

Spin noise of exciton polaritons in microcavities

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We develop a theory of spin fluctuations of exciton polaritons in a semiconductor microcavity under nonresonant unpolarized pumping. It is shown that the corresponding spin noise is sensitive to the scattering rates in the system, occupation of the ground state, statistics of polaritons, and interactions. The spin noise spectrum drastically narrows in the polariton lasing regime due to formation of a polariton condensate, while its shape can become non-Lorentzian owing to interaction-induced spin decoherence.

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Introduction. Quantum microcavity is a system where a semiconductor quantum well is placed between the Bragg mirrors making it possible to achieve a strong coupling between the light and matter. In this system the energy is coherently transferred back and forth between the photon trapped in the microcavity and the exciton, the elementary excitation of the semiconductor. First observed by Weisbuch *et al.*,¹ the strong coupling results in the formation of mixed light-matter quasiparticles, exciton polaritons, extensively studied since then.^{2,3} The exciton polaritons combine an extremely small effective mass, inherited from the photon, and strong interactions between themselves and with the environment due to the excitonic fraction in this quasiparticle. These quasiparticles are at the origin of several fascinating phenomena, which primarily include polariton lasing;⁴ macroscopic accumulation of polaritons in a single quantum state (nonequilibrium condensate) accompanied by spontaneous emission of coherent light by this state.

In semiconductor microcavities exciton polaritons are characterized by the spin projections $s_z = \pm 1$ onto the structure growth axis z corresponding to the right or left circular polarization of the photon and to the same spin component of the exciton. Superpositions of $s_z = \pm 1$ states give rise to the linear or elliptical polarization of exciton polaritons. Polariton lasers represent a model bosonic system where the polarization of light emitted by a microcavity directly corresponds to the spin state of quasiparticles. This allows one to study the polariton spin dynamics by optical methods.⁵ A great amount of prominent spin-related effects has been realized in microcavities, including self-induced Larmor precession,^{6,7} linear polarization inversion,⁸ optical spin Hall effect,^{9–11} spin Meissner effect,^{12–14} spin multistability,^{15,16} etc. (see Refs. 2 and 17 for reviews). Hence, quantum microcavities became a solid state playground to study interacting and, generally, nonequilibrium Bose systems.

In polariton lasers, the spontaneous symmetry breaking results in the appearance of a stochastic vector polarization including circular polarization related to the spin of a polariton condensate.^{18,19} The Stokes vector (pseudospin) of a condensate may be pinned to one of the crystal axes,²⁰ due to the structure anisotropy. In this case the time-averaged

polarization of emission of polariton lasers is defined by pinning, while its momentary value fluctuates. The spin fluctuations give rise to the “spin noise” which may be studied by optical measurements. The spin fluctuations, being inherent to any system at equilibrium or not, were first observed in 1981 in Na vapor.²¹ The studies of spin fluctuations have become an important part of spintronics. The measurements of spin noise provide crucial information on the spin dynamics of carriers, excitons, and nuclei, including spin precession and relaxation rates and statistics of spin fluctuations. These characteristics are hardly accessible otherwise.^{22–27}

Here we study theoretically the spin noise spectra of polariton lasers both below and above the laser threshold and demonstrate that they are extremely sensitive to the mean occupation number of the condensate and to the statistics of polaritons. A similar approach can be used for a variety of systems where condensates of spin polarized bosons interact with an incoherent reservoir. An appropriate example is given by magnon condensates in crystals, either strongly driven by periodic fields²⁸ or imposed to sufficiently strong static magnetic fields at a low temperature.²⁹ Here the noise is seen as fluctuations in the sample magnetization. The finite-temperature Bose-Einstein condensates of atoms in optical traps with synthetic spin-orbit coupling represent another important example of the direct applicability of our approach.³⁰

