Dynamical Suppression of Spontaneous Emission

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We study the influence of a strong resonant driving field on the spectral properties of a single, cavity-confined, two-level atom. Under conditions of atom-cavity resonance, the lines of the Mollow resonance fluorescence triplet are found to narrow with increasing driving-field strength, indicating a dynamical decoupling of the atom from the vacuum field.

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Cavity electrodynamics has been one of the central topics of quantum optics in recent years, and interesting new aspects of collective and single-atom behavior have come to light. These effects include, for example, the enhancement[1] and suppression[2,3] of spontaneous emission, vacuum Rabi splittings,[4,5] photon antibunching,[6,7] optical bistability,[8] and squeezing.[7,9] In the present Letter, we discuss new effects related to the spontaneous decay of cavity atoms in the presence of a strong driving field. We find that a resonant driving field may, under the conditions outlined below, act to suppress dynamically the rate of spontaneous emission and dramatically modify the spectrum of resonance fluorescence.

An important aspect of the present problem is the finite response time of the reservoir (photon modes centered at the cavity line) responsible for spontaneous emission in the cavity. The existence of a finite reservoir response time means that the atom-reservoir interaction may be regarded as non-Markovian. In contrast to the situation in free space, where non-Markovian quantum electrodynamic effects are predicted but are practically nonobservable,[10] we are confronted here with a situation in which the non-Markovian aspects of the problem approach dominance.

Let us start with a description of the physical situation that we have in mind. A single two-level atom is placed in the center of a cavity and driven by a laser field. We denote the atomic and laser frequencies by ω_0 and ω_L, respectively, and assume, for simplicity, that $\omega_L=\omega_0$. Inside the cavity, the density of electromagnetic field modes is assumed to display periodic peaks as a function of frequency (as, for example, in an optical confocal resonator), and, for simplicity, these peaks are assumed to possess a Lorentzian profile. One of the cavity mode peaks is taken to be exactly resonant with the atom, and adjacent cavity mode peaks are ignored. The results obtained here depend critically on the existence of spectral peaks in the cavity mode structure; however, the exact shape of the peaks is important only in terms of details.

Our aim will be to calculate the stationary resonance fluorescence spectrum[11] of the atom as a function of driving-field strength, and to note those novel features that result from the finite memory time of the cavity photon reservoir. The Hamiltonian of our system reads

$$H=\omega_0\left(\frac{\hat\sigma_3}{2}\right)+\frac{\Omega}{2}(e^{i\omega_0 t}\hat\sigma+e^{-i\omega_0 t}\hat\sigma^\dagger)+\int|k_c|a^\dagger_{k_c}a_{k_c}dk_c+\int|k_b|b^\dagger_{k_b}b_{k_b}d^3k_b$$
$$+\int g_c(k_c)(a^\dagger_{k_c}\hat\sigma+\hat\sigma^\dagger a_{k_c})dk_c+\int g_b(k_b)(b^\dagger_{k_b}\hat\sigma+\hat\sigma^\dagger b_{k_b})d^3k_b, \quad (1)$$

where the σ's are the usual Pauli matrices describing a two-level atom, and Ω is the Rabi frequency of the external driving field. The operators $(a^\dagger_{k_c},a_{k_c})$ and $(b^\dagger_{k_b}, b_{k_b})$ correspond, respectively, to modes associated and unassociated with the cavity resonance. The coupling constants g in Eq. (1) are proportional to the appropriate photon mode densities. Since $|g_b(k_b)|^2$ is needed only in the neighborhood of ω_0, it may be treated as a constant. As specified above, $|g_c(k)|^2$ is taken for convenience only to be a Lorentzian of half width Γ.

As a first step in calculating the fluorescence spectra, we derive modified Bloch equations to describe the time evolution of the atomic observables σ, $\sigma^\dagger$, and σ_3. In doing so, it is convenient to introduce the reservoir response functions

$$\int dk_c|g_c(k_c)|^2\exp\{i[|k_c|-\omega_c]t\}\simeq\gamma_c\Gamma e^{-\Gamma|t|}, \quad (2a)$$

$$\int d^3k_b|g(k_b)|^2\exp\{i[|k_b|-\omega_c]t\}\simeq\gamma_b\delta(t), \quad (2b)$$

where ω_c is the cavity resonance frequency (assumed equal to ω_0). Note that Eq. (2a) expresses the fact that photon modes associated with the cavity resonance have a finite response time Γ^{-1}, while Eq. (2b) indicates that the modes not associated with the cavity resonance respond instantaneously. As will become clearer below, the interaction of the atom with the cavity (background) modes contributes an amount γ_c (γ_b) to the overall undriven spontaneous emission rate γ_s. We can thus write

