PHYSICAL REVIEW A VOLUME 55, NUMBER 3 MARCH 1997

Atom-atom correlations via radiation-field coupling

Yves Gontier
Commissariat à l'Energie Atomique, Service des photons et des Molécules, Centre d'Etudes de Saclay,
91191 Gif-sur-Yvette Cedex, France

(Received 25 October 1995; revised manuscript received 15 October 1996)

It is shown that two identical systems, not interacting with each other but resonantly coupled with the same intense radiation field, can be correlated. The case of a pair of two-level atoms is quantitatively investigated by calculating the probability for the emission of a photon by one of the two atoms. One finds that the presence of the other atom strongly affects the shape of the resonance curve. It is concluded that this effect, called signal transmission by means of optical correlations, if observable, could be the way for transmitting information between the atoms without resorting to any modulation of the field amplitude. [S1050-2947(97)07503-3]

PACS number(s): 42.50.Ct, 42.50.Fx, 42.50.Lc, 32.80.Wr

I. INTRODUCTION

The correlations between distant particles have provided many subjects of active research for the last 50 years or more. In the most general case, the correlations are the way for a system to know the state of another remote system. The correlations between the elements of a physical system arise (i) when these elements interact with each other (overlapping of the wave functions); (ii) when they do not interact with each other but are coupled to a radiation field; (iii) when they interact with each other and with a radiation field. An example of the first case is provided by intra-atomic correlations in many-electron atoms. The second situation, which is the topic of interest throughout this paper, is typically that of cooperative effects, while the third case can be illustrated by atomic or molecular collisions in the presence of a radiation field. By considering the last two cases involving physical systems (hereafter represented by atoms) strongly coupled to a radiation field, one is faced with the problem of space dependence. In this respect, an important class of correlations concerns those that are drastically space dependent. This dependence is a consequence of the dipole-dipole interaction between two or several atoms while correlations are produced by the exchange of the photons generated by the decay of the atoms themselves. In the theory of dispersion forces, it is shown that this last process makes the lifetime of any excited atom strongly affected by the presence of the other atoms [1].

The present paper is mainly concerned with correlations that are independent of space. They come from exchanges of photons between atoms either confined within volumes whose dimensions are small compared to the wavelength of the field or separated from each other by multiples of the wavelength. In both cases, all the atoms are exposed to the same field. The former model has proved efficient in predicting the presence of superradiant states among the complete set of states built up from the combination of individual states of excited two-level atoms [2].

The starting point of our study is an interpretation of cooperative emission in terms of photon exchanges. Thus, a photon emitted by any atom (say atom No. 1) is reabsorbed by a neighboring atom (atom No. 2) or stimulates the emission of a photon by the atom No. 2 according to whether this last atom is in the lower state or in the upper (excited) state. The process can be continued and the atom No. 2 can play the same role as the atom No. 1 for another atom No. 3, etc. Such absorption-emission processes are independent of space and involve the field generated by the atoms themselves via induced emission.

Here, we assume that the field to which the atoms are strongly coupled is that of an intense external radiation field (laser). In this case, the photons involved in the exchanges are supplied by the laser. At high intensity, induced emission prevails over spontaneous emission which can be neglected. The use of a laser to induce correlations is interesting in several respects: (i) it enables the maintenance of correlation during longer time intervals and over larger distances, and (ii) it provides a way to control the strength of the correlations by varying the intensity, the phase or the polarization of the beam at the positions of the atoms.

