}_2\mathbf{R}_2 t_2)]_{(+,-)}\right\rangle. \quad (5)$$

The Heisenberg operator of the fluctuation of the $\varkappa q$-moment of the atom Wigner density matrix is defined as

$$\delta\hat{\rho}_j^{\varkappa q}(\mathbf{PR}t) = \sum_{n',n}(-1)^{j-n'}\Pi_j \begin{pmatrix} j & \varkappa & j \\ -n' & -q & n \end{pmatrix} \times \sum_{\mathbf{p}} e^{i\mathbf{pR}/\hbar} a^+_{n'\mathbf{P}-\mathbf{p}/2}(t)\, a_{n\mathbf{P}+\mathbf{p}/2}(t) - \rho_j^{\varkappa q}(\mathbf{PR}t), \quad (6)$$

where $\rho_j^{\varkappa q}(\mathbf{P}\cdot\mathbf{R}t)$ are the PM density matrices of the atoms; j is the total electron angular momentum and n is its projection; $\mathbf{P}$ is the momentum and $\mathbf{R}$ the coordinate of the atom; $a^+_{\dots}(t)$ and $a_{\dots}(t)$ are the Heisenberg creation and annihilation operators of the atoms whose time evolution is determined by the Hamiltonian of the atomic subsystem. We use the notation[15] $\Pi_{XY\dots} = [(2X+1)(2Y+1)\dots]^{1/2}$.

As shown in Ref. 10, in the case of probing in the absorption band the main contribution to the sum (4) is made by the first term. The qualitative form of the photocurrent noise spectrum is in this case the same as when an equilibrium medium is probed. In the optical transparency region, the main contribution to the sum (4) is made by the second term. The photocurrent correlator then takes the form:

$$\left\langle \frac{1}{2}[\delta i(t_1), \delta i(t_2)]_+\right\rangle \approx 3e^2\zeta^2 \sum_{\varkappa_1,q_1}\sum_{\varkappa_2,q_2}\sum_{K_1,Q_1}\sum_{K_2,Q_2}\iint d^3R_1\, d^3R_2 \times \iint \frac{d^3P_1}{(2\pi\hbar)^3}\frac{d^3P_2}{(2\pi\hbar)^3}\int_{-\infty}^{\infty}\frac{d\omega_1}{2\pi}\frac{d\omega_2}{2\pi}\{S^{00}_{K_1Q_1}(\omega_1 P_{1z};\varkappa_1 q_1) \times A^{00}_{K_2Q_2}(\omega_2 P_{2z};\varkappa_2 q_2) \times \theta(t_1-t_2) - A^{00}_{K_1Q_1}(\omega_1 P_{1z};\varkappa_1 q_1) S^{00}_{K_2Q_2} \times (\omega_2 P_{2z};\varkappa_2 q_2)\theta(t_2-t_1)\}\exp(-l(\omega_1) - l(\omega_2))K^{\varkappa_1 q_1;\varkappa_2 q_2}_{(a)j;j} \times (\mathbf{P}_1\mathbf{R}_1 t_1;\mathbf{P}_2\mathbf{R}_2 t_2)(\hbar\omega)^{-2} I^{K_1Q_1}(\omega_1,\mathbf{R}_{1\perp}) I^{K_2Q_2}(\omega_2,\mathbf{R}_{2\perp}). \quad (7)$$

Here $I^{K_iQ_i}(\omega_i, \mathbf{R}_{i\perp})$ is the spectral density of the polarization tensor of the radiation at the entrance to the cell, and is defined in a coordinate frame with z axis along the propagation direction,

$$I^{K_iQ_i}(\omega_i, R_{i\perp}) = -c\sum_{p,q} C^{K_iQ_i}_{1p1q}(\mathbf{E}^*(\mathbf{R}_{i\perp}))_p \times (\mathbf{E}(\mathbf{R}_{i\perp}))_q\, \delta(\omega_i-\omega), \quad (8)$$

where $(\mathbf{E}(\mathbf{R}_{i\perp}))_q$ are the cyclic components of the complex field amplitude at the point whose coordinate in the plane perpendicular to the z axis is $\mathbf{R}_{i\perp}$. In the case of Gaussian radiation the function $I^{K_iQ_i}(\omega_i, \mathbf{R}_{i\perp})$ is defined as the product of the polarization density matrix and the spectral intensity of the radiation.

