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PROBE TRANSPARENCY IN A TWO-LEVEL MEDIUM EMBEDDED BY A SQUEEZED VACUUM

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ABSTRACT

Effect of the detuning on the probe absorption spectra of a two-level system with and without a classically driven field in a squeezed vacuum is investigated. For a strong squeezing, there is a threshold which determines the positions and widths of the absorption peaks, for the squeezed parameter M. In a large detuning, the spectra exhibit some resemblance to the Fano spectrum. The squeezing-induced transparency occurs at the frequence $2\omega_L - \omega_A$ in the minimum-uncertainty squeezed vacuum. This effect is not phase-sensitive.

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I. Introduction

With the successful generation of squeezed light in quite a few laboratories [1], a great deal of attention has been drawn to a reexamination of fundamental atomic radiative processes and the role of the squeezed light in their interaction with atomic system. The first investigation by Gardiner [2] showed that the two quadratures of the polarization of a two-level atom interacting with a squeezed vacuum decay at vastly different rates, where the decay of an atomic polarization quadrature may be inhibited. The modifications of the resonance fluorescence spectrum of such a system were studied by Carmichael, Lane and Walls [3], who showed that for large classical applied field strengths the spectrum is a triplet, as in the absence of the squeezed vacuum [4], but that the central peak of the triplet has a width that depends strongly on the relative phase of the driving field and the squeezed vacuum when the squeezing is strong. The nonsecular effects in the resonance fluorescence were also reported by Smart and Swain [5,6], where anomalous resonance fluorescence spectra (including dispersive profiles and hole burning at the centre, etc.) were found under certain parameter values. This subject is also extented into the cavity situation [7].

Recently, Ritsch and Zoller [8], and An, *et al* [9] have considered the weak probe beam absorption spectrum of a strongly driven two-level atom system embedded in a squeezed vacuum. They have shown that the absorption spectrum is phase-sensitive and shows subnatural linewidths, when both driving and squeezed fields are tuned to the atomic resonance. The squeezing-induced enhancement of a refractive index with zero absorption [10] and gain without inversion [11] have also been studied on resonance. These results are due to the influence of phase-sensitive field correlations on atomic dipole dynamics

On the other hand, as there is a non-zero detuning between the atomic transition frequency and the carried frequency of the broadband squeezed vacuum, the resonance fluorescence spectrum exhibits a strong asymmetry [6]. The effect of the detuning on the subnatural linewidths of the absorption spectrum is also mentioned by An, Sargent III and Walls [9], without giving the details.

More recently, Ficek and Dalton [12] have also considered the role of the detuning in the absorption spectrum of a two-level atom without a coherently driven field, in a Fabry-Perot microcavity with injecting a squeezed light. An interesting phenomenon, the squeezing-induced transparency, is obtained. Its potential applications have been discussed [13], and many schemes for generating this phenomenon have also been proposed [14].

In the present paper we focus our attention on the effect of the detuning, in detail on the weak probe absorption of a two-level atomic system, in the absence of and in the presence of a classically driven field, embedded by a broadband squeezed vacuum. In the strongly perfect squeezing and large detuning limits, the absorption spectra show Fano-like profiles, and the transparency effect for the probe field is obtained at the frequence $2\omega_L - \omega_A$.

II. Equations of Motion

The model to be considered is composed of a two-level atom having an excited state $|1>$ and a ground state $|0>$ with a transition frequency ω_A, driven by a monochromatic laser field of frequency ω_L with a complex amplitude E_L, and embedded in a broadband squeezed vacuum bath which is assumed to be centred about the frequency ω_L. After the system has reached a steady-state, the medium is illuminated by a weak (nonsaturating) probe field at frequency ω_p with an amplitude E_p. The time evolution equations for the density matrix elements of the sytem read [5, 6, 8, 9]

$$\frac{d\rho_{11}}{dt} = -\gamma(N+1)\rho_{11} + \gamma N\rho_{00} + iV_{10}\rho_{01} - iV_{01}\rho_{10} \tag{1a}$$

$$\frac{d\rho_{00}}{dt} = -\gamma N\rho_{00} + \gamma(N+1)\rho_{11} - iV_{10}\rho_{01} + iV_{01}\rho_{10} \tag{1b}$$

$$\frac{d\rho_{10}}{dt} = -(\gamma_A + i\omega_A)\rho_{10} - \gamma Me^{-2i\omega_L t}\rho_{01} - iV_{10}(\rho_{11} - \rho_{00}) \tag{1c}$$

where γ is the spontaneous decay rate of the atom in a normal vacuum, and $\gamma_A = \gamma(N + \frac{1}{2})$, and N and M are parameters that describe the squeezing with the relationship $|M|^2 \leq N(N+1)$. The interacting energy is given by

$$V_{10} = V_{01}^{*} = \mu_{10}(E_L e^{-i\omega_L t} + E_p e^{-i\omega_p t}) \tag{2}$$

with μ_{10} being the dipole transition moment.

The equations of motion cannot readily be solved. In steady-state the off-diagonal density matrix element ρ_{10} exhibits harmonic oscillations not only at frequencies ω_L and ω_p, but at $n\omega_L \pm m\omega_p$ as well, where n and m are integers. However, if the probe beam is very weak, we can get a solution that is exact for the applied field E_L and is correct to the lowest order in the amplitude E_p of the weak probe field. Then we find that ρ_{10} oscillates at three dominant frequencies ω_L, ω_p and $2\omega_L - \omega_p$. Therefore, ρ_{10} in steady-state may be expressed in terms of Fourier amplitude, denoted by $\rho_{10}(\omega_i)$, as

$$\rho_{10}(t) = \rho_{10}(\omega_L)e^{-i\omega_L t} + \rho_{10}(\omega_p)e^{-i\omega_p t} + \rho_{10}(2\omega_L - \omega_p)e^{-i(2\omega_L - \omega_p)t} \tag{3}$$

where $\rho_{10}(\omega_L)$ and $\rho_{10}(\omega_p)$ are associated with atomic complex polarizations that yield index and absorption (or amplification) characteristics for the pump and probe modes. respectively,however, $\rho_{10}(2\omega_L - \omega_p)$ represents a wave-mixing response, which gives rise to generation of an optical wave with frequency $2\omega_L - \omega_p$.

Substituting Equs. (2) and (3) into Equ. (1) and keeping all orders in E_L and first order in E_p, the Fourier amplitudes can be solved algebraically as

$$\rho_{10}(\omega_L) = \frac{\gamma^2 M' + (2\gamma_A \gamma M' - \frac{1}{2}\Omega^2)(\rho_{11} - \rho_{00})^{dc}}{i\Omega(i\Delta - \gamma_A - \gamma M')} \tag{4a}$$

$$\rho_{10}(\omega_p) = \frac{\frac{i}{2}\Omega_p'(\gamma_A - i\delta + i\Delta)(\rho_{11} - \rho_{00})^{dc} + \frac{i}{2}\Omega(\gamma_A - i\delta + i\Delta + \gamma M')(\rho_{11} - \rho_{00})^{+\delta}}{\gamma^2 |M|^2 - (\gamma_A - i\delta)^2 - \Delta^2} \tag{4b}$$

$$\rho_{10}(2\omega_L-\omega_p)=\frac{\frac{i}{2}\Omega_p'^{*}\gamma M'(\rho_{11}-\rho_{00})^{dc}+\frac{i}{2}\Omega(\gamma_A+i\delta+i\Delta+\gamma M')(\rho_{11}-\rho_{00})^{-\delta}}{\gamma^2|M|^2-(\gamma_A+i\delta)^2-\Delta^2} \tag{4c}$$

where $(\rho_{11}-\rho_{00})^{dc}$ is the steady-state saturated populatation inversion induced by the applied field