$$\gamma_s = \gamma_b + \gamma_c. \tag{3}$$

With use of Eq. (2) and the Heisenberg equations for our system, modified Bloch equations are derived by the elimination of photon operators through a first-order expansion in γ_b and γ_c (Born approximation), with it kept in mind that the Markov property does not hold for the cavity resonance modes. In performing the expansion, we assume that both photon reservoirs are initially in the vacuum state (no fluorescence photons). The modified Bloch equations[12] take the form

$$\dot{\sigma}(t) = (i\Omega/2)\sigma_3(t) - \gamma_c\Gamma\int_0^t e^{-\Gamma(t-t')}\cos[\Omega(t-t')]\sigma(t')dt' - i\gamma_c\Gamma\int_0^t e^{-\Gamma(t-t')}\sin[\Omega(t-t')]\{[\sigma_3(t')+1]/2\}dt' - \gamma_b\sigma(t), \tag{4a}$$

$$\dot{\sigma}^\dagger(t) = -(i\Omega/2)\sigma_3(t) - \gamma_c\Gamma\int_0^t e^{-\Gamma(t-t')}\cos[\Omega(t-t')]\sigma^\dagger(t')dt' + i\gamma_c\Gamma\int_0^t e^{-\Gamma(t-t')}\sin[\Omega(t-t')]\{[\sigma_3(t')+1]/2\}dt' - \gamma_b\sigma^\dagger(t), \tag{4b}$$

$$\dot{\sigma}_3(t) = i\Omega[\sigma(t)-\sigma^\dagger(t)] - 2\gamma_c\Gamma\int_0^t e^{-\Gamma(t-t')}\{1+\cos[\Omega(t-t')]\}\{[\sigma_3(t')+1]/2\}dt' - i\gamma_c\Gamma\int_0^t e^{-\Gamma(t-t')}\sin[\Omega(t-t')][\sigma(t')-\sigma^\dagger(t')]dt' - 2\gamma_b[\sigma_3(t)+1]. \tag{4c}$$

The following comments should be made about the above equations. (i) They are valid only in the sense of the Born expansion in γ_c and γ_b. This means that γ_c and γ_b should be much smaller than Ω and Γ. These equations do not describe, therefore, the regime of vacuum Rabi splitting, $\Omega=0$, $\Gamma \simeq \gamma_c$. They may also in principle lead to some nonphysical effects (such as a negative power spectrum). Such effects will be, however, of order $(\gamma_c/\Omega)^2$ or $(\gamma_c/\Gamma)^2$ and may be corrected in the course of a more systematic expansion. (ii) They contain the usual (Markovian) contribution from the off-resonance (background) modes. (iii) Terms associated with the cavity modes have characteristic convolution-type memory integrals.[13] The memory extends over the cavity response time Γ^{-1}. (iv) In the limit $\Gamma \gg \Omega, \gamma_c$, Eqs. (4) reduce to the usual Bloch equations with $\gamma_s = \gamma_b + \gamma_c$ [as indicated earlier in Eq. (3)].

By adding Eqs. (4a) and (4b), we obtain an expression for the component of the Bloch vector $u = \sigma + \sigma^\dagger$, which is parallel to the driving-field vector. In the limit of small γ_c, u can be written as

$$u(t) = u(0)\exp[-\gamma(\Omega)t], \tag{5}$$

where $\gamma(\Omega) = \gamma_b + \{\gamma_c\Gamma^2/(\Gamma^2+\Omega^2)\}$. Interestingly, for $\Omega \gg \Gamma$, $\gamma(\Omega) \simeq \gamma_b$. We conclude that in the high-driving-field limit, the atom-cavity mode interaction does not result in a misalignment of the Bloch vector from the driving-field direction if it is so aligned initially. *In effect, one can dynamically switch off the effect of the cavity vacuum field on the component of the Bloch vector along the driving field.* As will be discussed elsewhere, this fact leads to a means of generating strongly squeezed atomic fluorescence. If the background mode density is small, as might be the case in suitably designed optical or microwave cavities, the u component of the Bloch vector may be stable for extended periods.

The other two components of the Bloch vector will (to the leading order in γ_c) undergo Rabi oscillations at the dynamically shifted frequency

$$\Omega'(\Omega) = \Omega[1 + \gamma_c\Gamma/2(\Gamma^2+\Omega^2)], \tag{6a}$$

and decay at the rate

$$\gamma'(\Omega) = \gamma_c[1 + \Gamma^2/2(\Gamma^2+\Omega^2)] + \tfrac{3}{2}\gamma_b. \tag{6b}$$

The decay rates of these components will thus be suppressed from $3(\gamma_c+\gamma_b)/2$ (for $\Omega \ll \Gamma$) to $\gamma_c + 3\gamma_b/2$ (for $\Omega \gg \Gamma$). Evidently, only the u component of the Bloch vector can be completely decoupled from the cavity modes.