Our discussion holds for atoms contained in a cell that is much smaller than a wavelength of the field, as well as for atoms that are distant from each other by a multiple of the wavelength. These two cases are completely equivalent (such an event was not considered as unreasonable by Dicke in his early papers). We call reciprocal areas, the regions of space where identical atoms that are suitably phased interact with the same field. Any such area containing one or several atoms behaves like isolated atoms. Thus, all the reciprocal areas are correlated with each other in the same way as the atoms are correlated inside each one. Since the correlations are free of any spatial dependence, these areas can be mutually separated by arbitrarily large distances. However, owing to the fact that these correlations are due to the strict identity of the reciprocal areas one may expect, like for individual atoms, a perturbation of one of them can be felt by the others and recorded by suitable measurement devices. This effect, which we name signal transmission by means of optical correlations (STOC), if observable, may provide a way for transmitting information from a point to other distant points of the space without modulating the intensity of the laser beam. This is somewhat different from the case where the intensity of the field is so small that the depletion or the increase of the photon number caused by absorption-emission of anyone of the atoms influences the other atom-field interactions. A possible experimental realization of reciprocal areas consists of glass cells containing the same vapor and separated from each other by multiples of the radiation wavelength. The most important difficulties in such experiments are (i) to get the same field (amplitude, phase, polarization) at the cell positions, and (ii) to hold the atoms at well defined positions (reciprocal areas) in the laser beam where the ponderomotive forces appearing at high intensity tend to repel the atoms from their initial positions. This effect cannot be avoided but its consequences upon correlations are not so serious as they seem. Due to their inertia, the time the atoms spend to leave the reciprocal areas is larger than the duration of the quantum interactions. So, the correlations produced by such interactions have enough time to occur. However, the random motion of the atoms makes highly probable that some of them cross the reciprocal areas before being repelled. By this way the correlations are maintained by the time the pulse operates.

Notice that so far we have considered the case where the correlations are generated by the photons emitted by the atoms themselves or those provided by a laser field. This implies they disappear when all the atoms have decayed toward the lower state or when the laser is switched off. Such processes are to be distinguished from those involving correlations which persist after the interaction has ceased. The most famous example is provided by the *gedankenexperiment* of Bohm [3], which is an illustration of the problem raised by the paper of Einstein, Rosen, and Podolsky [4]. In this case the correlation is the consequence of events that occurred in the past and do not require the presence of a medium (laser) for the transmission of the information. These correlations, which are intimately related to the nonlocality of quantum mechanics, have been extensively discussed in the past and fall out of the scope of the present work.

We reduce the problem to that of two identical two-level atoms interacting with the same radiation field delivered by a laser. The theory is made within a fully quantized model in order to use our technique of resummation of perturbation series. To formulate mathematically the conditions giving rise to correlations, we introduce the spin-$\frac{1}{2}$ formalism in Sec. II A, while the resummation of the four-operator perturbation series is done by using the resolvent operator. In the numerical analysis of Sec. III, we calculate the probability for the resonant emission of a photon by the atom No. 1 in the presence of the atom No. 2, as a function of the energy gap of the two atomic levels. In Sec. III, an experimental model is discussed.

II. ATOM-ATOM CORRELATIONS

A. Spin-$\frac{1}{2}$ formalism

To begin, we consider the simple case of two two-level atoms which do not interact with each other but are resonantly coupled with a radiation field whose characteristics are the same as the positions of the atoms. This is typically the case of two neighboring atoms lying far apart in order to avoid any overlapping of their wave functions. Within the dipole approximation, we discard the space dependence of the operators. We assume that the interaction between the jth atom and the field proceeds via the exchange of a photon labeled k. This label k accounts for all the features of the field in a well-defined region of space. Any change of the field parameters will correspond to another value of k. Since two-level atoms are involved, it is convenient to use the formalism of spin-flip operators. Within this framework, we can formulate mathematically the conditions for correlation in terms of commutators of the atom-field operators.

The Hamiltonian of a two-atom system resonantly coupled with a quantized radiation field is (in atomic units)

$$H=\sum_{j=1}^{2} \omega_{0j}S_3^j+\sum_{j,k} (\alpha_k^j S_+^j a_k+\alpha_k^{j*}S_-^j a_k^+)+\sum_k \omega_k a_k^+ a_k , \tag{2.1}$$

where the rotating-wave approximation is made and $S_\pm^j$ are the spin-flip operators [5] obeying the following commutation relations:

$$[S_\pm^j ,S_\mp^{j'}]_-=2S_3^j\delta_{jj'} \tag{2.2}$$

and

$$[S_3^j ,S_\pm^{j'}]_-=\pm S_\pm^j\delta_{jj'} . \tag{2.3}$$

Although the operators $S_\pm$ are represented by matrices whose elements are independent of the atomic parameters, like Dicke [2], we associate the index j with these operators to designate the atoms.