$$(\rho_{11}-\rho_{00})^{dc}=\frac{\gamma(\gamma^2|M|^2-\gamma_A^2-\Delta^2)}{-2\gamma_A(\gamma^2|M|^2-\gamma_A^2-\Delta^2)+\Omega^2(\gamma_A+\gamma|M|cos\Phi)} \tag{5}$$

and $(\rho_{11}-\rho_{00})^{\pm\delta}$ represents a coherent population pulsation yielded by the probe polarization, oscillating harmonically at the pump-probe beat frequency

$$(\rho_{11}-\rho_{00})^{+\delta}=\frac{i\Omega_p'\{[\gamma^2|M|^2-(\gamma_A-i\delta)^2-\Delta^2]\rho_{10}^{*}(\omega_L)-\frac{i}{2}\Omega(\gamma_A-i\delta+i\Delta+\gamma M'^{*})(\rho_{11}-\rho_{00})^{dc}\}}{(2\gamma_A-i\delta)[\gamma^2|M|^2-(\gamma_A-i\delta)^2-\Delta^2]-\Omega^2(\gamma_A-i\delta+\gamma|M|cos\Phi)} \tag{6a}$$

$$(\rho_{11}-\rho_{00})^{-\delta}=\left[(\rho_{11}-\rho_{00})^{+\delta}\right]^{*} \tag{6b}$$

where $\Omega=2|\mu_{10}E_L|, \Omega_p'=2|\mu_{10}E_p|e^{i(\phi_p-\phi_L)}$, ϕ_p and ϕ_L are the phases of the probe and pump fields, respectively, without loss of generalities, we take $\phi_p-\phi_L=0$ thereafter. $M'=|M|e^{-i\Phi}$ with $\Phi=2\phi_L-\phi$ being the relative phase between the applied field and the squeezed vacuum. $\Delta=\omega_L-\omega_A$ and $\delta=\omega_p-\omega_L$ are the detunings between the applied laser field (or the squeezed vacuum) and the atom, and the probe.

It is shown from Eqs. (4b) and (4c) that the pump mode and the squeezed mode at the central frequence ω_L are scattered into the sideband modes at frequencies $\omega_L\pm\delta$ (i.e., the probe and the nonlinear response) by the population pulsations $(\rho_{11}-\rho_{00})^{\pm\delta}$, respectively.

III. Probe Absorption Spectrum

In optical media involving resonant absorption the linear susceptibility $\chi(\omega_p)$ of the probe field at frequency ω_p may be expressed in terms of the off-diagonal elements of the density matrix ρ_{10}

$$\chi(\omega_p) = \frac{\mathcal{N}\mu_{10}}{\epsilon_0 E_p}\rho_{10}(\omega_p) \tag{7}$$

where $\mathcal{N}$ is the number density of atoms, ϵ_0 being the free space permitivity.

The real and imaginary parts of the susceptibility make contributions to the refractive index and absorption coefficient, respectively. Therefore, the scaled probe absorption spectrum may be defined as

$$W(\omega_p) = Im\left[\frac{\rho_{10}(\omega_p)}{\Omega'_p}\right] \tag{8}$$

A. In the Absence of a Classically Driven Field

If there is no driven field, the probe response may read

$$\rho_{10}(\omega_p) = \frac{-i\Omega'_p\gamma(\gamma_A - i\delta + i\Delta)}{4\gamma_A[\gamma^2|M|^2 - (\gamma_A - i\delta)^2 - \Delta^2]} \tag{9}$$

By inspecting, it is easily found that there is a threshold at $\gamma|M| = |\Delta|$, where the probe absorption has different characters below and above the point.

Below threshold, $\gamma|M| < |\Delta|$, the absorption spectrum takes the form

$$W(\omega_p) = \frac{\gamma}{8\mu}\left[\frac{\Delta + \mu}{\gamma_A^2 + (\mu + \delta)^2} - \frac{\Delta - \mu}{\gamma_A^2 + (\mu - \delta)^2}\right] \tag{10}$$

where $\mu = \sqrt{\Delta^2 - \gamma^2|M|^2}$

The absorption has two lines. One is centred at frequence $\omega_L - \mu$, and the other at $\omega_L + \mu$. Both have the same width. On the other hand, it is easy to see that the positions of both spectral components are shifted towards centre for $\frac{(r|M|)^2}{2\Delta}$ when $\Delta^2 \gg (r|M|)^2$ due to the presence of the squeezed vacuum. Figure 1 is the absorption spectra for some values of parameters N, M and Δ. It is worthwhile to note that the spectra show some resemblance to the Fano spectrum [15] and the probe transparency takes place at the frequence $2\omega_s - \omega_A$ when $\Delta \gg 1$ and $N \gg 1$ under the perfect squeezing. It is evident from eq. (10) that the first line always shows absorption and the second one always exhibits gain if we take $\Delta > 0$. When $N \gg 1$, the two spectra are broaden so that they overlap. Competition between the absorption and gain may lead to the suppression of the absorption at the frequence $2\omega_s - \omega_A$, thus the transparency occurs.

Above threshold, $\gamma|M| > |\Delta|$, the scaled absorption spectrum is

$$W(\omega_p) = \frac{\gamma}{4\gamma_A}\frac{1}{2\mu'}\left[\frac{\Delta\delta + \mu'(\mu' + \gamma_A)}{(\mu' + \gamma_A)^2 + \delta^2} - \frac{\Delta\delta + \mu'(\mu' - \gamma_A)}{(\mu' - \gamma_A)^2 + \delta^2}\right] \tag{11}$$

where $\mu' = \sqrt{\gamma^2|M|^2 - \Delta^2}$

The absorption spectrum is composed of two spectral components which are located at centre, but with different widths. Each component can be separated into dispersive part from absorptive part. It is the dispersive characters that give rise to the Fano-like profiles in the absorption spectra. See fig. 2. The probe transparency is also at frequency $2\omega_s - \omega_A$.

In fact, when $|M| = \sqrt{N(N+1)}$, the absorption spectrum value at $2\omega_s - \omega_A$ is

$$W(\omega_p = 2\omega_s - \omega_A) = \frac{1}{1 + [4\Delta(2N+1)]^2} \tag{12}$$

Thus $W(2\omega_s - \omega_A) \approx 0$ when $\Delta \gg 1$ and $N \gg 1$, regardless of below or above threshold, that is, that the reponse of the medium to the probe mode is transparent.

Here is in a position to look at the role of the squeezed vacuum in its interaction with an atom. On the one hand, the squeezed vacuum modifies dramatically the decay rates of the two quadrature components of the polarization [2]. One rate is very small and the other (including atomic inversion) becomes extremely large in a strong completely squeezing limit. That is, that one quadrature is essentially unchanged, while the other decays rapidly to zero and the inversion to its stationary value $-\frac{1}{2N+1}$. On the other hand, from Eqs. (4c) and (6b), the wave-mixing reponse reads

$$\rho_{10}(2\omega_L - \omega_p) = i\frac{\gamma^2 M'\Omega_p'^*}{4\gamma_A[(\gamma_A + i\delta)^2 + \Delta^2 - \gamma^2|M|^2]} \tag{13}$$

which is proportional to the squeezed parameter M. As we know, in the standard vacuum, the nonlinear reponse occurs only in the presence of a driven field. So Eq. (13) may mean that a squeezed vacuum has some of the characteristics of the driving field.

