Noise in dead zones

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We investigate the quantum noise gained by a weak probe laser interacting with a driven two-level system in the Doppler-broadening limit in which the so-called "dead zone" of the absorption spectrum is formed. We find that the plot of the noise as a function of the pump-probe detuning (which we shall refer to as the "noise spectrum") is hardly affected by the Doppler broadening for the case of a resonant pump. It consists of a large central peak flanked by two smaller peaks at the edges of the dead zone, which are at the Rabi sideband frequencies. Thus the dead zone, which is produced when absorption and gain regions of the absorption spectrum are averaged by Doppler broadening, is a rather noisy region, indicating the occurrence of many competing absorption and emission processes. The probe noise spectrum for a detuned pump is quite different in the presence and absence of Doppler broadening as is the absorption spectrum. On Doppler broadening the absorption spectrum consists of a dead zone centered at the Rayleigh resonance frequency with asymmetric absorption peaks on either side. The noise spectrum has a peak at the center with two sidebands at the edges of the dead zone. The sidebands are of almost the same intensity for short path lengths, but the peak on the side of the three-photon scattering becomes larger than the other peak as the path length increases.

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I. INTRODUCTION

Recently [1–3], there has been a revival of interest in the zones of almost zero probe absorption ("dead zones"), which occur when a driven two-level system is interrogated by a weak copropagating beam in the limit of large Doppler broadening. These dead zones, first predicted by Baklanov and Chebatoev [4], were described in detail by Khitrova, Berman, and Sargent [5]. Essentially, the two points of zero absorption that occur for a resonantly driven two-level system at the probe frequency $\omega_2 = \pm 2V_{ba}$ (where $2V_{ba} = d_{ba}E_1/\hbar$ is the Rabi frequency of the driving field) are replaced in the Doppler limit [limit in which the Doppler width $D > 2V_{ba} > (1/T_2)$] by a zone of almost zero absorption between the two points $-2V_{ba} < \omega_2 < 2V_{ba}$. The renewed interest [1–3] in dead zones has mainly been concerned with the behavior of the refractive index within the dead zone. It was predicted by Scully and co-workers [6] that a variety of coherently prepared driven three- and four-level systems can exhibit a maximum in their refractive index at a point of zero absorption. It was also shown [2,3,7,8] that the much simpler resonantly driven two-level system displays the same behavior. Subsequently, the effect of Doppler broadening on both the absorption and refractive index of this system was explored [1–3]. It was found that the minimum and maximum values of the refractive index that occur at points of nonabsorption in the absence of Doppler broadening are scarcely affected by Doppler broadening. Thus it seems that the dead zone is only inactive with respect to absorption, but not to dispersion. The origin of the dead zone and the associated behavior of the refractive index have recently been explained using velocity-dependent dressed states by Ling and Barbay [1]. (One should note that the behavior of the absorption [9,10] and of the absorption and refractive index [3], in the absence of Doppler broadening, has also been discussed in terms of dressed states.)

Another reason for the renewed interest [1] in dead zones is the apparent similarity between the dead zone and the region of transparency [termed electromagnetically induced transparency (EIT)] observed by Harris and co-workers [11] when an atomic system is driven by an intense field. This similarity is particularly impressive in the calculations of Moseley *et al.* [12] and Gea-Banacloche and co-workers [13] of the Doppler-broadened probe absorption and refractive index in a cascade three-level system interacting with a strong pump and weak counterpropagating probe. Again the extrema in the refractive index were found to occur at the edges of the region of near transparency. The calculations of Moseley *et al.* [12] were performed in order to explain their observation of focusing and defocusing of the weak probe within the EIT window, induced by the transverse intensity profile of the strong pump. In recent papers [14], we discussed the related problem of the nonlinear refractive index cross modulation produced in a two-level system interacting with an intense off-resonance laser and a second moderately strong laser tuned near one of the Rabi sidebands of the system. As a consequence, depending on the transverse profiles of the two lasers, self-focusing (defocusing) may be induced in the self-defocusing (focusing) nonlinear medium. Some recently observed anomalies in conical emission [15] could be explained qualitatively by these effects.

