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Stochastically-Induced Quantum Interference in Coherently Driven Two-Level Atoms

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Abstract. It is shown that stochastic perturbation can create phase correlation in a coherently driven two-level system. This correlation results in a noise-induced interference in the two-level system. One possible manifestation of the described interference is a narrow structure in resonance fluorescence spectra. Theoretical interpretation is well confirmed by preliminary experimental results.

Keywords: quantum interference, resonance fluorescence, phase noise
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1. Introduction

Stochastic noise in general destroys coherence in physical systems and consequently eliminates signatures of interference phenomena. Simple familiar examples are the Young's double-slit interference experiment in classical optics and the quantum beat or level-crossing experiments in atomic physics. In the Young double-slit experiment, the contrast of interference light fringes decreases if the relative phase of interfering beams is affected by noise, e.g. associated with a limited coherence of light. The atomic analogue of the Young experiment is the scattering of light on a system possessing superposition states, e.g. when an atom initially in ground state g is excited to a superposition of states e and e' and decays to state g'. Two kinds of scattering experiments can be performed with such atomic systems. One is when levels e and e' are nondegenerate and their superposition is prepared by appropri-

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ately short light pulse. In the decay following the pulsed excitation there are two components of frequencies $\omega_{e'g}$ and ω_{eg}, which beat together and produce an oscillation at the frequency $\omega_{e'e}$, the so-called *quantum beats* [1]. Quantum beats signify quantum interference between indistinguishable channels of the light scattering between levels g and g' via the intermediate states e, e' that are the counterparts of two slits in the Young experiment. Such oscillations are washed out if some phase noise, e.g. due to collisions, perturbs the phases of the atomic dipoles associated with levels e and e'. Another kind of experiment can be performed with a CW excitation by changing the spacing between levels e and e'. This can be easily accomplished if e and e' are Zeeman sublevels affected by a magnetic field. At a given value of the external magnetic field, levels e and e' may cross, i.e. become energetically degenerate, $\omega_{e'e} = 0$. Again, degenerate states e, e' become counterparts of two slits in the Young experiment [2]. Such a *level crossing* of excited-state sublevels results in stationary interference contribution to the resonance fluorescence. Note that such level crossings always occur at zero magnetic field, which is known as the Hanle effect [3]. Similarly as in the case of quantum beats, the level-crossing resonance broadens and disappears if a phase noise perturbs the free evolution of the atomic dipoles.

Apart from numerous situations where noise destroys the interference, there are cases when noise creates interference. The most important example is the *pressure-induced resonance* which, like quantum beats and level-crossing resonances, also occurs in three-level systems [4].

In this paper we demonstrate that stochastic perturbation can create quantum interference also in a simpler, two-level system. We describe results of our theoretical and experimental study of collisionally perturbed resonance fluorescence (RF) and attribute narrow features of the RF spectra to the noise-induced interference.

2. Theoretical Model

2.1. Atomic dynamics

We use the model of a two-level atom with states $|g\rangle$ and $|e\rangle$, respectively, driven by a laser field of frequency ω_L and subjected to a stochastic perturbation. We assume that the stochastic perturbation affects atomic resonance frequency, $\omega_a(t) = \omega_{a0} + \delta\omega(t)$, with $\langle \delta\omega(t)\delta\omega(t')\rangle = 2\Gamma\delta(t-t')$. This class of noise corresponds to many realistic situations, e.g. elastic, dephasing collisions.

In the unperturbed atomic basis $\{|e\rangle, |g\rangle\}$ the Hamiltonian of the system has the form

$$\hat{H}_{AL} = \hbar[\omega_L - \omega_a(t)]\hat{S}_z + \frac{\hbar}{2}\Omega(\hat{S}^+ + \hat{S}^-)\,, \tag{1}$$

with

$$\hat{S}_z = \begin{pmatrix} \frac{1}{2} & 0 \\ 0 & \frac{1}{2} \end{pmatrix}, \quad \hat{S}^+ = \begin{pmatrix} 0 & 0 \\ 1 & 0 \end{pmatrix}, \quad \hat{S}^- = \begin{pmatrix} 0 & 0 \\ 1 & 0 \end{pmatrix}, \tag{2}$$

where Ω is the Rabi frequency of the coherent atom–light interaction.