Model. The dynamics of the exciton-polariton spin doublet can be described with the pseudospin approach where the density matrix of polariton state with the wave vector \mathbf{k} can be written as $\hat{\rho}_{\mathbf{k}} = N_{\mathbf{k}} \hat{I} + \mathbf{S}_{\mathbf{k}} \cdot \hat{\boldsymbol{\sigma}}$. Here \hat{I} is the unit 2×2 matrix, $\hat{\boldsymbol{\sigma}}$ is a pseudovector composed of Pauli matrices, $N_{\mathbf{k}}$ is the spin-average occupancy of the state \mathbf{k} , and $\mathbf{S}_{\mathbf{k}}$ is the pseudospin of polaritons in this state. In what follows we focus on the dynamics of the ground state, corresponding to $\mathbf{k} = 0$, treat all other states as a reservoir, and omit the \mathbf{k} subscript in the notations. The pseudospin components S_{α} , where $\alpha = \{x, y, z\}$ is the Cartesian index, corresponds to the polarization of emission: S_z/N gives its circular polarization degree, while S_x/N and S_y/N give the linear polarization degree in the axes frames (xy) and $(x'y')$ rotated by 45° with respect to each other.² The dynamics of the occupation number and

the polariton pseudospin is governed by the set of kinetic equations¹⁷

$$\frac{dN}{dt} + \frac{N}{\tau_0} + \mathcal{Q}_n\{N, \mathbf{S}\} = 0, \quad (1)$$

$$\frac{d\mathbf{S}}{dt} + \mathbf{S} \times \boldsymbol{\Omega} + \frac{\mathbf{S}}{\tau_0} + \mathcal{Q}_s\{\mathbf{S}, N\} = 0, \quad (2)$$

where τ_0 is the lifetime of polaritons in the ground state and $\boldsymbol{\Omega}$ is the effective magnetic field arising from anisotropy of the system and from polariton-polariton interactions; its explicit form will be given below. The scalar, $\mathcal{Q}_n\{N, \mathbf{S}\}$, and vector, $\mathcal{Q}_s\{\mathbf{S}, N\}$, collision integrals describe the arrival and departure of the particles into the ground state (see Refs. 17 and 6 for general expressions). Here we adopt their simplest form:

$$\mathcal{Q}_n\{N, \mathbf{S}\} = W^{\text{out}}N - W^{\text{in}}(1 + N), \quad (3)$$

$$\mathcal{Q}_s\{\mathbf{S}, N\} = (W^{\text{out}} - W^{\text{in}} + \tau_s^{-1})[\mathbf{S} - \mathbf{S}_0(\boldsymbol{\Omega})]. \quad (4)$$

Here W^{out} and W^{in} are out- and in-scattering rates, related to the presence of the reservoir, as schematically illustrated in Fig. 1(a). In particular, W^{in} is proportional to the occupation of the reservoir and it is determined by the pumping rate. The term proportional to $1 + N$ describes the stimulated transitions due to the bosonic statistics of quasiparticles. In the collision integral for the polariton spin, $\mathcal{Q}_s\{\mathbf{S}, N\}$, τ_s is the spin-relaxation time and $\mathbf{S}_0(\boldsymbol{\Omega})$ is the steady-state spin induced by the effective field $\boldsymbol{\Omega}$. If polariton exchange with reservoir is efficient, one can introduce the effective temperature T of the polariton system. The steady-state spin is $-\langle N \rangle \boldsymbol{\Omega} / |\boldsymbol{\Omega}|$ for $T \ll \Omega$ and 0 for $T \gg \Omega$. In what follows, unless otherwise specified, we consider the case of high temperatures, where $\langle \mathbf{S} \rangle \equiv 0$, and the occupancy of the polariton ground state

given by the balance of in- and out-scattering processes is

$$\langle N \rangle = \frac{W^{\text{in}}}{\tau_0^{-1} + W^{\text{out}} - W^{\text{in}}}, \quad (5)$$

where the condition $W^{\text{in}} < \tau_0^{-1} + W^{\text{out}}$ should hold.³¹ The straightforward generalization of our approach to account for the entire ensemble of polaritons following Refs. 17 and 6 using the full density matrix is not expected to yield qualitatively different results.