We may qualitatively interprete the predicted effect from the viewpoint of the role of the driving-like field of the squeezed vacuum. When the squeezed vacuum interacts with a two-level atom, it creates two virtual levels one at $\omega_A + \Delta$ above the excited level $|1>$, and the other at $\omega_A + \Delta$ below the ground level $|0>$, where the detuning $\Delta = \omega_s - \omega_A$, between the squeezed vacuum and the atom, is assumed to be more than zero. The photons from the squeezed vacuum will be scattered inelastically at the two frequencies $\omega_A + 2\Delta$ and ω_A. Scattering of a photon at the first frequency is accompanied by the absorption of a pair of correlated squeezed photons. The process leaves the atom in the excited state, from which it may spontaneously emit a photon at the second frequency. Since the population in the ground level is slightly more than that in the excited state, the line at ω_A will show absorption, while the line at $\omega_A + 2\Delta$ can show gain. So the net absorption at $\omega_A + 2\Delta$ will decrease. The transparence can occur in certain cases.

B. In the Presence of a Coherently Driven Field

If the atom is driven by a laser field, the absorption can be written as, making use of Eqs.(4)-(6).

$$W(\omega_p) = Im\left[\frac{E(\delta)}{D(\delta)}\right] \tag{14}$$

where

$$E(\delta) = \frac{1}{2}\Omega(\gamma_A - i\delta + i\Delta + \gamma M')\rho_{10}^*(\omega_L) - \frac{i}{2}\left[(2\gamma_A - i\delta)(\gamma_A - i\delta + i\Delta) + \frac{1}{2}\Omega^2\right](\rho_{11} - \rho_{00})^{dc} \tag{15a}$$

$$D(\delta) = (2\gamma_A - i\delta)[(\gamma_A - i\delta)^2 + \Delta^2 - \gamma^2|M|^2] + \Omega^2(\gamma_A - i\delta + \gamma|M|cos\Phi) \tag{15b}$$

The characteristics of the absorption spectrum is determined by the roots of $D(\delta)$. In general it is difficult to get its roots. Here our main interest is of case of a strong minimium-uncertainty squeezed vacuum , i.e. $N \gg 1$, $|M| \approx N + \frac{1}{2}$, then

$$D(\delta) = (2\gamma_A - i\delta)[(\gamma_A - i\delta)^2 + \Delta^2 + \Omega^2 - \gamma^2|M|^2] - \gamma_A\Omega^2(1 - cos\Phi) \tag{16}$$

If we take $\phi = 0$, $D(\delta)$ has three roots as follows

$$\delta_0 = -i2\gamma_A \tag{17a}$$

$$\delta_\pm = i\left(-\gamma_A \pm \sqrt{(\gamma|M|)^2 - (\Delta^2 + \Omega^2)}\right) \tag{17b}$$

This shows the absorption spectrum, similar to resonance fluorescence spectrum [3], has a three-resonance structure. The width of the central component is $2\gamma_A$, nevertheless the positions and widths of the other depend on values of the factor $\sqrt{(\gamma|M|)^2 - (\Delta^2 + \Omega^2)}$. When $(\gamma|M|)^2 > (\Delta^2 + \Omega^2)$, the factor contributes a real value, corresponding two resonance components are also located at centre, but with different widths. However, when the factor is an imaginary number, the other two resonances are separated from centre. There is a threshold in parameter $|M|$, which determines the spectral structure.

Below threshold, $(\gamma|M|)^2 < (\Delta^2 + \Omega^2)$, the absorption spectrum takes the form

$$W(\omega_p) = \frac{a_0(q_0 + \delta)^2 + b_0}{(2\gamma_A)^2 + \delta^2} + \frac{a_+(q_+ + \delta + \nu)^2 + b_+}{\gamma_A^2 + (\delta + \nu)^2} + \frac{a_-(q_- + \delta - \nu)^2 + b_-}{\gamma_A^2 + (\delta - \nu)^2} \tag{18}$$

where

$$\nu = \sqrt{\Delta^2 + \Omega^2 - (\gamma|M|)^2} \tag{19a}$$

$$a_0 = \frac{\gamma(\Delta^2 - \Omega^2)}{4(\Delta^2 + \Omega^2)^2} \tag{19b}$$

$$q_0 = \frac{\Delta\Omega^2}{2(\Delta^2 - \Omega^2)} \tag{19c}$$

$$b_0 = a_0 \left[4\gamma_A^2 - \left(\frac{\Omega^2}{2(\Delta^2 - \Omega^2)} \right)^2 \right] \tag{19d}$$

$$a_\pm = \mp \frac{\gamma}{8\nu\,(\Delta^2 + \Omega^2)^2} \left[\Delta(\Delta^2 + \frac{1}{2}\Omega^2) \pm \nu(\Delta^2 - \Omega^2) \right] \tag{19e}$$

$$q_\pm = \frac{(\nu^2(\Delta^2 - \frac{1}{2}\Omega^2) \pm \nu\Delta^3 + \frac{1}{2}\gamma_A^2\Omega^2)}{\left[\mp\nu(\Delta^2 - \Omega^2) - \Delta(^2 + \frac{1}{2}\Omega^2)\right]} \pm \nu \tag{19f}$$

$$b_\pm = \pm \frac{1}{2\nu(\Delta^2 + \Omega^2)} \left[-(q_\pm \mp \nu)^2 a_\pm - \frac{\gamma}{2}\Delta^3 \right] \tag{19g}$$

The spectrum can be decomposed of three parts, the central spectrum and two Rabi sideband spectra, where the width of the central one is twice of the two sideband ones. Each may be disconnected into a Fano spectrum [15] and an absorption spectrum. However, since each width is very large so that they overlap, the total absorption does not display a well-separated three-resonance profile. It is interesting that the numerical results, plotted in Figs. 3, 4, have shown that the spectrum has shaped in a Fano-like profile and a zero- absorption at frequence $2\omega_L - \omega_A$, in a large atom-pump detuning limit.

Above threshold, $(\gamma|M|)^2 > (\Delta^2 + \Omega^2)$, the absorption spectrum can be also decomposed of a Fano spectrum and an absorption spectrum,

$$W(\omega_p) = \frac{a_0(q_0 + \delta)^2 + b_0}{(2\gamma_A)^2 + \delta^2} + \frac{A_+(Q_+ + \delta)^2 + B_+}{(\gamma_A + \nu')^2 + \delta^2} + \frac{A_-(Q_- + \delta)^2 + B_-}{(\gamma_A - \nu')^2 + \delta^2} \tag{20}$$

where

$$\nu' = \sqrt{(\gamma|M|)^2 - (\Delta^2 + \Omega^2)(21a)}$$

$$A_{\pm} = \frac{\gamma(\gamma_A \pm \nu')}{4\nu'(\Delta^2 + \Omega^2)^2} \left[\frac{\nu'\Delta^2}{2\gamma_A} \mp (\Delta^2 - \frac{\Omega^2}{2}) \right] \quad (21b)$$

$$Q_{\pm} = \frac{\Delta \left[\pm\nu'(\Delta^2 + \frac{1}{2}\Omega^2) - \gamma_A(\Delta^2 - \frac{1}{2}\Omega^2) \right]}{2\left[2\gamma_A(\Delta^2 - \frac{1}{2}\Omega^2) \mp \nu'\Delta^2 \right]} \quad (21c)$$

$$B_{\pm} = A_{\pm} \left[-Q_{\pm}^2 + \frac{2\gamma_A^2(\gamma_A \pm \nu')(\Delta^2 - \Omega^2)}{2\gamma_A(\Delta^2 - \frac{1}{2}\Omega^2) \mp \nu'\Delta^2} \right] \quad (21d)$$

Similar to below threshold, the spectrum may also consist of three components, but they are located at the centre with different widths. Figs. 5 and 6 display absorption spectra above threshold in a large detuning case. They also shape Fano-like profiles, and a probe transparency takes place at frequence $2\omega_L - \omega_A$.