In the present paper, we consider an aspect of the dead zones that has not yet been discussed, namely, the quantum noise gained by a weak probe whose frequency ω_2 lies within the dead zone as it propagates through a medium of driven two-level atoms. We recently showed [16,17] that, in the absence of Doppler broadening, a weak probe propagating through a medium of resonantly driven two- or degenerate Λ-shaped three-level atoms [18] has three peaks in its "noise spectrum" (for brevity, we refer to the plot of the probe noise as a function of the pump-probe detuning as the noise spectrum), situated near points of probe nonabsorption. This is true for all values of $\alpha_0 L$ (where α_0 is the absorption coefficient and L is the propagation length) for the three-level system that is population trapped. However, for the two-level system, the noise spectrum changes from three to two peaks as the value of $\alpha_0 L$ increases, provided the pump Rabi frequency is sufficiently small. We interpreted the peaks in the noise spectrum at points of probe nonabsorption as indicative of the fact that, at these frequencies, a large number of absorption processes compete with an equally large number of emission events so that the average absorption is zero. Here we shall demonstrate that the probe noise spectrum, in common with the refractive index spectra, is only weakly affected by Doppler broadening. This is true because the dead zone, in analogy with the points of nonabsorption, results from the averaging out of many absorption and emission events.

The probe noise spectrum for a detuned pump is, however, quite different in the presence and absence of Doppler broadening. In the absence of Doppler broadening, it consists of a peak at the three-photon scattering (TPS) frequency, a smaller peak at the extraresonant frequency, and an even smaller peak at the absorption frequency due to the fact that the probe at this frequency has already been significantly absorbed on propagation. In the Doppler limit, the outer peaks move to the edge of the dead zone and are of the same intensity for small propagation lengths so that the spectrum is similar to that obtained for a resonant pump. The probable reason for this is the dominant contribution of the resonant velocity subgroup to the noise. As the propagation length increases, the peak on the side of the TPS frequency gradually becomes more intense than the peak on the side of the absorption frequency.

The noise properties of a weak quantum probe interacting with a driven two-level system were described by Boyd and co-workers [19]. They were interested in the effect of collisions on the noise associated with the amplification of the weak probe when the system is driven near resonantly by a strong pump [20]. We adopt the same formalism derived from the quantum theory of multiwave mixing developed by Agarwal [21]. We note that such a theory has also been derived using different formalisms by Sargent, Holm, and Zubairy [22] and Reid and Walls [23]. The equivalence between these theories and that of Agarwal has been discussed by Agarwal and Boyd [24].

II. GENERAL DERIVATIONS

We treat a two-level system interacting with a strong pump and a weak probe. The energy level scheme and geometry of the system are given in Fig. 1. The frequency offsets of the pump and probe frequencies $\omega_{1,2}$ from the resonance frequency ω_{ba} are given by $\Delta_{1,2}=\omega_{ba}-\omega_{1,2}$ and the pump-probe detuning by $\delta=\Delta_1-\Delta_2=\omega_2-\omega_1$.

We assume the pump to be linearly polarized and much stronger than the probe and write it (in the interaction picture) as a classical field, ignoring its depletion due to interaction with the medium:

$$\vec{E}_1(\vec{r},t)=\vec{\varepsilon}_1\exp[i(\vec{k}_1\cdot\vec{r}-\omega_1 t)]+\text{c.c.} \tag{1}$$

The much weaker probe is taken to be a quantum field written as

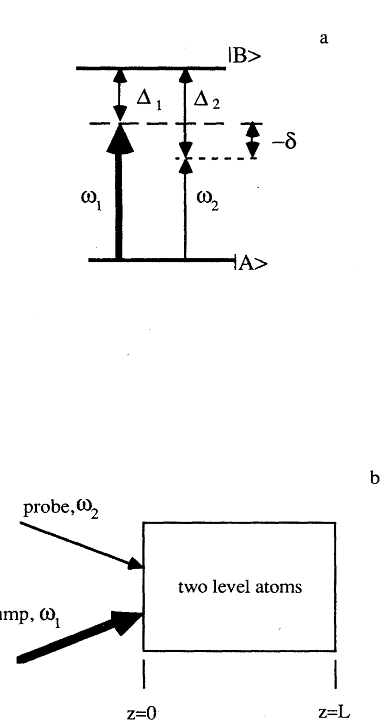

FIG. 1. (a) Two-level system. (b) Geometry of the atom-laser system.