The fluctuations of the condensate occupation number $\delta N(t) \equiv N(t) - \langle N \rangle$ and pseudospin $\delta \mathbf{S}(t) \equiv \mathbf{S}(t) - \langle \mathbf{S} \rangle$ are described by the correlation functions, namely, $\mathcal{K}(t) \equiv \langle \delta N(t') \delta N(t' + t) \rangle$ and $C_{\alpha\beta}(t) \equiv \langle S_\alpha(t') S_\beta(t' + t) \rangle$, where the angular brackets denote the averaging over the time t' for a given shift t . According to the general theory of fluctuations^{32–34} the correlation functions obey the same set of kinetic equations for t or t' dependence as fluctuating quantities.³⁵ The solution of Eqs. (1) and (2) in the absence of effective magnetic fields and interactions ($\boldsymbol{\Omega} \equiv 0$) results in the exponential time decay of correlations and isotropic spin fluctuations:

$$\mathcal{K}(t) = \mathcal{K}(0)e^{-|t|/\tau_c}, \quad C_{\alpha\beta}(t) = \delta_{\alpha\beta} C_{\alpha\alpha}(0)e^{-|t|/\tau_{c,s}}, \quad (6)$$

where $\delta_{\alpha\beta}$ is the Kronecker δ symbol, single time correlators (mean-square fluctuations) $\mathcal{K}(0) = \langle (\delta N)^2 \rangle$, $C_{\alpha\alpha}(0) = \langle S_\alpha^2 \rangle_0$ will be found below, while the particle-spin correlations $\langle \delta N(t') S_\alpha(t' + t) \rangle$ vanish and will be disregarded. We introduced the correlation times τ_c and $\tau_{c,s}$ according to

$$\frac{1}{\tau_c} = \frac{1}{\tau_0} + W^{\text{out}} - W^{\text{in}} = \frac{\tau_0^{-1} + W^{\text{out}}}{1 + \langle N \rangle}, \quad \frac{1}{\tau_{c,s}} = \frac{1}{\tau_c} + \frac{1}{\tau_s}. \quad (7)$$

Equations (6) and (7) clearly show that the particle number and spin fluctuations of exciton polaritons decay exponentially with time and the correlation time of the particle number fluctuations is τ_c , while the spin fluctuations vanish faster, at $\tau_{c,s} < \tau_c$. The spin fluctuation spectra defined as $C_{\alpha\beta}^{(\omega)} \equiv \int_{-\infty}^{\infty} C_{\alpha\beta}(t) e^{i\omega t} dt$ are Lorentzian:

$$C_{\alpha\alpha}^{(\omega)} = C_{\alpha\alpha}(0) \frac{2\tau_{c,s}}{1 + \omega^2 \tau_{c,s}^2}, \quad (8)$$

with the half-width at half maximum $\tau_{c,s}^{-1}$ determined by the inverse spin correlation time. It follows from Eq. (7) that the correlation time of fluctuations is strongly enhanced and the spin noise spectrum is strongly narrowed, if $\langle N \rangle \gg 1$, i.e., where the ground state is macroscopically occupied. In particular, if $\tau_s \rightarrow \infty$ and $W^{\text{out}} \rightarrow 0$, which corresponds to the negligible spin flip in the ground state and negligible depletion of the condensate due to the scattering of polaritons back to the reservoir, we have for the correlation time of particle and spin fluctuations $\tau_c = \tau_0(1 + \langle N \rangle)$. The lifetime of exciton polaritons in the state-of-the-art structures varies from ~ 1 ps to ~ 100 ps with the corresponding maximum $\langle N \rangle$ being between 10^3 and 10^5 , yielding τ_c in the range of 1 ns–10 μ s. Hence, the spin noise frequencies range from MHz to GHz for the macroscopically occupied ground state. The typical spin noise spectra and temporal dependence of the correlators in Figs. 1(b) and 1(c) make the drastic effect of the ground state