As we know, when an atomic system is driven far-off resonance ($\Delta \gg \gamma$) in the standard vacuum, the absorption spectrum has well-separated three lines, centred at $\omega_A + 2\Delta$, ω_L and ω_A, respectively, which have a clear physical interpretation. However, instead of a strongly squeezed vacuum, although the spectrum can also be written in a three resonance form, the absorption does not show a well-separated three lines since each width is very large so that they overlap. Due to the competition among them, it exhibits a Fano profile and transparency effect finally, so long as the detuning Δ is not much more than the modified decay rate γ_A. This effect is not so easy to understand in the usual literature. Perhaps we should bravely make an assumption to interpret it. That is a strong squeezed vacuum would induce an atom to appear a continuous structure, due to its driving-field-like role and its modifing effect on the atomic decay. When the atom is illuminated by the probe light, it can directly jump from the ground state into the excited state via absorbing a probe photon, or transits into the continuum through absorbing a pair of correlated photon from the squeezed vacuum, then emits a photon and goes down the excited state finally.

The two pathways may interfere which may lead to zero absorption under certain conditions, i.e. the probe transparency.

It is worth to note that when $\Delta \gg 1$ and $N \gg 1$, the spectra are not obviously sensitive to the phase between the driving field and the squeezed vacuum, nevertheless, the variations of the phase only slightly modify the magnitude of the spectra. See Figs. 7 and 8.

IV. Summary

In this paper we have studied absorption spectra of a two-level atom in a squeezed vacuum. In the large squeezing and detuning limit, there is a threshold for the squeezed parameter M in absorption spectra. Below threshold, the position of the absorption peak is at one of Rabi sidebands, however, the absorption peak is located at centre above threshold. The probe spectra shape in Fano-like profiles and it is transparent for the probe at frequence $2\omega_L - \omega_A$ in both cases. The effect can appear in absence of and in presence of a driving field. It also is not phase-sensitive. An interpretation of this effect may be based on a braved assumption of which the strong squeezing may induce a continuum near atomic excited energy level in the large detuning limit, which emerges a new atomic absorption tunnel, which may interfere with the normal one. This interference can give rise to zero probe absorption under certain conditions.

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References

[1] H. J. Kimble, Phys. Rep., 219, 227(1992); and references therein.

[2] C.W.Gardiner, Phys. Rev. Lett. 56, 1917(1986).

[3] H.J.Carmichael, A.S.Lane and D.F.Walls, Phys. Rev. Lett., 58, 2539(1987); J. Mod. Opt., 34, 821(1987).

[4] B.R.Mollow, Phys. Rev., 188, 1969(1969).

[5] S. Smart and S. Swain, Phys. Rev. A 48, R50(1993); Opt. Commun., 99, 369(1993); S. Swain, "Anomalous resonance fluorescence spectra in a squeezed vacuum" (???).

[6] S. Smart and S. Swain, Phys. Rev. A 45, 6857(1992); 6863(1992); Quantum Opt., 5, 75(1993).

[7] A. S. Parkins and C. W. Gardiner, Phys. Rev. A 40, 3796(1989); A. S. Parkins, P. Zoller and H. J. Carmichael, Phys. Rev. A 48, 758(1993); P. R. Rice and L. M. Pedrotti, J. Opt. Soc. Am. B 9, 2008(1992).

[8] H. Ritsch and P.Zoller, Opt. Commun., 64, 523(1987); 66, 333(1988); Phys. Rev. A 38, 4657(1988).

[9] S. An, M. Sargent III and D. F. Walls, Opt. Commun., 67, 373(1988); S. An and M. Sargent III, Phys. Rev. A 39, 3998(1989).

[10] S. Swain and P. Zhou, in preparation.

[11] Z. Ficek, W. Smyth and S. Swain, Opt. Commun., (to be published).

[12] Z. Ficek and B.J.Dalton, Opt. Commun., 102, 231 (1993).

[13] P. L. Knight, Comments At. Mol. Phys., 15, 193(1984); K. M. Gheri, C. Saavedra and D. F. Walls, Phys. Rev. A 48, 3344(1993); M. O. Scully and M. Fleischhauer, Phys. Rev. Lett., 69, 1360(1992).

[14] S. E. Harris, J. E. Field and A. Imamoglu, Phys. Rev. Lett., 64, 1107(1990); M. O. Scully, Phys. Rev. Lett., 67, 1855(1991); M. Fleischhauer, C. H. Keital, M. O. SCully, C. Su, B. T. Ulrich and S. Y. Zhu, Phys. Rev. A 46, 1468 (1992); K. Hakuta, L. Marmet and B. P. Stoicheff, Phys. Rev. A 45, 5152(1992).

[15] U. Fano, Phys. Rev. A 124, 1866(1961).

FIGURE CAPTIONS

- Fig. 1: Probe absorption spectrum of a two-level atom damped by a broadband perfectly squeezed vacuum, in the absence of a classically driving field, below threshold, with parameters $\Delta = 2, N = 1, 1(a)$; $\Delta = 5, N = 1, 1(b)$; $\Delta = 15, N = 10, 1(c)$; $\Delta = 20, N = 10, 1(a)$.

- Fig. 2: Same as Fig. 1, but above threshold, with parameters $\Delta = 1, N = 1, 2(a)$; $\Delta = 5, N = 10, 2(b)$; $\Delta = 10, N = 10, 2(c)$; $\Delta = 15, N = 20, 2(a)$.

- Fig. 3: Probe absorption spectrum of a coherently driven two-level atom damped by a broadband perfectly squeezed vacuum, below threshold, with parameters $\Omega = 1, \phi = 0, N = 10$, and $\Delta = 11, 3(a)$; $\Delta = 15, 3(b)$; $\Delta = 20, 3(c)$; $\Delta = 30, 3(d)$.

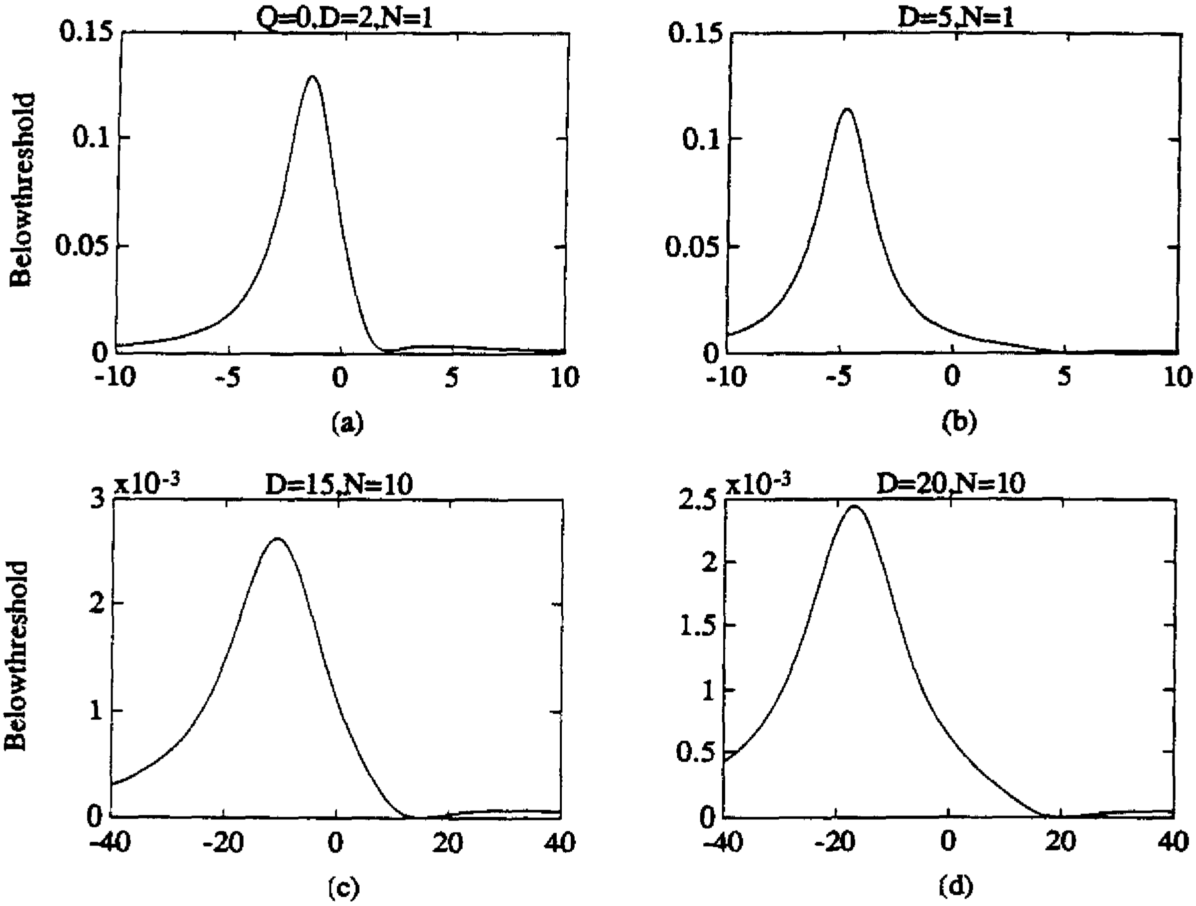
- Fig. 4: Same as Fig. 3, but with parameters $\Omega = 1, \phi = 0, \Delta = 15$, and $N = 1, 4(a)$; $N = 5, 4(b)$; $N = 10, 4(c)$; $N = 14, 4(d)$.