The decoupling of the atom from the vacuum field can be understood qualitatively if one models the vacuum field as a classical fluctuating field that adds to the externally applied driving field. For example, a stability analysis of the semiclassical Bloch vector motion shows that the decay of the u component is triggered exclusively by driving-field fluctuations at the frequencies $\omega_0 \pm \Omega$. In the cavity, for $\Omega \gg \Gamma$, vacuum-field–induced driving-field fluctuations are substantially suppressed at these frequencies. Alternatively, for $\Omega \gg \Gamma$, the laser and cavity vacuum fields can be pictured as forming an effective field fluctuating slowly on the time scale Ω^{-1}. The stability of the u component of the Bloch vector can then be seen to result from adiabatic following.

The above observations regarding the decay of the Bloch vector suggest that the atom's high-driving-field ($\Omega \gg \Gamma$) resonance fluorescence spectrum will differ in the following qualitative ways from the standard Mollow spectrum.[11] The central component of the spectrum (corresponding to the u component of the Bloch vector) will be dramatically narrowed compared to its $\Omega \ll \Gamma$ value if $\gamma_c \gg \gamma_b$. Under the same conditions, the sidebands should also be narrowed (though not as dramatically as the central peak) and shifted.[14]

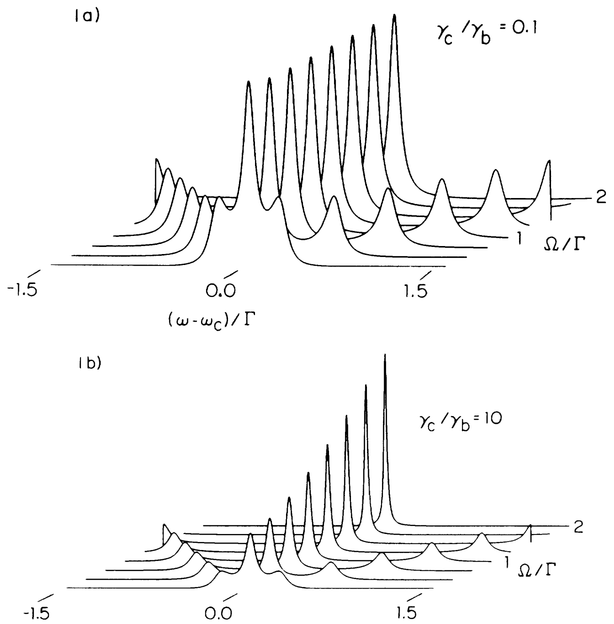

FIG. 1. Off-axis resonance fluorescence spectra of a single driven atom in a cavity. The driving field, atom, and cavity are exactly resonant. Horizontal: observation frequency; vertical: relative fluorescence intensity. Successive traces correspond to increasing driving-field strength. The quantity $\gamma_s = \gamma_c + \gamma_b = \Gamma/20$ throughout.

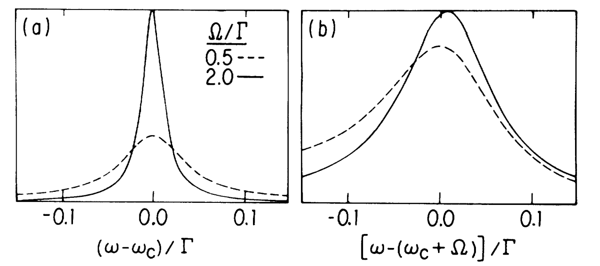

FIG. 2. High-resolution studies of selected peaks shown in Fig. 1(b). In (a), the strong narrowing of the central fluorescence peak is displayed, while in (b), the relatively minor narrowing of one of the symmetric side peaks is shown. Axes are described for Fig. 1.