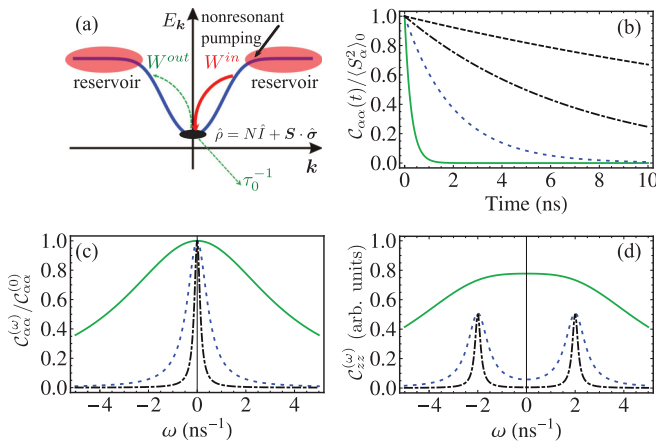


FIG. 1. (Color online) (a) Illustrative scheme of the pumping. The reservoir and the ground state are shown as well as the in- and out-scattering processes. (b) Temporal dependence of spin fluctuations calculated for $\langle N \rangle = 10$ [green (solid)], 100 [blue (dotted)], and 1000 [black (dash-dotted)]. Other parameters are $\tau_0 = 25$ ps, $W^{\text{out}} = 0$, and $\tau_s = 10$ ns. The dashed curve is calculated for $\tau_s \rightarrow \infty$ and $\langle N \rangle = 1000$. (c) Spin noise power spectra calculated for the same parameters. (d) Spin noise power spectra calculated with allowance for the anisotropic splitting $\Omega_a = 2$ ns⁻¹. Other parameters are the same as in panels (b) and (c).

occupation clear. The several orders of magnitude enhancement of the spin correlation time and narrowing of the noise spectrum is a general bosonic effect. Indeed, any fluctuation is amplified by the bosonic stimulation, factor $1 + N$ in Eq. (3).

The mean square of the particle and pseudospin fluctuations can be found using the master equation approach.³ In the absence of interactions and effective magnetic fields, the system is spin isotropic and can be described by the independent occupations of two orthogonal spin states, N_\uparrow and N_\downarrow , described by the same distribution functions $P(N_{\uparrow,\downarrow})$. Using these distribution functions one can express the mean-square fluctuations, that is, $t = 0$ correlators $\langle(\delta N)^2\rangle = \sum_{N_\uparrow, N_\downarrow} P(N_\uparrow)P(N_\downarrow)[(N_\uparrow + N_\downarrow)/2 - \langle N \rangle]^2$ and $\langle S_\alpha^2 \rangle_0 = \sum_{N_\uparrow, N_\downarrow} P(N_\uparrow)P(N_\downarrow)[(N_\uparrow - N_\downarrow)/2]^2$, as

$$\langle S_\alpha^2 \rangle_0 = \langle(\delta N)^2\rangle = \frac{1}{2}\langle N \rangle[1 + (g^{(2)} - 1)\langle N \rangle], \quad (9)$$

where $\langle S_\alpha^2 \rangle_0$ corresponds to isotropic fluctuations, and $g^{(2)}$ is the second-order coherence of a single state. In particular, $g^{(2)} = 2$ corresponds to the thermal statistics,³⁶ where particle and spin square fluctuations are $\propto \langle N \rangle(1 + \langle N \rangle)$ and grow quadratically with the ground state occupation. This situation is realized for equilibrium Bose gas³⁷ or if the ground state feedback on reservoir is negligible. By contrast, $g^{(2)} = 1$ corresponds to the coherent statistics, where the mean-square fluctuations are suppressed, being $\propto \langle N \rangle$. In the limit of low occupancy, $\langle(\delta N)^2\rangle = \langle N \rangle$ in agreement with the theory of classical gas. Moreover, the statistics of spin fluctuations can be determined, as a convolution of distribution functions $P(N - S_\alpha)$ and $P(N + S_\alpha)$. In the limit $\langle N \rangle \gg 1$ we obtain

$$p_{\text{coh}}(S_\alpha) = (\pi \langle N \rangle)^{-1/2} \exp(-S_\alpha^2/\langle N \rangle), \quad g^{(2)} = 1, \quad (10)$$

$$p_{\text{th}}(S_\alpha) = \langle N \rangle^{-1} \exp(-2|S_\alpha|/\langle N \rangle), \quad g^{(2)} = 2. \quad (11)$$

The full statistics of polariton condensates can be determined by numerical integration of Langevin equations for the condensate wave functions, as was done, e.g., in Ref. 19 for stochastic polarization under pulsed excitation. Here we resort to an analytical approach based on kinetic equations treating $g^{(2)}$ phenomenologically.