- Fig. 5: Same as Fig. 3, but above threshold, with parameters $\Omega = 2, \phi = 0, N = 20$, and $\Delta = 2, 5(a)$; $\Delta = 10, 5(b)$; $\Delta = 15, 5(c)$; $\Delta = 20, 5(d)$.

- Fig. 6: Same as Fig. 5, but with parameters $\Omega = 2, \phi = 0, \Delta = 10$, and $N = 10, 6(a)$; $N = 15, 6(b)$; $N = 20, 6(c)$; $N = 50, 6(d)$.

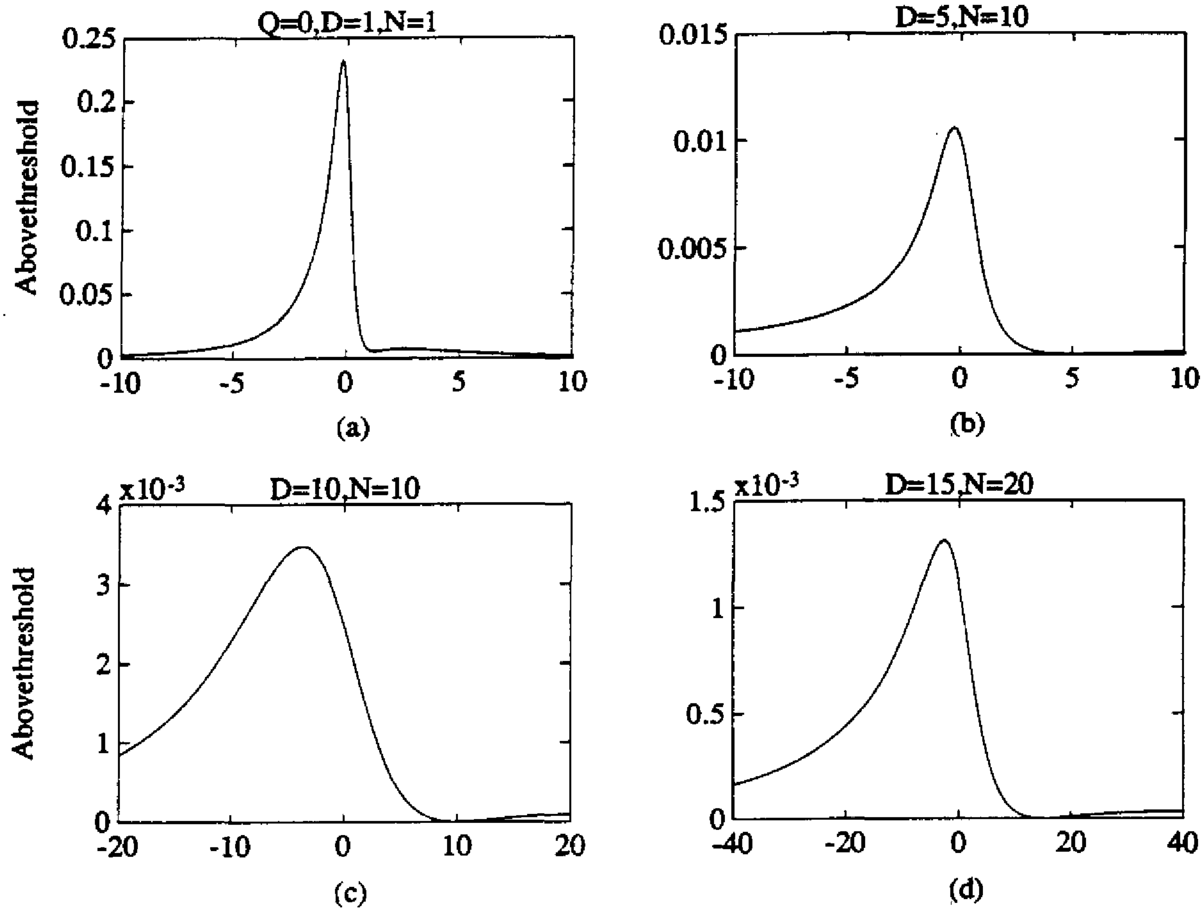
- Fig. 7: Same as Fig. 3, but with parameters $\Omega = 2, \Delta = 20, N = 10$, and $\phi = 0, 7(a)$; $\phi = \frac{1}{4}\pi, 7(b)$; $\phi = \frac{1}{2}\pi, 7(c)$; $\phi = \pi, 7(d)$.

- Fig. 8: Same as Fig. 5, but with parameters $\Omega = 2, \Delta = 10, N = 20$, and $\phi = 0, 8(a)$; $\phi = \frac{1}{4}\pi, 8(b)$; $\phi = \frac{1}{2}\pi, 8(c)$; $\phi = \pi, 8(d)$.