The antisymmetric density correlator in a coordinate frame with z' axis along the magnetic field, in the stationary spatially homogeneous state, takes the form

$$K_{(a)j;j}^{\varkappa_1 q_1';\varkappa_2 q_2'}(\mathbf{P}_1\mathbf{R}_1 t_1;\mathbf{P}_2\mathbf{R}_2 t_2)=n_0(2\pi\hbar)^3 f_0(\mathbf{P}_1)\delta(\mathbf{P}_1-\mathbf{P}_2)$$
$$\times\delta\left(\mathbf{R}_2-\mathbf{R}_1-\frac{\mathbf{P}_1}{m}(t_2-t_1)\right)$$
$$\times\sum_{\varkappa}\frac{1-(-1)^{\varkappa_1+\varkappa_2+\varkappa}}{2}(-1)^{2j}\Pi_\varkappa^2\Pi_j$$
$$\times\begin{pmatrix}\varkappa_1 & \varkappa_2 & \varkappa\\ q_1' & q_2' & 0\end{pmatrix}\begin{Bmatrix}\varkappa_1 & \varkappa_2 & \varkappa\\ j & j & j\end{Bmatrix}\exp(-i\omega_0 q_1'(t_1-t_2))\rho_j^{\varkappa 0}, \tag{9}$$

where $f_0(\mathbf{P}_1)$ is the Maxwell function, $\mathbf{n}_0$ is the density of the atoms, and m is the atom mass. Expression (9) corresponds to the limit of a strong magnetic field when $\omega_0\bar{\tau}\gg 1$, where $\bar{\tau}$ is the average flight time. The antisymmetric density correlator in (7) is defined in a coordinate frame with z axis along the propagation direction, and is connected with the correlator (9) by the transformation

$$K_{(a)j;j}^{\varkappa_1 q_1;\varkappa_2 q_2}(\mathbf{P}_1\mathbf{R}_1 t_1;\mathbf{P}_2\mathbf{R}_2 t_2)=\sum_{q_1',q_2'}D_{q_1 q_1'}^{\varkappa_1 *}(\alpha,\beta,\gamma)D_{q_2 q_2'}^{\varkappa_2 *}(\alpha,\beta,\gamma)$$
$$\times K_{(a)j;j}^{\varkappa_1 q_1';\varkappa_2 q_2'}(\mathbf{P}_1\mathbf{R}_1 t_1;\mathbf{P}_2\mathbf{R}_2 t_2), \tag{10}$$

where $D_{qq'}^{\varkappa}(\alpha,\beta,\gamma)$ is the Wigner D function.[15] The Euler angles α, β, and γ characterize the coordinate-frame rotation $(x,y,z)\to(x',y',z')$.

The matrices $S_{KQ}^{00}(\omega P_z;\varkappa q)$ and $A_{KQ}^{00}(\omega P_z;\varkappa q)$ describing the interaction of an atom with a light beam are particular cases of the general expressions (A1–A3) of Ref. 10, and take the form

$$S_{KQ}^{00}(\omega P_z;\varkappa q)=\sigma'(\omega,P_z)\delta_{K,\varkappa}\delta_{Q,-q}(-1)^{j+j'+\varkappa-q}$$
$$\times\frac{\Pi_\varkappa}{\Pi_{j1}}\begin{Bmatrix}j & j & \varkappa\\ 1 & 1 & j'\end{Bmatrix},$$
$$A_{KQ}^{00}(\omega P_z;\varkappa q)=i\sigma''(\omega,P_z)\delta_{K,\varkappa}\delta_{Q,-q}(-1)^{j+j'+\varkappa-q}$$
$$\times\frac{\Pi_\varkappa}{\Pi_{j1}}\begin{Bmatrix}j & j & \varkappa\\ 1 & 1 & j'\end{Bmatrix}, \tag{11}$$
$$\sigma'(\omega,P_z)+i\sigma''(\omega,P_z)=\frac{4\pi k}{\hbar}|d_{jj'}|^2$$
$$\times\left[-i\left(\omega-\omega_{j'j}-k\frac{P_z}{m}\right)+\frac{\gamma_{j'}}{2}\right]^{-1}. \tag{12}$$

Here $k=\omega/c$, $d_{jj'}$ is the reduced dipole-moment matrix element, j' is the total angular momentum of the excited state, P_z is the component of the atom momentum along the wave vector $\mathbf{k}$, $\gamma_{j'}$ is the reciprocal lifetime of the excited state, and $\omega_{j'j}$ is the frequency of the atomic transition.