Role of effective magnetic fields. We begin the discussion of the effective magnetic fields with the case of an anisotropic system where the polariton doublet is split into a pair of states, linearly polarized along the x and y axes. In such a case the vector $\mathbf{\Omega} = (\Omega_a, 0, 0)$ determines the anisotropic splitting,^{6,38–40} and we assume that the effective temperature T of the system exceeds the anisotropic splitting to neglect the steady spin polarization. The S_z and S_y rotate, while the dynamics of S_x remains purely dissipational. As a result, the fluctuations become anisotropic with nonzero spectrum components (cf. Ref. 41)

$$\frac{C_{xx}^{(\omega)}}{\langle S_\alpha^2 \rangle_0} = \frac{2\tau_{c,s}}{1 + \omega^2\tau_{c,s}^2}, \quad \frac{C_{yy}^{(\omega)}}{\langle S_\alpha^2 \rangle_0} = \sum_{\pm} \frac{\tau_{c,s}}{1 + (\omega \pm \Omega_a)^2\tau_{c,s}^2}, \quad (12)$$

$$C_{yz}^{(\omega)} = \frac{2i\omega\Omega_a\tau_{c,s}^2}{1 + \tau_{c,s}^2(\omega^2 + \Omega_a^2)} C_{yy}^{(\omega)}, \quad (13)$$

and $C_{zz}^{(\omega)} = C_{yy}^{(\omega)}$, $C_{yz}^{(\omega)} = [C_{zy}^{(\omega)}]^*$, $\langle S_\alpha^2 \rangle_0$ is given by Eq. (9), and $\tau_{c,s}$ is given by Eq. (7). Figure 1(d) shows the spin noise spectra calculated with the allowance for the anisotropic splitting for different occupations of the ground state. It is seen that the single peak is transformed into a two-peak structure even for a very small value of $\hbar\Omega_a$ taken in our calculation. This is because the spin correlation time $\tau_{c,s}$ depends strongly on the ground state occupation: For small pumping rates and small $\langle N \rangle$ the product $\Omega_a\tau_{c,s} \ll 1$ and the splitting is not visible, however, for larger pumping $\Omega_a\tau_{c,s}$ becomes comparable or larger than 1, making the anisotropic splitting $\hbar\Omega_a$ resolvable, as demonstrated in Fig. 1(d). If the temperature is so low that $T \ll \hbar\Omega_a$, the polaritons are predominantly polarized along the effective field direction x . In this case, $\langle S_x \rangle = -\langle N \rangle$, $\langle S_y \rangle = \langle S_z \rangle = 0$, while the mean-square fluctuations take the form $\langle(S_x - \langle S_x \rangle)^2\rangle = \langle(\delta N)^2\rangle$, $\langle S_y^2 \rangle = \langle S_z^2 \rangle = \langle N \rangle/2$. Moreover, the nontrivial single time correlation appears: $\langle S_y S_z \rangle = \langle S_z S_y \rangle^* = -i\langle N \rangle/2$. The spin noise spectrum in this case can be obtained in a similar way. Equation (12) holds here, albeit with different mean-square fluctuations found above.

Now let us consider the effect of polariton-polariton interactions on the spin noise. As follows from multiple experimental and theoretical works^{8,12–14,16,42} these interactions are strongly spin anisotropic: The exciton polaritons with the same z pseudospin components, i.e., with the same circular polarizations, repel each other efficiently due to the exchange interaction of electrons/holes with the same spin, while the polaritons with opposite circular polarizations can weakly attract each other; the latter is neglected. Hence, the interactions create an effective fluctuating field $\mathbf{\Omega} = (0, 0, \Omega_i)$, directed along the z axis. Its magnitude $\hbar\Omega_i \equiv \alpha_1 S_z$ ($\alpha_1 > 0$ is the interaction constant) is related to the fluctuations of the polaritons pseudospin z component. This effective field has two important effects on polariton spin dynamics and spin noise: (i) it induces precession of the pseudospin around the z axis, known as self-induced Larmor precession,^{6,7} and (ii) it suppresses fluctuations of the z pseudospin component, favoring linear polarization of the macro-occupied state.^{12,43}