Time evolution of the system is described by the following equation for density matrix ρ

$$\dot{\rho} = \frac{1}{i\hbar}[\hat{H}_{AL}, \rho] + \gamma \mathcal{L}_1 \rho + 4\Gamma \mathcal{L}_2 \rho \,, \tag{3}$$

where $\mathcal{L}_1\rho = \hat{S}^-\rho\hat{S}^+ - \frac{1}{2}\hat{S}^+\hat{S}^-\rho - \frac{1}{2}\rho\hat{S}^+\hat{S}^-$ and $\mathcal{L}_2\rho = \hat{S}_z\rho\hat{S}_z - \frac{1}{2}\hat{S}_z\hat{S}_z\rho - \frac{1}{2}\rho\hat{S}_z\hat{S}_z$ are operators describing the effects of spontaneous emission (rate γ), and the phase noise (amplitude Γ), respectively.

The system dynamics can be described also in the dressed-state basis $\{|1\rangle, |2\rangle\}$, defined as:

$$|1\rangle = \cos\theta\,|g\rangle + \sin\theta\,|e\rangle \,, \tag{4}$$

$$|2\rangle = -\sin\theta\,|g\rangle + \cos\theta\,|e\rangle \,, \tag{5}$$

with $2\theta = -\arctan(\Omega/\Delta)$, $\Delta = \omega_L - \omega_0$. The Hamiltonian of the coherently driven system with the stochastic perturbation takes the form

$$\hat{H}_{AL} = E_1\,|1\rangle\langle 1|\,E_2\,|2\rangle\langle 2| - \hbar\delta\omega_a(t)\frac{\Omega}{\sqrt{\Omega^2+\Delta^2}}\left(|1\rangle\langle 2| + |2\rangle\langle 1|\right). \tag{6}$$

The first part of $\hat{H}_{AL}$ describes a dressed atom unperturbed by noise ($E_{1,2} = \pm\frac{1}{2}\hbar\sqrt{\Omega^2+\Delta^2}$ are energies of the dressed states). The second term represents the mixing of the dressed states by stochastic perturbation.

2.2. Resonance fluorescence

We want to analyze the effect of noise on the spectrum of RF $S(\omega)$ which can be calculated as

$$S(\omega) = \mathrm{Re}\,G(\omega)\,, \qquad G(\omega) = \lim_{t\to\infty}\int_0^\infty e^{-i\omega\tau}\left\langle \hat{S}^+(t+\tau)\hat{S}^-(t)\right\rangle \mathrm{d}\tau\,, \tag{7}$$

where $\hat{S}^{\pm}(t)$ are time-dependent atomic operators. For calculating time evolution of $\hat{S}^{\pm}(t)$ we use a new simulation approach [5] which extends the quantum trajectory method [6].

Figure 1 represents simulated spectra corresponding to a resonant tuning of the driving laser, $\Omega = 50\gamma$ and for various values of the ratio between the stochastic noise and coherent driving strengths Γ/Ω (after [5]).

When coherent driving is stronger than the stochastic perturbation, $\Gamma/\Omega < 1$, the spectrum resembles the familiar Mollow triplet [7] with well-resolved components (Fig. 1a). With an increase of the noise amplitude, the Rabi sidebands are less resolved and when Γ becomes comparable to Ω, the spectrum takes the shape of a single, noise broadened resonance with a narrow dip at ω_L (Fig. 1b,c). The dip narrows down as Γ increases and asymptotically approaches natural linewidth γ, when $\Gamma/\Omega \to \infty$. For nonresonant excitation, $\Delta \neq 0$, the spectrum exhibits a similar dependence on Γ/Ω, yet the dip is transformed into a narrow Fano-like resonance, as shown in Fig. 2 for two cases: $\Gamma/\Omega = 0.2$ and $\Gamma/\Omega = 3$.