$$\vec{E}_2(\vec{r},t)=\beta_2\vec{\varepsilon}_2\hat{a}\,\exp[i(\vec{k}_2\cdot\vec{r}-\omega_2 t)]+\text{H.c.} \tag{2}$$

Here $\hat{a}$ is the annihilation operator for the probe, $\vec{\varepsilon}_2$ is the polarization unit vector of the probe, $\beta_2=-i(\omega_2/V)^{1/2}$, and V is the quantization volume. We assume Planck's constant to be equal to unity here and in the following derivations. The following assumptions are made in the derivation: (i) the dipole approximation; (ii) the rotating-wave approximation, with the applied and generated field frequencies assumed to be very close to the atomic resonance; (iii) the Markov approximation, with the time scales associated with the field dynamics assumed to be much longer than those associated with the atomic dynamics; and (iv) the probe and the generated fields are so weak that it is sufficient to keep only terms up to second order in the annihilation operator for the probe.

We write the equation of motion for the density operator for the coupled atom-field system as

$$\frac{\partial\hat{\rho}}{\partial t}=-i[\hat{H}_{0A}+\hat{H}_{0F}+\hat{H}_1,\hat{\rho}]+\hat{L}_A\hat{\rho}. \tag{3}$$

Here $\hat{H}_{0A}$ and $\hat{H}_{0F}$ are the unperturbed Hamiltonians for the atom and the field; $\hat{H}_1$ is the atom-field interaction Hamiltonian and $\hat{L}_A$ is the relaxation operator, which involves contributions from both spontaneous emission and atomic collisions.

The equation for the density operator for the field $\hat{\rho}_2$ is derived by tracing the density operator $\hat{\rho}$ over the atomic variables and by applying projector operator techniques. The following master equation results [21]:

$$\frac{\partial \hat{\rho}_2}{\partial t} = -\frac{|\beta_2|^2 N}{2}(\tilde{C}^{+-}(-i\delta)[\hat{a}^\dagger,\{\hat{a},\hat{\rho}_2\}] + \tilde{Q}^{+-}(-i\delta) \times [\hat{a}^\dagger,[\hat{a},\hat{\rho}_2]]) + \text{H.c.} \tag{4}$$

Here N is the number of atoms in the medium and

$$\tilde{C}^{+-}(i\delta) = \int_0^\infty d\tau\, e^{-i\delta\tau} C^{+-}(\tau) = \langle \hat{P}_2^+(i\delta)\hat{P}_2^-(0)\rangle - \langle \hat{P}_2^-(0)\hat{P}_2^+(i\delta)\rangle, \tag{5}$$

$$\tilde{Q}^{+-}(i\delta) = \int_0^\infty d\tau\, e^{-i\delta\tau} Q^{+-}(\tau) = \langle \hat{P}_2^+(i\delta)\hat{P}_2^-(0)\rangle + \langle \hat{P}_2^-(0)\hat{P}_2^+(i\delta)\rangle - 2\langle \hat{P}_2^+(i\delta)\rangle\langle \hat{P}_2^-(0)\rangle, \tag{6}$$

where $\hat{P}_2^+ = \hat{\mathbf{P}}^+ \cdot \vec{\varepsilon}_2^{\,*}$ and $\hat{P}_2^- = \hat{\mathbf{P}}^- \cdot \vec{\varepsilon}_2$ are operators related to the polarization that drives the probe.

This master equation describes the characteristic change in the probe's mode due to linear absorption and emission in the presence of the pump. The functions $\tilde{C}^{+-}$ and $\tilde{Q}^{+-}$ are polarization correlation functions [21]: $\tilde{C}^{+-}$ is proportional to the semiclassical susceptibility whereas $\tilde{Q}^{+-}$ corresponds to the quantum fluctuations of the atomic system. Specifically, the real part of $\tilde{C}^{+-}$ determines the absorption and its imaginary part determines the dispersion whereas $\tilde{Q}^{+-}$ contributes to the quantum noise.