It is instructive to start the analysis with the spin precession effect assuming that the effective temperature is high enough, $T \gg \alpha_1 \langle S_z^2 \rangle_0^{1/2}$. In addition, we neglect below the anisotropic splitting ($\Omega_a = 0$). Clearly, the effective field Ω_i induces dephasing of the in-plane pseudospin components. It follows from Eq. (2) that

$$\frac{dC_{\alpha\beta}(t)}{dt} + \frac{C_{\alpha\beta}(t)}{\tau_{c,s}} \pm \frac{\alpha_1}{\hbar} C_{\alpha\beta}(t) S_z(t) = 0. \quad (14)$$

Here the upper (lower) sign corresponds to $\alpha = x$ (y), correlations between z and in-plane components are disregarded, since the field $\mathbf{\Omega}$ does not couple them, and S_z can be considered as an independent parameter whose fluctuations are given by Eqs. (6) and (8). The set of linear equations (14) can be readily solved as

$$\frac{C_{\alpha\alpha}(t)}{\langle S_\alpha^2 \rangle_0} = e^{-|t|/\tau_{c,s}} \left\langle \exp \left[i \frac{\alpha_1}{\hbar} \int_0^{|t|} S_z(t_1) dt_1 \right] \right\rangle_z, \quad (15)$$

with $\langle \dots \rangle_z$ meaning the averaging over the fluctuations of S_z . The treatment of the general case is beyond the scope of this

work. Here we consider two limiting cases: the regimes of fast and slow fluctuations, respectively.

If $\alpha_1 \langle S_z^2 \rangle_0^{1/2} \tau_{c,s} / \hbar \ll 1$, the interaction-induced effective field changes much faster than the pseudospin rotates. This case corresponds to the motional narrowing and

$$\frac{C_{\alpha\alpha}(t)}{\langle S_{\alpha}^2 \rangle_0} = \exp\left(-\frac{|t|}{\tau_{c,s}} - \frac{\alpha_1^2 \langle S_{\alpha}^2 \rangle_0 \tau_{c,s} |t|}{\hbar^2}\right). \quad (16)$$

In this limit the spin fluctuations decay exponentially, resulting in the Lorentzian spectrum of spin noise. This spectrum is, however, anisotropic: The width of $C_{xx}^{(\omega)}$ and $C_{yy}^{(\omega)}$ is larger than that of $C_{zz}^{(\omega)}$ because the interaction-induced field $\mathbf{\Omega} \parallel z$ does not affect S_z .

In the opposite limit, where the fluctuations of the spin z component are slow enough and the in-plane spin components make several oscillations during the correlation time of S_z , i.e., $\alpha_1 \langle S_z^2 \rangle_0^{1/2} \tau_{c,s} / \hbar \gg 1$, the fluctuations of S_z and of the effective field Ω_i can be assumed frozen. In this case, the spin dephasing takes place on the time scale of Ω_i^{-1} . The particular t dependence of the correlator $C_{\alpha\alpha}(t)$ is determined by statistics of the condensate.⁴⁴ For $\langle N \rangle \gg 1$ we obtain for two important limiting cases of coherent and thermal statistics ($t \ll \tau_{c,s}$),

$$\frac{C_{\alpha\alpha}(t)}{\langle S_{\alpha}^2 \rangle_0} = \exp(-\Gamma_{\text{coh}}^2 t^2), \quad g^{(2)} = 1, \quad (17)$$

$$\frac{C_{\alpha\alpha}(t)}{\langle S_{\alpha}^2 \rangle_0} = \frac{1}{1 + \Gamma_{\text{th}}^2 t^2}, \quad g^{(2)} = 2, \quad (18)$$