In Eq. (2.1), the coefficient α_k^j is expressed in terms of the single-photon flux F'/F_0' (F'=flux/photon number, F_0=3.22×10^{34} cm^{-2} s^{-1} [6]) and the dipole matrix element corresponding to the jth atom with photon polarization $\vec{\varepsilon}_k$ is

$$\alpha_k^j=-i\left(\frac{F'}{F_0}\right)^{1/2} \omega_k^{1/2}{}_j\langle\pm|\vec{r}_j\cdot\vec{\varepsilon}_k|\mp\rangle_j , \tag{2.4}$$

where for clarity the upper and the lower states of the jth atom are denoted by $|+\rangle_j$ and $|-\rangle_j$, respectively.

The Hamiltonian of Eq. (2.1) can be rewritten into the following more convenient form:

$$H=H_0^A+H_F+\sum_{j=1}^{2}\sum_k (V_{j,k}^+ + V_{j,k}^-), \tag{2.5}$$

where

$$H_0^A=\omega_{01}S_3^1+\omega_{02}S_3^2, \tag{2.6}$$

$$H_F=\sum_k \omega_k a_k^+ a_k , \tag{2.7}$$

and

$$V_{j,k}^-=(V_{j,k}^+)^*=\alpha_k^j S_+^j a_k , \quad j=1,2, \tag{2.8}$$

where for generality, the energy-level separations are assumed to be different. From Eqs. (2.2) and (2.3) one obtains the commutation relations for the operators $V^\pm$

$$[V_{j,k}^\pm ,V_{j,k}^\pm]_-=0 \tag{2.9}$$

and

$$[V^{\pm}_{j,k}, V^{\mp}_{j',k'}]_- = \mp \alpha^{j}_{k} \alpha^{j'}_{k'} \{2S^{j}_{3} a_k a^{+}_{k'} \delta_{jj'} + S^{j}_{\pm} S^{j}_{\mp} \delta_{kk'}\}. \tag{2.10}$$

Equation (2.9) is unconclusive because the commutator always vanishes. We note that the first term on the right-hand side of Eq. (2.10) largely prevails over the second one because it is proportional to the intensity. For a monomode field, it does not vanish if $j' = j$ and $k' = k$, i.e., when identical atoms see the same field. This equation can be considered as being the mathematical formulation of the correlation, which comes from the fact that the photon absorption (emission) that is taking place in one of the two undistinguishable atom-field system is not independent of the emission (absorption) induced in the other system.

With these notations, we can build up all the possible naked states (in contrast to the dressed states) for a system of two correlated atoms. We find that these states are

$$|\mathrm{Correl}\rangle = \begin{cases} |a_1 a_2\rangle \\ \dfrac{1}{\sqrt{2}} (|a_1 b_2\rangle + |b_1 a_2\rangle). \\ |b_1 b_2\rangle \end{cases} \tag{2.11}$$

Equation (2.11) display the only three radiative states that can be built up from the four states of a couple of two-level atoms. The plots of the energy value of such a two-atom system interacting with a radiation field as a function of the single-atom energy-level separation provide an energy spectrum identical to that of a single three-level atom. It consists of two hyperbola branches located symmetrically on both sides of a straight line [8].

B. Resummed theory

Since the effect we study requires high-intensity radiation fields, we must resort to a nonperturbative model. To get reliable results, we make an exact resummation of a perturbation series. In the case of two noncommuting operators (absorption and emission operators of a single atom), the technique has been presented in detail previously [7]. Here we are faced with a more complicated situation since one has to handle four noncommuting operators (absorption and emission operators of two atoms). Such a problem has been solved for a general case [8]. For the sake of brevity, only the salient results will be recalled and truncated continued fractions will be used.