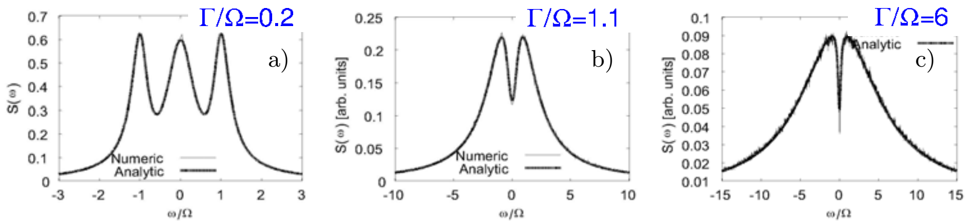

Fig. 1. RF spectra for various noise magnitudes obtained from analytic calculations (solid) and numerical simulations (dotted) for $\Delta = 0$ and $\Omega = 20\gamma$. The vertical scale is in arbitrary units, the same for all three plots. Frequency ω is defined relative to ω_L

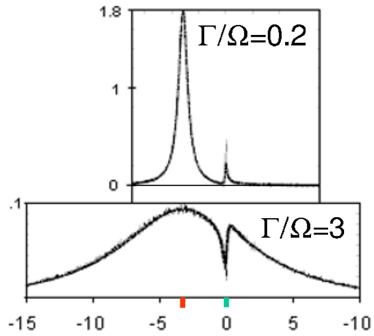

Fig. 2. RF spectra, $S(\omega)$, for two different values of Γ/Ω calculated for $\Delta/\Omega = 3$ and $\Omega = 20\gamma$ (same ω and vertical scales for both graphs). The vertical ticks mark the atomic resonance and laser light frequencies

This resemblance to the Fano resonance is a signature of the interference nature of the narrow feature seen for $\Gamma > \Omega$ [8].

3. Experimental

The above theoretical predictions have been verified in a preliminary experiment with a setup outlined in Fig. 3 [9]. Light beam from a pulsed dye laser (about 6 ns pulses of 10–100 μJ energy) is tuned close to the 553.5 nm resonance line of Ba (6 $^1S_0$–$^1P_1$ transition). The beam is directed to an oven containing Ba vapors (density about 10^{13} cm^{-3}) and some 1–10 Torr of Ar buffer gas. A side window in the oven allows the study of the fluorescence light which is spectrally analyzed with a resolution better than 10 GHz. Figure 4 depicts typical RF spectra recorded for various laser pulse intensities, barium densities and buffer gas pressures.

For sufficiently intense light, the fluorescence spectrum splits into two components (Fig. 4a), rather than in three members of the collisionless Mollow triplet.

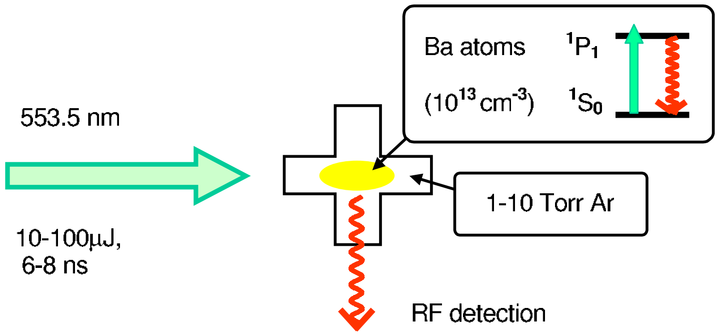

Fig. 3. Experimental setup for studying RF spectra of Ba atoms colliding with buffer gas atoms. Laser pulses (6–8 ns long and 10–100 μJ intense) are tuned close to the 553.5 nm resonance line. RF spectra are recorded by a side window with about 6 GHz resolution