The optical thickness $l(\omega)$ is defined as

$$l(\omega)=\frac{n_0 L}{3(2j+1)}\int\sigma'(\omega;P_z)f_0(\mathbf{P})\frac{d^3P}{(2\pi\hbar)^3}, \tag{13}$$

where L is the cell length.

If the deviation of the probing radiation from the resonance transition is substantially larger than either the homogeneous or the inhomogeneous line broadening, the contribution of the symmetric density correlator discarded in (7) is of order

$$\gamma_{j'}/|\omega-\omega_{j'j}|\ll 1.$$

We emphasize once more that the symmetric density correlator can be neglected in principle only for a nonequilibrium medium. In an equilibrium gas, the antisymmetric density correlator, and accordingly the right-hand side of (7), is zero.[1)]

3. ANALYSIS OF POSSIBLE EXPERIMENTS

Consider the experiment schematically illustrated in Fig. 1. The probing radiation passes through the polarizer P and through a gas cell with the active-atoms vapor and is incident on a photodetector. The magnetic field $\mathbf{H}_0$ is perpendicular to the propagation direction. In the atomic system we assume all the PM, but not the density matrix, to be aligned (see Sec. 2).

If the probing radiation is linearly polarized at an angle α with the magnetic field, substitution of (9)–(12) in (7) yields for the photocurrent correlator

$$\left\langle\frac{1}{2}[\delta i(t_1),\delta i(t_2)]_+\right\rangle=\frac{5}{2}e^2\zeta^2 J^2 e^{-2l(\omega)}\sigma'(\omega)\sigma''(\omega)(-1)^{2j}$$
$$\times\Pi_j^{-1}\begin{Bmatrix}1 & 1 & 2\\ j & j & j'\end{Bmatrix}^2\sum_{\varkappa}\frac{1-(-1)^{\varkappa}}{2}\rho_j^{\varkappa 0}\Pi_\varkappa^2\begin{pmatrix}2 & 2 & \varkappa\\ 2 & -2 & 0\end{pmatrix}\begin{Bmatrix}2 & 2 & \varkappa\\ j & j & j\end{Bmatrix}$$
$$\times\left[-p^2\sin^2(2\alpha)a_1(t_1-t_2)+\frac{1}{4}(1-p\cos 2\alpha)^2 a_2(t_1-t_2)\right], \tag{14}$$

where J is the average photon-flux density in the incident radiation and p is the degree of linear polarization. The radiation-frequency detuning $\Delta\omega=\omega-\omega_{j'j}$ is assumed to exceed greatly the line and the radiation-spectrum widths. The dependences of the cross sections $\sigma'(\omega,P_z)$ and $\sigma''(\omega,P_z)$ on the atom momentum can then be neglected.

The functions $a_q(\tau)$ determine the time dependence of the photocurrent correlator and take the form

$$a_q(\tau)=n_0\iint d^3R_1\,d^3R_2\left(\frac{m}{2\pi T}\right)^{3/2}\frac{1}{|\tau|^{3/2}}\exp\left(-\frac{m(\mathbf{R}_1-\mathbf{R}_2)^2}{2T\tau^2}\right)$$
$$\times\sin(q\omega_0|\tau|)\mathcal{F}(\mathbf{R}_{1\perp})\mathcal{F}(\mathbf{R}_{2\perp}). \tag{15}$$

The dimensionless function $\mathcal{F}(\mathbf{R}_\perp)$ characterizes the spatial distribution of the intensity in the beam cross section. Without calculating the integral (15) in detail, we note that in the Fourier representation the functions $a_q(\tau)$ correspond to dispersive profiles $a_q(\Omega)$ whose centers are the resonance frequencies $\Omega_q=q\omega_0$ $(q=1,2)$.