These correlation functions are calculated with the aid of the quantum regression theorem [25], which we summarize as follows: If

$$\langle G_i(t_1)\rangle = \sum_j g_{ij}(t_1 - t_2)\langle G_j(t_2)\rangle,$$

where $t_1 > t_2$, then

$$\langle G_k(t_2) G_i(t_1)\rangle = \sum_j g_{ij}(t_1 - t_2)\langle G_k(t_2) G_j(t_2)\rangle,$$

where G_j is a complete group of operators, g_{ij} are c functions, and the system is assumed to be Markovian.

In order to calculate these expectation values, we first recall that for the two-level system, the polarization operator is

$$\hat{\mathbf{P}} = \hat{\mathbf{P}}^+ + \hat{\mathbf{P}}^- = d_{ba}^* \vec{\varepsilon}_1^{\,*} |a\rangle\langle b| + d_{ba}\vec{\varepsilon}_1 |b\rangle\langle a|. \tag{7}$$

Assuming that the probe is in the $\vec{\varepsilon}_R$ direction, it follows that

$$\hat{P}_2 = \hat{P}_2^+ + \hat{P}_2^- = \hat{\mathbf{P}}^+ \cdot \vec{\varepsilon}_R^{\,*} + \hat{\mathbf{P}}^- \cdot \vec{\varepsilon}_R = d_{ba}^* |a\rangle\langle b| + d_{ba}|b\rangle\langle a|. \tag{8}$$

The quantum regression theorem requires the knowledge of the expectation values of the polarization operators. These expectation values involve the calculation of the elements of the density matrix $\rho_{ba} = \langle |a\rangle\langle b|\rangle$. In fact, the terms $\langle \hat{P}_2^+(0)\hat{P}_2^-(0)\rangle$ and $\langle \hat{P}_2^+(0)\rangle$ and their Hermitian conjugates always reduce to a sum of terms of the form $\langle |a\rangle\langle b|\rangle$. In order to calculate these matrix elements we turn to the Bloch equations [26] for the density matrix elements of this system, whose general form is given by $\dot{\psi}(\tau) = A\Psi(\tau) + B$, where Ψ is the vector of the density-matrix elements, A is a matrix, and B is a vector.

For the closed two-level system, the familiar equations [21] are

$$\Psi = \begin{pmatrix} \rho'_{ab} \\ \rho'_{ba} \\ \frac{1}{2}(\rho_{bb} - \rho_{aa}) \end{pmatrix},$$

$$A = \begin{pmatrix} -\frac{1}{T_2} + i\Delta_1 & 0 & 2iV_{ba}^* \\ 0 & -\frac{1}{T_2} - i\Delta_1 & -2iV_{ba} \\ iV_{ba} & -iV_{ba}^* & -\frac{1}{T_1} \end{pmatrix},$$

$$B = \begin{pmatrix} 0 \\ 0 \\ \frac{\eta}{2T_1} \end{pmatrix}. \tag{9}$$

Here $\rho_{aa} + \rho_{bb} = 1$ and $\rho'_{ab} = e^{-i\omega_1 t}\rho_{ab}$; η is the equilibrium value of the atomic inversion $\rho_{bb} - \rho_{aa}$ in the absence of the pump, V_{ba} is the pump Rabi frequency, and T_1 and T_2 are the longitudinal and transverse relaxation times.

When the Laplace transform is applied to this equation, it is possible to evaluate $\Psi(z) = U(z)\Psi(0)$, where $\Psi(0) = -A^{-1}B$ is the steady-state solution and $U(z) = (zI - A)^{-1}$, where I is the identity matrix. When $z = -i\delta$ the result gives the matrix elements required for the calculation of the polarization correlation functions.

In terms of the U matrix and the steady-state value of the density matrix elements $\Psi(0)$, the polarization correlation functions are thus given by

$$\tilde{C}^{+-} = |d_{ba}|^2\{-2U_{2,2}\Psi_3 + U_{2,3}\Psi_1\}, \tag{10}$$

$$\tilde{Q}^{+-} = |d_{ba}|^2\{-2U_{2,1}\Psi_1^2 + U_{2,2}(1 - 2\Psi_1\Psi_2) - 2U_{2,3}\Psi_1\Psi_3\}, \tag{11}$$

in agreement with Agarwal [21], where we have dropped the (0) from the steady-state density-matrix elements in these equations for simplicity.