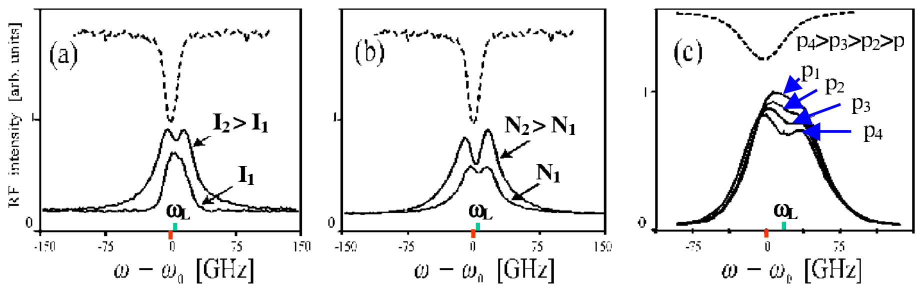

Fig. 4. Dependences of the RF spectra on: (a) light intensity – the dip at ω_L appears only when the light is sufficiently intense ($I_1 \approx 8$ μJ, $I_2 \approx 800$ μJ, $N_{\mathrm{Ba}} \approx 10^{12}$ cm^{-3}, $p_{\mathrm{Ar}} \approx 5$ Torr); (b) atomic density – for low barium density the dip occurs at ω_L, at higher densities the reabsorption at ω_0 overwhelms the interference dip ($N_1 \approx 10^{12}$ cm^{-3}, $N_2 \approx 10^{13}$ cm^{-3}); (c) the interference dip at $\omega_L \neq \omega_0$ requires sufficient collisional perturbation, it disappears when the rate of collisions, i.e. buffer gas pressure, is too low ($p_1 < p_2 < p_3 < p_4$). Upper broken-line plots are reference absorption spectra for ω_0 calibration

Observation of the double peak, or the dip structure, requires sufficient collisional perturbation, $\Gamma \geq \Omega$, which is provided by a sufficiently high buffer gas pressure (Fig. 4c). For a given light intensity, the splitting of the resonance components (the dip depth and width) increases with increasing barium density (Fig. 4b) and/or buffer gas pressure (Fig. 4c). The dip is clearly not due to reabsorption since it appears at ω_L, rather than ω_0, and at Ba densities lower than those necessary for reabsorption (Fig. 4b). Within the range of available laser energies, we have not seen triplet spectra. Apart from that, the measured spectra are very similar to those calculated (Figs. 1), convoluted with an instrumental profile of our spectrometer.

4. Noise-Induced Interference

In Ref. [5] we have shown that stochastic perturbation resulting in fluctuations of an atomic resonance frequency is also responsible for correlating phases of atomic dressed states. This correlation can be visualized by analyzing the phase difference $\Delta\phi = \phi_1 - \phi_2$ which stabilizes around 0 and π [5] (ϕ_1 and ϕ_2 are phases of the dressed-state components in atomic wavefunction $|\Phi\rangle = a_1 e^{i\phi_1}|1\rangle + a_2 e^{-i\phi_2}|2\rangle$).

In this paper we use yet another interpretation of the noise-induced interference in terms of indistinguishable spontaneous emission channels. The inelastic part of the resonance fluorescence spectrum, Eq. (7), can be represented by three contributions associated with three different poles s_+, s_- and s_0

$$S(\omega) = \mathrm{Re}\left(\frac{A_+}{\omega - s_+} + \frac{A_-}{\omega - s_-} + \frac{A_0}{\omega - s_0}\right). \tag{8}$$

Detailed expressions for the poles and amplitudes A_+, A_-, and A_0 are given in Ref. [5]. Here, we concentrate exclusively on the physical interpretation of their properties. Figure 5 represents the real and imaginary parts of $s_\pm$ versus Ω/Γ. One can distinguish two distinct regimes: the coherent regime where $\Omega > \Gamma$, and the opposite strong-noise regime where $\Omega < \Gamma$. In the coherent regime, $\mathrm{Re}\, s_\pm$ have two distinct values which represent positions of well-resolved sidebands of the Mollow triplet. The sideband widths, determined by $-\mathrm{Im}\, s_\pm$, are equal to $\Gamma + 3\gamma/4$. In the strong-noise regime, positions of the sidebands are degenerate at ω_L, while their widths are different: one approaches $3\gamma/2$, the second 2Γ. Also, the signs of $A_\pm$ are opposite, so that when $\Omega > \Gamma > \gamma$, the spectrum consists of a wide background and a narrow dip, as seen in our simulations (Fig. 1b,c) and in the experiment (Figs. 4).