Q=0,D=2,N=1
D=5,N=1
D=15,N=10
D=20,N=10
Belowthreshold
Belowthreshold
x10-3
x10-3
(a)
(b)
(c)
(d)

Fig.1



Q=0,D=1,N=1
0.25
0.2
0.15
0.1
0.05
0
Abovethreshold
-10
-5
0
5
10
(a)
D=5,N=10
0.015
0.01
0.005
0
-10
-5
0
5
10
(b)
4
x10-3
D=10,N=10
3
2
1
0
Abovethreshold
-20
-10
0
10
20
(c)
1.5
x10-3
D=15,N=20
1
0.5
0
-40
-20
0
20
40
(d)

Fig.2

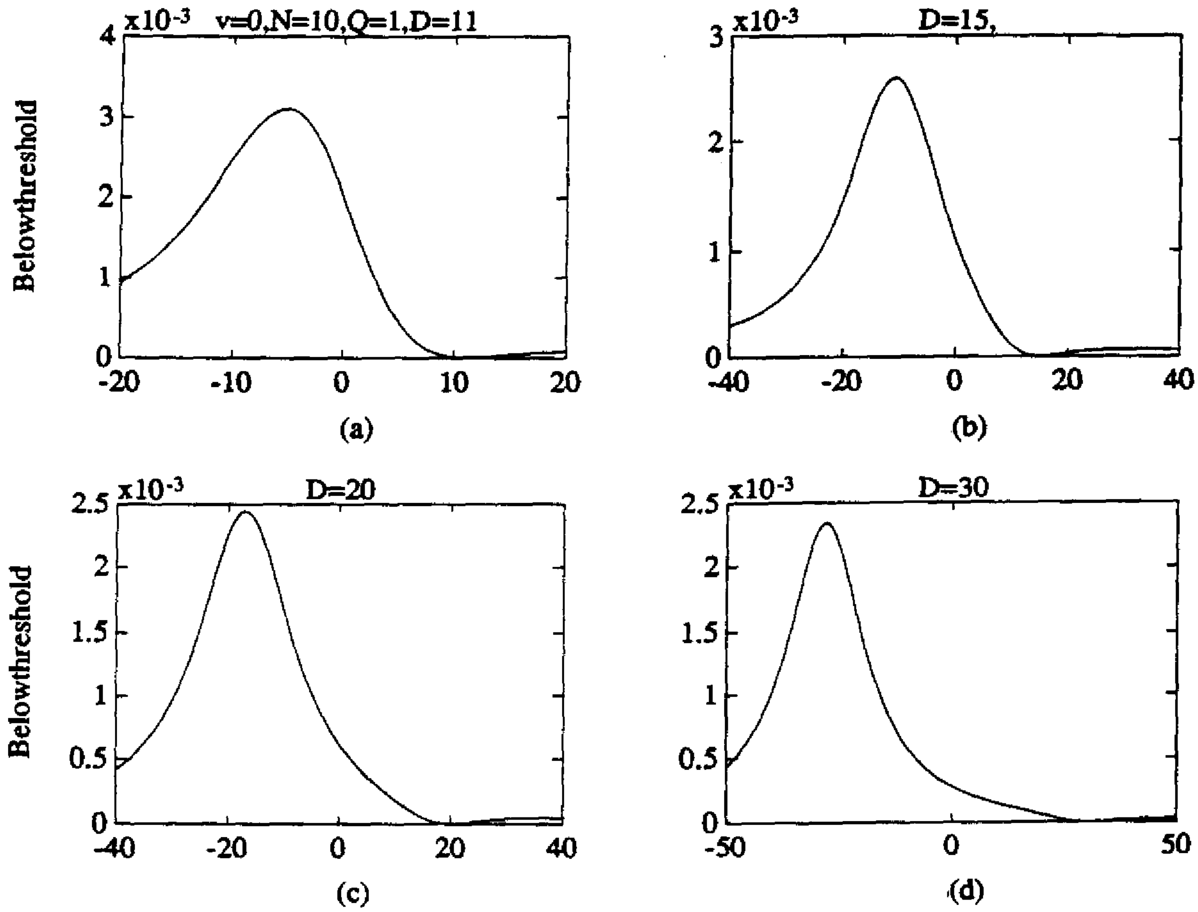
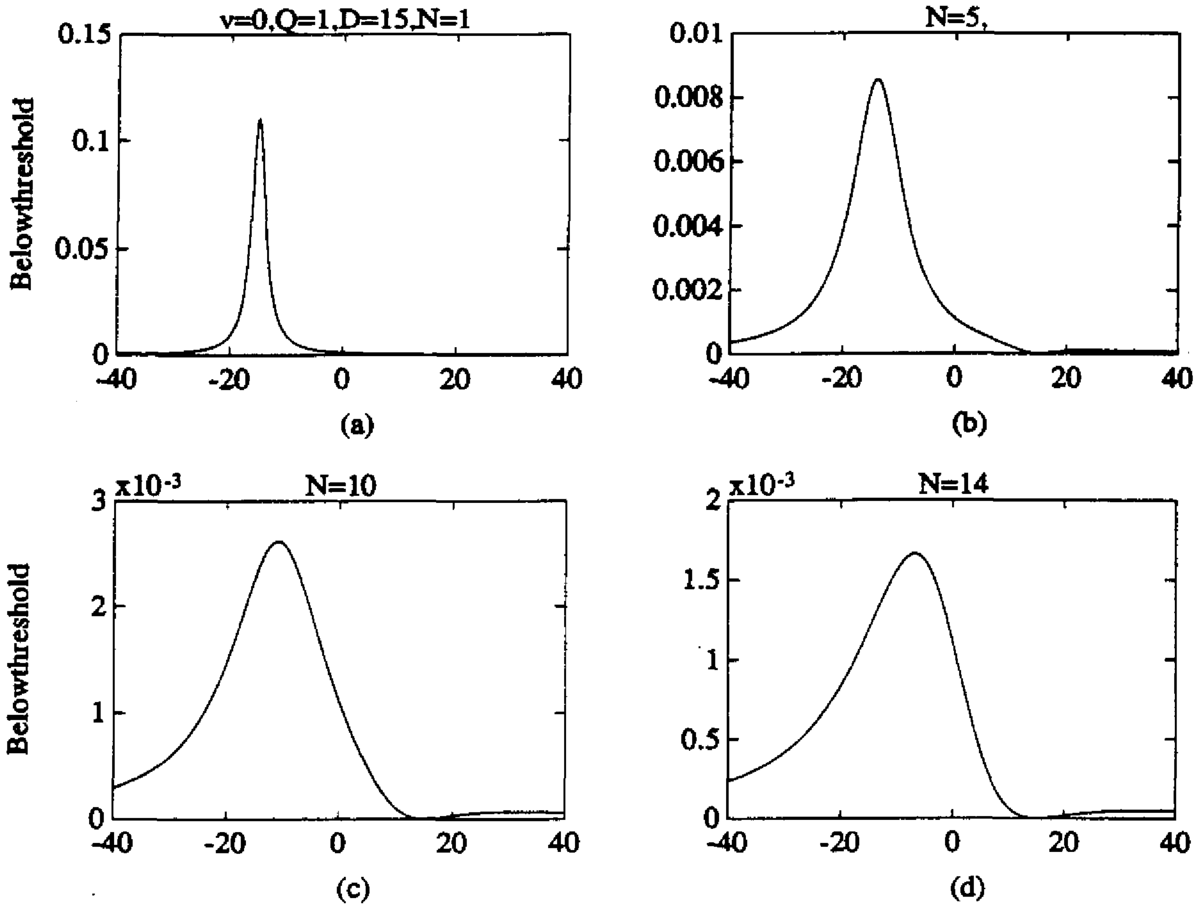

x10-3 v=0,N=10,Q=1,D=11
Belowthreshold
4
3
2
1
0
-20
-10
0
10
20
(a)
x10-3 D=15,
3
2
1
0
-40
-20
0
20
40
(b)
x10-3 D=20
Belowthreshold
2.5
2
1.5
1
0.5
0
-40
-20
0
20
40
(c)
x10-3 D=30
2.5
2
1.5
1
0.5
0
-50
0
50
(d)

Fig.3



v=0,Q=1,D=15,N=1
Belowthreshold
0.15
0.1
0.05
0
-40
-20
0
20
40
(a)
N=5,
0.01
0.008
0.006
0.004
0.002
0
-40
-20
0
20
40
(b)
x10-3
N=10
Belowthreshold
3
2
1
0
-40
-20
0
20
40
(c)
x10-3
N=14
2
1.5
1
0.5
0
-40
-20
0
20
40
(d)