A qualitative form of the photocurrent fluctuation spectrum $\langle|\delta i_\Omega|^2\rangle$ with allowance for shot noise is shown in

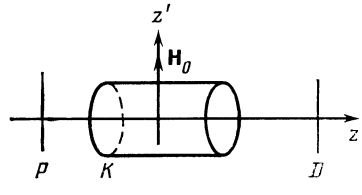

FIG. 1. Experimental setup for the observation of a radiation-fluctuation spectrum; K—gas cell, P—polarizer, D—photodetector, $\mathbf{H}_0$—external magnetic field.

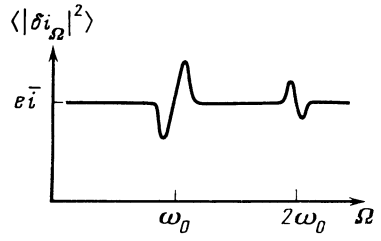

FIG. 2. Qualitative form of the photocurrent fluctuation spectrum in the case of probing in the transparency band; ω_0 is the Zeeman-splitting frequency and $e\bar{i}$ is the shot-noise level.

Fig. 2. The width of the profiles corresponding to the informative part of the photocurrent correlator is determined by the parameter τ_c^{-1} where τ_c is the correlation time characterizing the damping of the function $a_q(\tau)$. In order of magnitude, τ_c is equal to the average collision time $\bar{\tau}\sim a/\bar{v}$, where a is the beam diameter and $\bar{v}$ is the average thermal velocity.

The appearance, in the photocurrent-noise spectrum, of regions in which the fluctuations become smaller than the Poisson level is evidence that the probing radiation has acquired quantum-statistical properties. In our problem the fluctuation spectrum depends strongly on the polarization properties of the radiation. If the polarization and magnetic-field directions coincide ($p=1$, $\alpha=0$) the right-hand side of (14) vanishes. This is physically natural, for in a coordinate frame with z' axis along $\mathbf{H}_0$ the light interacts only with PM fluctuations having $\varkappa=q'=0$. The spectrum of PM fluctuations with $q'=0$ is centered at the zero frequency, in the vicinity of which the contribution of the antisymmetric density correlator is zero.

For the photocurrent correlator we have in the case of circularly polarized radiation

$$\left\langle \frac{1}{2}[\delta i(t_1),\delta i(t_2)]_+ \right\rangle = e^2\xi^2 J^2 e^{-2l(\omega)}\sigma'(\omega)\sigma''(\omega)(-1)^{2j}\Pi_j^{-1} \times \sum_{\varkappa}\frac{1-(-1)^{\varkappa}}{2}$$
$$\rho_j^{\varkappa 0}\Pi_{\varkappa}^2\left[-\frac{3}{2}\begin{pmatrix}1 & 1 & \varkappa\\ 1 & -1 & 0\end{pmatrix}\begin{Bmatrix}1 & 1 & \varkappa\\ j & j & j\end{Bmatrix}\begin{Bmatrix}1 & 1 & 1\\ j & j & j'\end{Bmatrix}^2\xi_2^2 a_1(t_1-t_2) + \frac{5}{8}\begin{pmatrix}2 & 2 & \varkappa\\ 2 & -2 & 0\end{pmatrix}\begin{Bmatrix}2 & 2 & \varkappa\\ j & j & j\end{Bmatrix}\begin{Bmatrix}1 & 1 & 2\\ j & j & j'\end{Bmatrix}^2 a_2(t_1-t_2)\right].$$

The Stokes parameter ξ_2 is indicative of the degree of circular polarization.[16] For circularly polarized radiation, quantum properties appear in the photocurrent noise spectrum even in the case of atoms with total angular momentum $j=1/2$, in view of the orientation fluctuations. In this case alkali atoms constitute a suitable working medium for experiments, owing to their high saturated vapor density at room temperature.

The photocurrent correlator for radiation of arbitrary polarization is obtained by adding to (14) a contribution proportional to ξ_2^2 in (16). Note that in these experimental situations the sign of the informative photocurrent correlator, together with the locations of the regions with less or more shot noise in Fig. 2, is determined by the sign of the projection of odd-rank PM on the magnetic field direction.