Since we describe the atomic dynamics as a non-Markovian stochastic process, the quantum regression theorem cannot be employed, and the validity of the intuitive conclusions drawn above must be carefully scrutinized. This fact is brought out in the complexity of the equations for the two-time, dipole-dipole, correlation functions which are normally involved in the calculation of the power spectrum. It turns out that the equations do not take the same form as Eq. (4). In fact, the correlation functions $\langle \hat{\sigma}^\dagger(t)\hat{\sigma}(t')\rangle$, $\langle \hat{\sigma}(t)\hat{\sigma}(t')\rangle$, and $\langle \hat{\sigma}_3(t)\hat{\sigma}^\dagger(t')\rangle$ couple in a nontrivial manner to correlations of the type $\langle \hat{\sigma}^\dagger(t)\hat{P}_+(t')\rangle$, $\langle \hat{\sigma}(t)\hat{P}_+(t')\rangle$, and $\langle \hat{\sigma}_3(t)\hat{P}_+(t')\rangle$, where $\hat{P}_+ = (1+\hat{\sigma}_3)/2$ is a projection onto the upper atomic state.[13] It is surprising to find that the conjectures made earlier concerning the stationary spectrum do nonetheless hold. The easiest way to see that is to use the alternative definition of the power spectrum,

$$S_i(k) = \lim_{T\to\infty} \frac{1}{T}\langle \epsilon^\dagger_{k_i}(T)\epsilon_{k_i}(T)\rangle, \tag{7}$$

where $\epsilon = a$ or b depending on whether the fluorescence is into the cavity $(i=c)$ or background $(i=b)$ modes, respectively. In the latter case, integration over all solid angles has been performed. Equation (7) relates spectra to the single-time correlation functions which do fulfill equations of the same form as Eq. (4).

Equations (4) may be written in the Laplace-transformed form

$$G(z)\tilde{\sigma}(z) = K(z) + \sigma(0), \tag{8}$$

where $G(z)$ is a 3×3 resolvent matrix, $\tilde{\sigma}(z) = [\tilde{\sigma}(z), \tilde{\sigma}^\dagger(z), \tilde{\sigma}_3(z)]$ is a Laplace-transformed Bloch vector, and $K(z)$ contains the contributions from inhomogeneous terms in Eq. (4). With use of Eq. (8), the incoherent part of the spectra [Eq. (7)] may be written as

$$S_c(k) = \mathrm{Re}\left\{ \frac{\gamma_c^2\Gamma^2}{\Gamma^2+k^2} \sum_{i=1}^{3} G_{2i}^{-1}(ik)\left[K_i(ik)\sigma_{\mathrm{st}} + \tfrac{1}{2}\delta_{i2}(\sigma_{3\mathrm{st}}+1) - \delta_{i3}\sigma_{\mathrm{st}}\right]\right\}, \tag{9a}$$

$$S_b(k) = (\gamma_b^2/\gamma_c^2)[(\Gamma^2+k^2)/\Gamma^2]S_c(k), \tag{9b}$$

where $G^{-1}(z)$ is the inverse matrix of $G(z)$, δ_{ij} denotes the Kronecker delta, and σ_{st} ($\sigma_{3\mathrm{st}}$) represent stationary values of corresponding atomic variables. Evaluating Eq. (9b) in the limit $\Omega, \Gamma \gg \gamma_c, \gamma_b$, one finds that the width of the central peak is given by $\gamma(\Omega)$ [Eq. (5)], and its height is

$$h_{\mathrm{cent}} = \frac{\gamma_b^2(\Gamma^2+\Omega^2)}{4[\gamma_b(\Gamma^2+\Omega^2)+\gamma_c\Gamma^2]}. \tag{10}$$

Under the same conditions, the width of the side peaks is found to be $\gamma'(\Omega)$ [Eq. (6b)], while their height is

$$h_{\mathrm{side}} = \frac{\gamma_b^2}{8}\frac{\gamma'(\Omega)}{\gamma'^2(\Omega)+[\Omega'(\Omega)-\Omega]^2}. \tag{11}$$

The sidebands are shifted from the central frequency by $\Omega'(\Omega)$ [Eq. (6a)].

In Fig. 1, we show off-axis resonance fluorescence spectra calculated from Eq. (9b) for a variety of illustrative situations. In Fig. 1(a), where the influence of the cavity is weak, variations of the resonance fluorescence spectra from the standard Mollow form[11] are essentially unobservable. In Fig. 1(b), on the other hand, the central peak of the spectrum shows a strong narrowing as the driving-field Rabi frequency exceeds the cavity-resonance halfwidth Γ. In Fig. 2, we focus on the behavior of selected peaks shown in Fig. 1(b).

In closing, we note that experimental study of the effect described here should provide interesting insight into the fundamental problem of atomic dynamics in the presence of coupling to finite-bandwidth, arbitrary–line-shape reservoirs. Furthermore, the conditions relevant to Fig. 1(b) should be experimentally realizable in optical or microwave cavities. A more detailed analysis which also addresses the statistical properties of the atomic fluorescence and accounts for atom-field detuning will be presented elsewhere.

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