In order to examine the gain and the noise in these systems, we now follow Gaeta and co-workers [19] and convert the master equation into a Langevin equation for the annihilation operator $\hat{a}$:

$$\frac{d\hat{a}}{dz} = -\alpha\hat{a} + \hat{f}(z), \tag{12}$$

where

$$\alpha = \alpha_0 \tilde{C}^{+-}(-i\delta) \equiv [2\pi N\omega_2 n/(h/2\pi)Vc]\tilde{C}^{+-}(-i\delta),$$

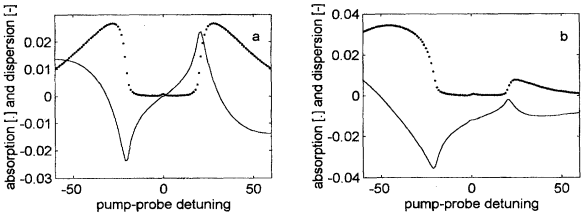

FIG. 2. Typical absorption $(\cdot\ \cdot\ \cdot\ \cdot)$ and dispersion (—) (proportional to $n-1$) spectra plotted as a function of δT_2 in the Doppler limit: (a) the pump is resonant with the atomic transition $\Delta_1 T_2=0$ and (b) the pump is detuned from the atomic resonance frequency such that $\Delta_1 T_2=-40$. Other parameters are $1/T_1=2(1/T_2)$, $V_{ba}T_2=10$, and $DT_2=50$.

$z=tc/n$ determines a steady-state traveling-wave situation, c/n is the phase velocity of the probe, and $\hat{f}(z)$ is the fluctuation operator such that $\langle \hat{f}(z)\rangle=0$. This equation is easily derived by multiplying the master equation from the left by $\hat{a}$ and taking the trace over the field variables.

Second-order correlation functions of $\hat{f}(z)$ can be derived in a similar way from the master equation by multiplying it by the number operator. Thus, in agreement with Gaeta and co-workers [19], the second-order correlation functions for the fluctuation (noise) operator are

$$\langle \hat{f}^\dagger(z)\hat{f}(z')\rangle=\alpha_0 \mathrm{Re}[\tilde{Q}^{+-}(-i\delta)-\tilde{C}^{+-}(-i\delta)]\delta(z-z'), \tag{13a}$$

$$\langle \hat{f}(z)\hat{f}^\dagger(z')\rangle=\alpha_0 \mathrm{Re}[\tilde{Q}^{+-}(-i\delta)+\tilde{C}^{+-}(-i\delta)]\delta(z-z'). \tag{13b}$$

Integration of the Langevin equation over the interaction length leads to the following equation [19], which is the familiar equation for a linear amplifier (see, e.g., [27]):

$$\hat{a}(L)=g\hat{a}(0)+\hat{F}\equiv e^{-\alpha L}\hat{a}(0)+\int_0^L \hat{f}(z)e^{\alpha(z-L)}dz, \tag{14}$$

in which g is the gain of the probe and $\hat{F}$ is the noise operator related to the probe amplification in a strongly driven atomic system. Also, according to Gardiner [28],

$$[\mathrm{noise}(\hat{a})]\alpha\langle(\hat{a}^\dagger-\langle\hat{a}^\dagger\rangle)(\hat{a}-\langle\hat{a}\rangle)\rangle\equiv\langle\hat{f}^\dagger\hat{f}\rangle \tag{15}$$

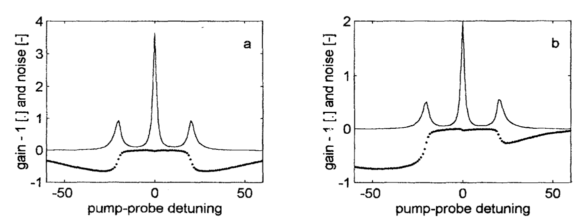

FIG. 3. Typical gain $(\cdot\ \cdot\ \cdot\ \cdot)$ and noise (—) spectra plotted as a function of δT_2 in the Doppler limit for $\alpha_0 L=40$: (a) the pump is resonant with the atomic transition $\Delta_1 T_2=0$ and (b) the pump is detuned from the atomic resonance frequency such that $\Delta_1 T_2=-40$. Other parameters are $1/T_1=2(1/T_2)$, $V_{ba}T_2=10$, and $DT_2=50$.