where the dephasing rates are $\Gamma_{\text{coh}}^2 = \alpha_1^2 \langle N \rangle / 4\hbar^2$ and $\Gamma_{\text{th}}^2 = \alpha_1^2 \langle N \rangle^2 / 4\hbar^2 = \langle N \rangle \Gamma_{\text{coh}}^2$. In this limit, the temporal dependence of the spin fluctuations is directly related to the ground state statistics: Gaussian fluctuations of S_z described by $p_{\text{coh}}(S_z)$ in Eq. (10) result in the Gaussian decay of the in-plane spin components, while the sharply peaked $p_{\text{th}}(S_z)$ in Eq. (11) for thermal statistics results in the slow power-law decay of the fluctuations due to high probability of small S_z , corresponding to low precession rates. As a result, the noise spectrum of the in-plane pseudospin components deviates strongly from Lorentzian. Fourier transform of Eqs. (17) and (18) yields

$$\frac{C_{\alpha\alpha}^{(\omega)}}{\langle S_{\alpha}^2 \rangle_0} = \frac{\sqrt{\pi}}{\Gamma_{\text{coh}}} \exp(-\omega^2 / 4\Gamma_{\text{coh}}^2), \quad g^{(2)} = 1, \quad (19)$$

$$\frac{C_{\alpha\alpha}^{(\omega)}}{\langle S_{\alpha}^2 \rangle_0} = \frac{\pi}{\Gamma_{\text{th}}} \exp(-|\omega| / \Gamma_{\text{th}}), \quad g^{(2)} = 2. \quad (20)$$

The calculated nonanalytical dependence of the spin fluctuations for the thermal statistics at $\omega \rightarrow 0$ is related to the power-law temporal decay of the fluctuations. The typical temporal dependence of the in-plane pseudospin correlators

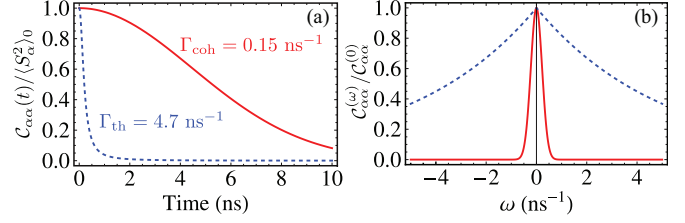


FIG. 2. (Color online) (a) Temporal dependence of the in-plane pseudospin fluctuations calculated with allowance for interactions according to Eqs. (17) and (18) for $\langle N \rangle = 1000$, $\alpha_1 = 10^{-2} \text{ ns}^{-1}$ (Ref. 42) for coherent statistics [red (solid)] and thermal statistics [blue (dotted)]. (b) Spin noise power spectra calculated for the same parameters as in (a).

and noise power spectra are plotted in Figs. 2(a) and 2(b), respectively. The figure clearly shows different qualitative behavior of the noise of the in-plane pseudospin components for different statistics of the polaritons. We stress that the drastic difference in the decoherence times and noise spectral widths for coherent and thermal statistics is related to the different dependencies of $\langle S_z^2 \rangle$ on the ground state occupancy $\langle N \rangle$ [see Eq. (9)].

Note, that if $\alpha_1 \langle S_z^2 \rangle$ is comparable with or larger than the effective temperature T of the system, the fluctuations of the pseudospin z component would be suppressed. This effect can be modeled by multiplying the distribution function of S_z in Eqs. (10) and (11) by the Boltzmann factor $\exp(-\hbar\alpha_1 S_z^2 / T)$ for the probability of thermal fluctuations.¹² As a result, in the limit of $T \rightarrow 0$ we obtain $\langle S_z^2 \rangle = T / 2\hbar\alpha_1$. In this case, the dephasing rate, which determines the spin noise spectral width, can be estimated as $\sim \sqrt{\alpha_1 T / \hbar}$.

Conclusions. We have developed an analytical theory of spin fluctuations of polaritons in microcavities in the lasing regime and demonstrated that the spin noise spectra, being a fundamental property of any spin system, qualitatively depend on the occupation numbers, statistics, and interactions between the particles. Various regimes of spin noise have been identified. Thus, spin noise spectroscopy allows one to study a large variety of spin-related properties of bosonic systems in a single experiment. Experimental verification of these predictions can be done by Fourier spectroscopy of Kerr or Faraday rotation. Extension of this model to other spin-polarized bosonic systems is straightforward.

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