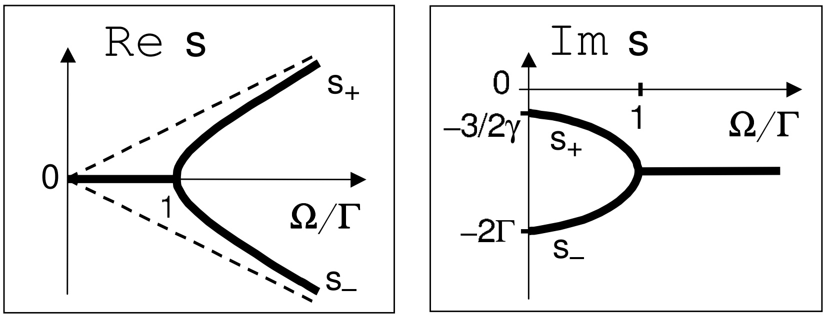

Fig. 5. Dependences of $\mathrm{Re}\, s_\pm$ and $\mathrm{Im}\, s_\pm$ on Ω/Γ. In the coherent regime ($\Omega > \Gamma$), $s_\pm$ are complex numbers corresponding to frequencies of the two sidebands placed symmetrically around the central component at ω_L and having width $\Gamma + 3\gamma/4$. In the strong-noise regime ($\Omega < \Gamma$), $s_\pm$ are entirely imaginary numbers corresponding to sidebands with degenerate frequencies and different widths

In the dressed-atom picture, the three contributions to Eq. (8) can be associated with three transitions between the dressed states, such as shown in Fig. 6.

Figure 6a represents the coherent regime with narrow, nondegenerate dressed states and, consequently, distinguishable spontaneous transitions yielding a three-peaked Mollow spectrum. Figure 6b corresponds to the strong-noise regime where the dressed levels are unresolved with respect to their width, hence the spontaneous emission lines are degenerate too. In such case, transition from any of the upper states, e.g. $|1\rangle''$ to the lower one $|1\rangle'$ can proceed either by a collisional (noise) mixing $|1\rangle'' - |2\rangle''$ followed by a photon emission to $|1\rangle'$, or by a reverse time ordering between collision and photon emission: $|1\rangle'' \rightarrow |2\rangle' \rightarrow |1\rangle'$. These two transition channels are indistinguishable, so the net probability exhibits interference contribution. This noise-induced interference is responsible for the narrow spectral features in observed spontaneous emission in the strong-noise regime. There are also two other possible emission channels: one direct emission: $|1\rangle'' \rightarrow |1\rangle'$ with no collision, second: $|1\rangle'' \rightarrow |2\rangle'' \rightarrow |2\rangle' \rightarrow |1\rangle'$ associated with two collisions. These two channels, however, differ substantially and do not contribute to the interference.

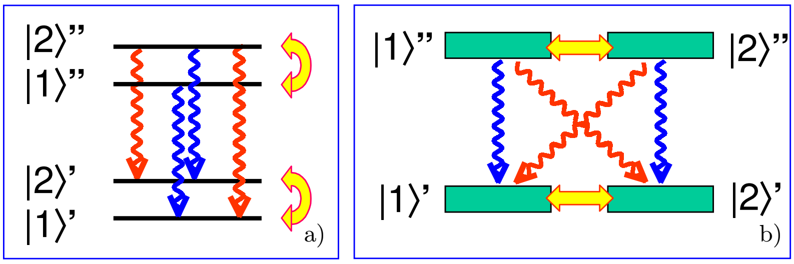

Fig. 6. Energy level diagram illustrating all possible transitions between the dressed-atom states. Wavy arrows represent the spontaneous emission and double arrows the state mixing by stochastic perturbation. If atomic dressed states are well resolved (a), the spontaneous emissions are spectrally well-resolved transitions. In the opposite case (b), the spontaneous emission channels overlap. In such a case, indistinguishable emission channels occur which lead to noise-induced interference

5. Conclusions

We have shown that stochastic perturbation can create phase correlation in a coherently driven two-level system. This correlation results in a noise-induced interference that affects atomic RF spectra. We have verified our theoretical predictions experimentally and observed RF spectra of collisionally perturbed atoms with a distinct narrow dips immune to the collisional broadening. It is interesting to extend the discussion to the case of a phase noise of the driving field. In Ref. [10] we discussed such a case and observed competition between two kinds of noise. In a future work we will also analyze the relation of the observed phenomena to the Quantum Zeno Effect.

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Notes

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