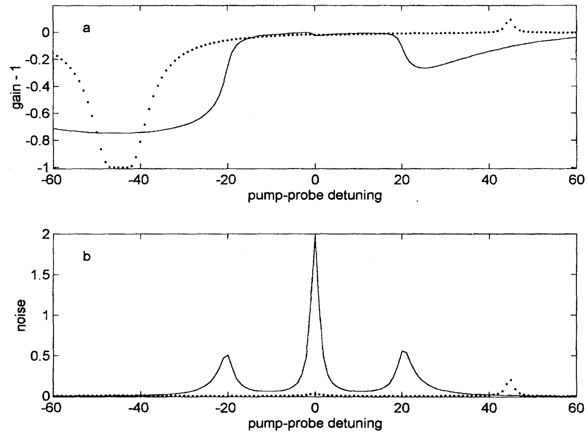

FIG. 4. For a detuned pump (a) the gain spectrum plotted as a function of δT_2 with (—) and without $(\cdot\ \cdot\ \cdot\ \cdot)$ Doppler broadening and (b) the noise spectrum plotted as a function of δT_2 with (—) and without $(\cdot\ \cdot\ \cdot\ \cdot)$ Doppler broadening. The parameters are $1/T_1=2(1/T_2)$, $V_{ba}T_2=10$, $\Delta_1 T_2=-40$, and $DT_2=50$.

and therefore over the interaction length, the amount of quantum noise introduced by the atomic fluctuations is given by

$$\langle \hat{F}^\dagger\hat{F}\rangle=\frac{1-|g|^2}{2\ \mathrm{Re}[\tilde{C}^{+-}(-i\delta)]}\mathrm{Re}[\tilde{Q}^{+-}(-i\delta)-\tilde{C}^{+-}(-i\delta)]. \tag{16}$$

Finally, the Doppler-broadened spectra for copropagating pump and probe beams are obtained by averaging numerically over all the correlation functions according to

$$\{\cdots(\Delta_1,\delta)\}_D=\left(\frac{1}{\pi D^2}\right)^{1/2}\int_{-\infty}^{\infty}d\Delta_1'\{\cdots(\Delta_1',\delta)\}\times\exp\left[-\frac{(\Delta_1-\Delta_1')^2}{D^2}\right], \tag{17}$$

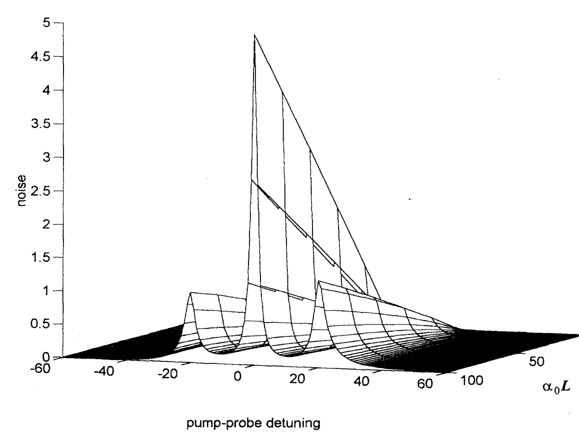

FIG. 5. Noise spectrum plotted as a function of δT_2 and of the propagation length for a detuned pump. The parameters are $1/T_1=2(1/T_2)$, $V_{ba}T_2=10$, $\Delta_1 T_2=-40$, and $DT_2=50$.

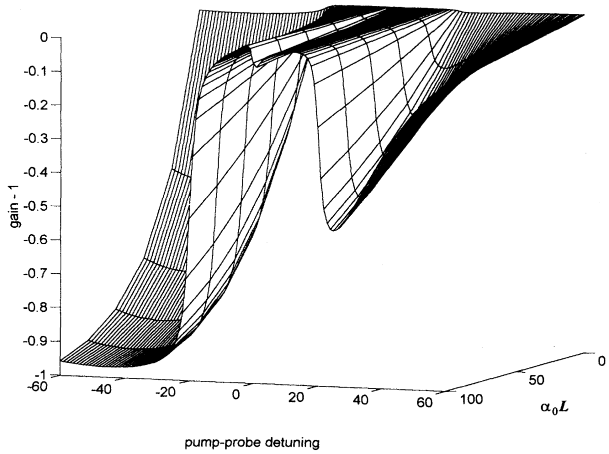

FIG. 6. Gain spectrum plotted as a function of δT_2 and of the propagation length for a detuned pump. The parameters are $1/T_1 = 2(1/T_2)$, $V_{ba}T_2 = 10$, $\Delta_1 T_2 = -40$, and $DT_2 = 50$.