Fig.4

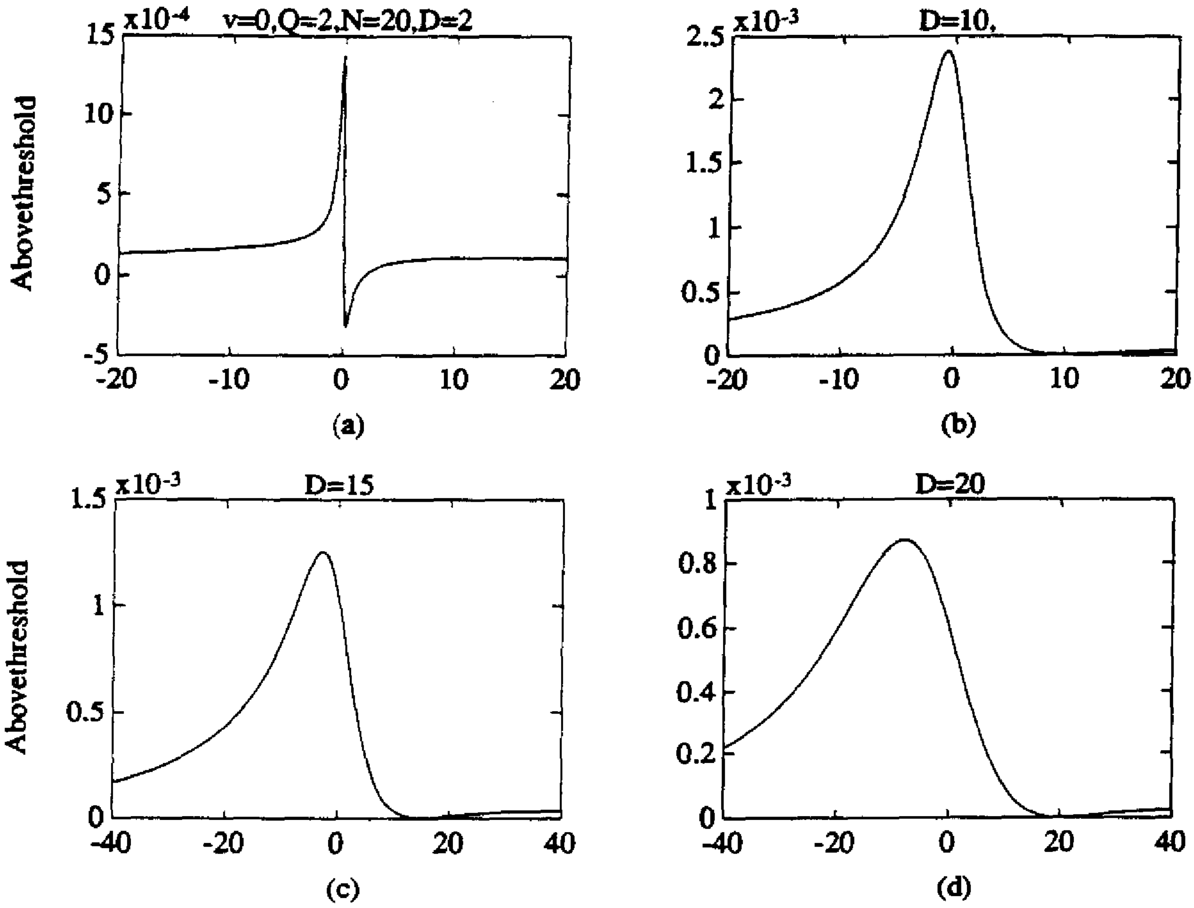

x10-4
v=0,Q=2,N=20,D=2
Abovethreshold
15
10
5
0
-5
-20
-10
0
10
20
(a)
x10-3
D=10,
2.5
2
1.5
1
0.5
0
-20
-10
0
10
20
(b)
x10-3
D=15
Abovethreshold
1.5
1
0.5
0
-40
-20
0
20
40
(c)
x10-3
D=20
1
0.8
0.6
0.4
0.2
0
-40
-20
0
20
40
(d)

Fig.5

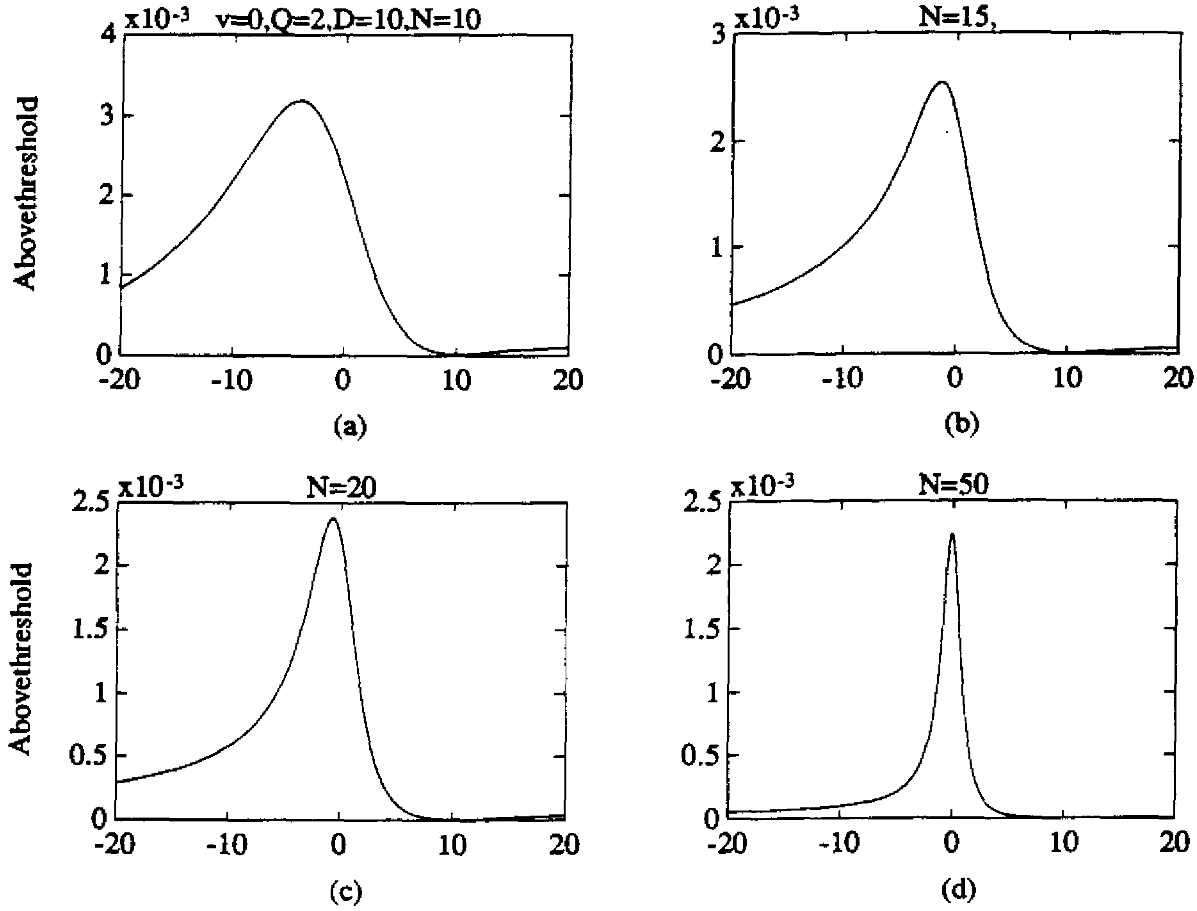

x10-3 v=0,Q=2,D=10,N=10
Abovethreshold
4
3
2
1
0
-20
-10
0
10
20
(a)
x10-3 N=15,
3
2
1
0
(b)
x10-3 N=20
Abovethreshold
2.5
2
1.5
1
0.5
0
(c)
x10-3 N=50
(d)