4. DISCUSSION AND ESTIMATES

Let us determine the order of magnitude of the observed quantum features in the photocurrent noise spectrum. We consider for simplicity atoms with $j=j'=1/2$. The amplitude of the resonance at the frequency $\Omega_1=\omega_0$ relative to the shot-noise background is characterized by the dimensionless parameter

$$\eta = 3^{1/2}\xi\xi_2^2|\rho^{10}|l(\omega)e^{-l(\omega)}J\sigma''(\omega)\tau_c. \tag{17}$$

Assuming $\xi=\xi_2=1$, $\rho^{10}=3^{-1/2}$ (the maximum value), and $l(\omega)=1$ we obtain

$$\eta = J_{tr}\sigma''(\omega)\tau_c \sim \psi = \bar{J}\sigma''(\omega)\tau_c, \tag{18}$$

where J_{tr} is the photon flux density in the transmitted radiation, and $\bar{J}$ is the average photon-flux density in the cell. The quantity ψ has the physical meaning of the average phase change of the wave function of an atom crossing the probing beam, and $\Delta=\bar{J}\sigma''(\omega)$ describes the shift of the atomic energy level by the nonresonant field. Recognizing that we have neglected in the calculation the reaction of the radiation on the state of the medium, we must regard the parameter ψ, which is the upper bound of η, as a small number.

Note that in the case $\psi\gtrsim 1$ (recognizing that $\sigma'(\omega)\ll\sigma''(\omega)$ in the transparency band) the mean values of the PM of the density matrix of the atoms of the medium can remain unchanged when acted upon by the probing radiation. Indeed, the effects of optical pumping by a probing beam are characterized by a parameter $w=\bar{J}\sigma'(\omega)\tau_c$ which can be small under these conditions. The same parameter corresponds to the relative contribution of the symmetric photocurrent correlator (see Eq. (4)), which we have neglected in the calculation. Increasing ψ up to $\gtrsim 1$ and maintaining $w\ll 1$, we can in principle increase the depth of the dip in the photocurrent noise spectrum. The growth of the dip depth is restricted mainly by the following effects: In the diagrammatic expansion (3), the free-atom Green's functions in the ground state are replaced by Green's functions of atoms "dressed" by the external probing radiation. The phase of the Green's function corresponding to the return line (proportional to the occupation numbers of the atomic states), and the phase of the retarded Green's function that joins these space–time points, acquire different increments, since the atoms interact with radiation localized in space. Finally, at high radiation intensities such that $\psi\gtrsim 1$, the growth of the informative photocurrent correlators saturates due to the appearance of additional oscillating factors that decrease the integrals over the vertices of the diagrams in (3).

We estimate now the real values of the physical parameters at which $\eta\sim\psi\sim 1$. We take Na vapor to be the working substance. We choose the detuning $\Delta\omega$ of the probing-radiation frequency from the resonant atomic frequency of the D_1 line to be of the order of the absorption Doppler width Γ_D. In this case $\sigma'(\omega)\sim 10^{-14}$ cm^2 and $\sigma''(\omega)\sim 3\cdot 10^{-12}$ cm^2. The optical thickness of a cell of length $L\sim 1$ cm reaches $l\sim 1$ at an atom density $n_0\sim 10^{14}$ cm^{-3}. The correlation time for probing by a beam of diameter $a\sim 0.1$ cm is of the order of $\tau_c\sim 10^{-5}$ s. The mean free path of the atoms is $\lambda_0\sim 1$ cm $\gg a$. The photon current density in the beam is $J\sim(\sigma''(\omega)\tau_c)^{-1}\sim 3\cdot 10^{16}$ cm$^{-2}\cdot$s^{-1}, which corresponds to

a radiation intensity $I \sim 10$ mW cm^{-2}. The estimated parameters of the atom concentration and of the radiation intensity are perfectly attainable in experiment.

We conclude by citing one more interpretation of our results, pointed out by I. V. Sokolov. Distinguished among the terms of the diagrammatic expansion (3) of the two-photon Green's function are some having as factors single-photon Green's functions connected with the fluctuation correlators of the quadrature components of the transmitted radiation. The sub-Poisson statistics of the photocurrent noise, observed in the present study, are a consequence of the transition of the electromagnetic field emerging from the gas cell into a squeezed state. The probing radiation then acts simultaneously as a pump wave in the four-wave process that ensures an inhomogeneous evolution of the quadrature components, and as a heterodyne wave.

[1)]We neglect the small Boltzmann factor $\hbar\omega_0/T$, which determines the nonuniformity of the population of the Zeeman sublevels in the equilibrium case.

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