where D is the $(1/e)$ Doppler half-width and then substituting the results in the relevant expressions.

III. RESULTS AND DISCUSSIONS

Typical probe absorption and dispersion spectra, calculated in the Doppler limit for the case of a resonant pump, are shown in Fig. 2(a). The spectra are in agreement with previous semiclassical treatments of the two-level system [2,3,7,8]. The dead zone is centered at $\delta = 0$ (that is, $\omega_2 = 0$) and lies between the Rabi sideband frequencies $\delta = \pm 2V_{ba}$, which are points of nonabsorption in the absence of Doppler broadening. The dispersion spectrum is broadened only slightly by the Doppler broadening, so that the dead zone is inactive only with respect to absorption but not to dispersion. This point will be reinforced further by the probe noise spectra.

When the pump is detuned from resonance with the atomic transition [see Fig. 2(b)], a dead zone of length $4V_{ba}$ centered at $\delta = 0$ is again obtained. Here, however, $\delta = 0$ corresponds to the center of the dispersive extraresonant feature [29,30], which appears in the absence of Doppler broadening. Thus the dead zone for the case of a detuned pump derives, in part, from the averaging out of this feature. In addition, the absorptive regions that flank the dead zone are now asymmetrical, with higher absorption on the side at which the absorption peak appears in the absence of Doppler broadening [31]. Similarly, the absolute value of the minimum in the dispersion on that side is greater than that of the maximum on the other side.

Probe gain and noise spectra are shown in Figs. 3(a) and 3(b) for the same parameters as in Figs. 2(a) and 2(b), respectively, for the case of $\alpha_0 L = 40$. The gain spectra for both the resonant pump [Fig. 3(a)] and detuned pump [Fig. 3(b)] exhibit the same dead zone bounded by absorptive regions as the absorption spectra of Figs. 2(a) and 2(b). The noise spectrum for the case of a resonant pump [Fig. 3(a)] is a slightly broadened version of that obtained in the absence of Doppler broadening. It consists of a central peak at $\omega_2 = 0$ and two peaks centered at the Rabi frequencies, which correspond to the points of nonabsorption in the absence of Doppler broadening and to the edges of the dead zone in the presence of Doppler broadening. Thus both the points of nonabsorption and the dead zone derive from the averaging out of many competing absorption and emission events.

In Figs. 4(a) and 4(b) we compare the probe gain and noise spectra for a detuned pump in the presence and absence of Doppler broadening. We see in Fig. 4(a) that the gain feature that appears, in the absence of Doppler broadening [31], at the three-photon scattering frequency $\delta = (\Delta_1^2 + 4V_{ba}^2)^{1/2}$ is replaced by an absorption feature at the nearest edge of the dead zone and that the absorption feature moves from $\delta = -(\Delta_1^2 + 4V_{ba}^2)^{1/2}$ to the other edge of the dead zone. The noise spectrum [Fig. 4(b)] consists, in the absence of Doppler broadening, of a peak at the three-photon scattering frequency, a smaller peak at the extraresonant frequency, and an extremely small peak at the absorption frequency due to the fact that nearly all the probe has been absorbed on propagation. In the Doppler limit, the outer peaks move to the edges of the dead zone and are almost of the same intensity for small propagation distances, with the peak on the side of the TPS frequency gradually overtaking the other sideband as $\alpha_0 L$ increases (see Fig. 5), as one would expect from the spectrum in the absence of Doppler broadening. The gain spectrum is shown as a function of $\alpha_0 L$ in Fig. 6: as the propagation length increases, the dead zone becomes more prominent due to increased absorption at the sidebands.

ACKNOWLEDGMENTS

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