Fig.6

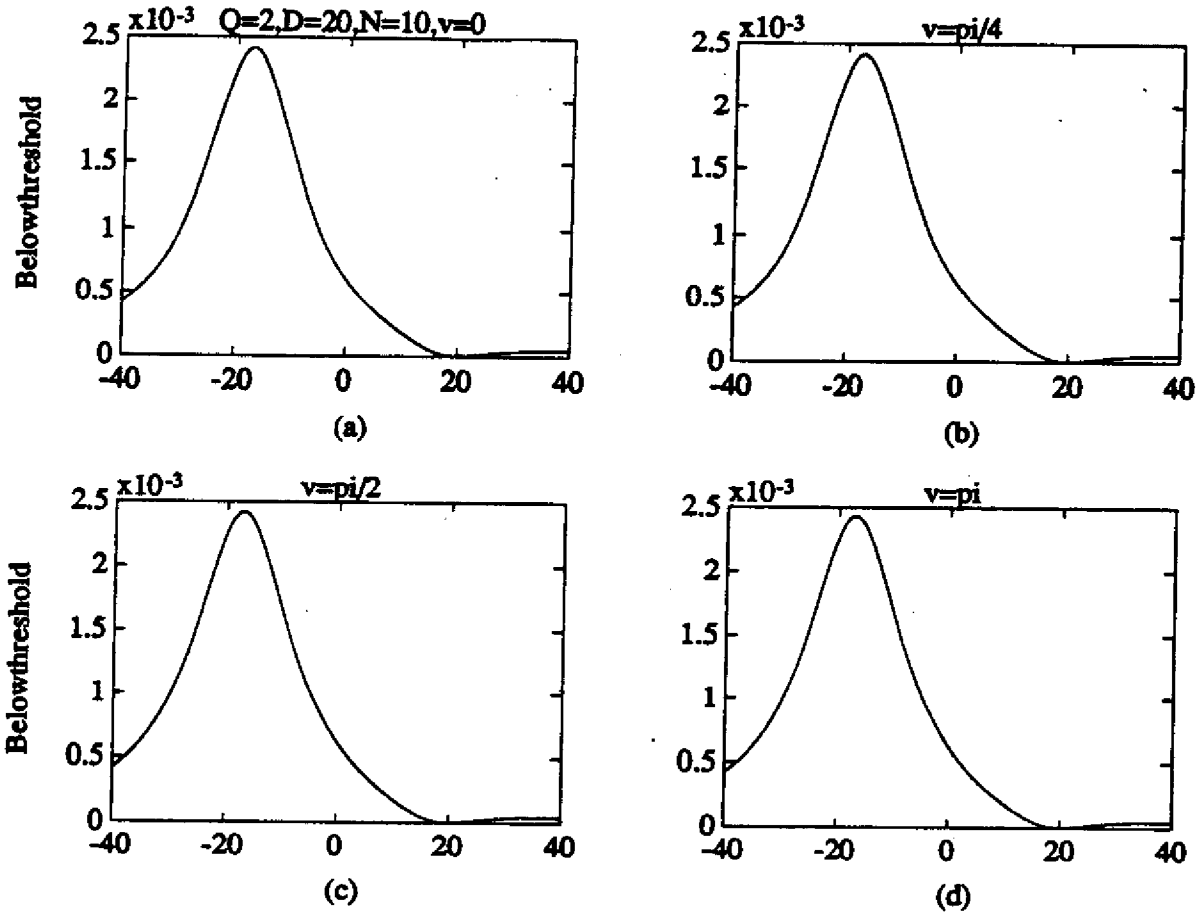

x10-3
Q=2,D=20,N=10,v=0
Belowthreshold
2.5
2
1.5
1
0.5
0
-40
-20
0
20
40
(a)
x10-3
v=pi/4
(b)
x10-3
v=pi/2
(c)
x10-3
v=pi
(d)

Fig.7

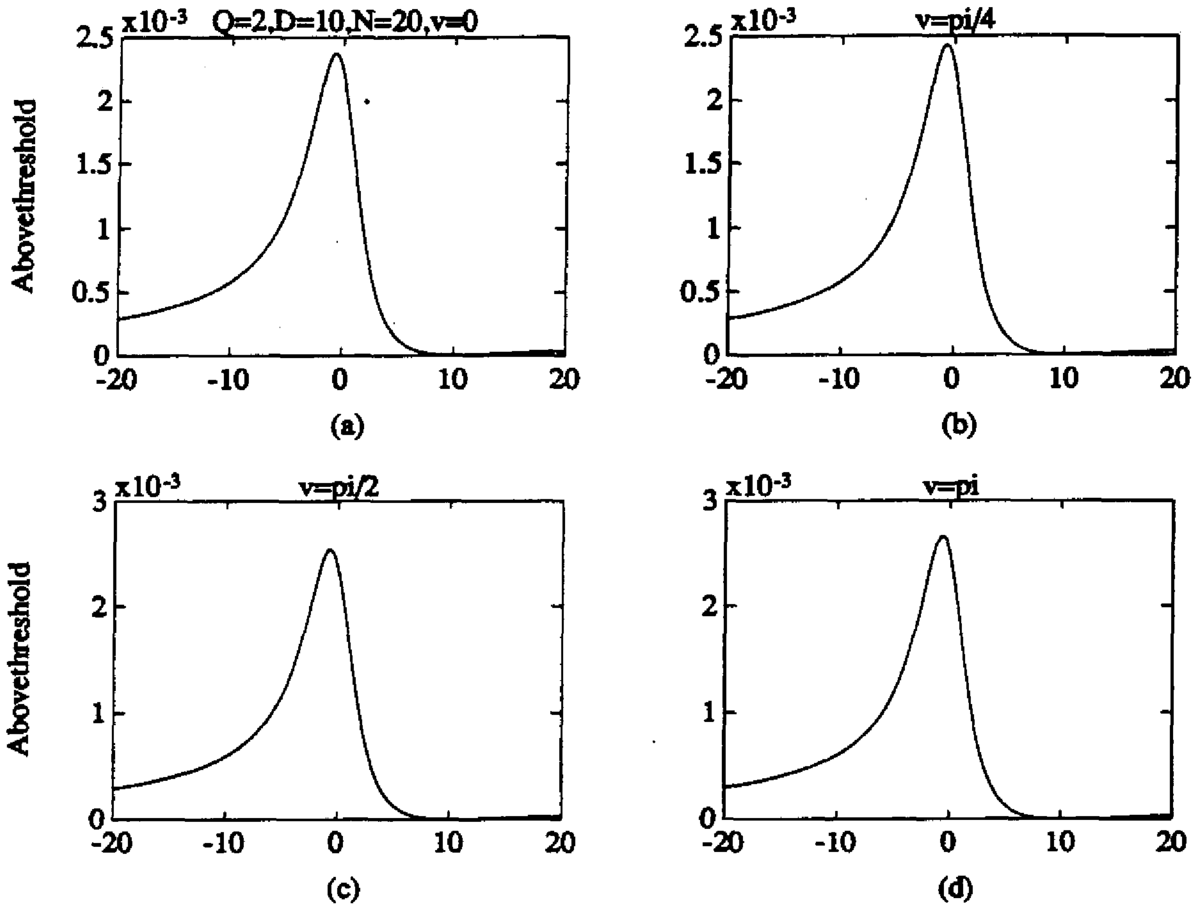

x10-3 Q=2,D=10,N=20,v=0
Abovethreshold
(a)
x10-3 v=pi/4
(b)
x10-3 v=pi/2
Abovethreshold
(c)
x10-3 v=pi
(d)

Fig.8