In general, the behavior of any system can be predicted once the time evolution operator is known. This operator can be calculated from the resolvent operator by means of the inversion integral

$$U(t) = \frac{1}{2\pi i} \oint e^{-izt} G(z) dz, \tag{2.12}$$

where

$$G(z) = \frac{1}{z - H}. \tag{2.13}$$

We refrain from presenting in more detail the resolvent theory, which is well known and widely utilized [9]. According to Eqs. (2.5) and (2.13), $G(z)$ can be expressed as

$$G(z) = G_0(z) + G_0(z) H_I G(z), \tag{2.14}$$

where

$$H_I = \sum_{j=1}^{2} (V^{+}_{j} + V^{-}_{j}) \tag{2.15}$$

and

$$G_0(z) = \frac{1}{z - (H^{AT}_{0} + H^{F})}. \tag{2.16}$$

In Eq. (2.15), the subscript k refers to the field state is dropped because in the operator expressions, it is assumed that the field is the same everywhere. The solution of Eq. (2.14) is obtained by an iteration technique which provides infinite series of increasing powers of the interaction H_I.

The problem one has to solve is to calculate the resonant emission probability of atom No. 1 in the presence of atom No. 2 (which also can resonantly absorb and emit photons of the field). Initially, the two atoms are assumed to be in their upper states and we consider the final state of the system where only atom No. 1 has decayed in its lower state (atom No. 2 remaining in its upper state). Thus, it is not necessary to symmetrize the initial and the final states of the system which are $|a_1, a_2, n\rangle$ *and* $|b_1, a_2, n+1\rangle$, respectively, i.e., atom No. 1 emits a photon by making a transition from the upper state a_1 toward the lower state b_1, atom No. 2 remains in its initial state a_2 while the photon number n is increased by one unit. Notice that the influence of atom No. 2 on atom No. 1 is independent of the choice of the states. It will be shown in foregoing accounts that the atoms can be correlated whatever the initial and final atomic states may be.

The matrix element one has to calculate is

$${}^{(k+1)}G_{b_1 a_2, a_1 a_2}(z) = \langle b_1 a_2, n+1 | {}^{(k+1)}G(z) | a_1 a_2, n\rangle, \tag{2.17}$$

where ${}^{(k+1)}G(z)$ is the operator describing the resonant emission of a photon by atom No. 1. As a result of the resummation, this operator can be expressed as

$${}^{(k+1)}G(z) = G(z) B G^{+}(z), \tag{2.18}$$

where

$$G(z) = \frac{1}{1 - \rho^{+}(z) - \rho^{-}(z)}, \tag{2.19}$$

$$G^{\pm}(z) = \frac{1}{1 - \rho^{\pm}(z)}, \tag{2.20}$$

and

$$\rho^{+}(z)=BG^{+}(z)A+YG^{+}(z)X, \tag{2.21}$$

$$\rho^{-}(z)=AG^{-}(z)B+XG^{-}(z)Y. \tag{2.22}$$

The absorption-emission operators A, B, X, and Y appearing in Eqs. (2.18), (2.21), and (2.22) are defined by $A=G_0V_1^{-}$, $B=G_0V_1^{+}$, $X=G_0V_2^{-}$, and $Y=G_0V_2^{+}$, respectively. To get the computational formulas, we replace $G^{\pm}(z)$ in Eqs. (2.21) and (2.22) by their values obtained from Eq. (2.20) by iteration, then substitute the expressions obtained into Eq. (2.19).

These expressions are other versions of the general formulas [8] which have been simplified to make the problem tractable. The loss of accuracy is negligible within the intensity range we consider since the major part of the contribution comes from the operators $\rho^{\pm}$.

III. QUANTITATIVE ANALYSIS

A. Analytic formulation

Equations (2.17)–(2.22) enable one to self-consistently compute to all orders the probability for the resonant net emission of a photon by atom No. 1. To make the calculation tractable without resorting to exaggerated computation time, we have retained, to each order, the most significant contributions. In addition, the number of iterations of the continued fractions has been limited to that ensuring the stability of the results. According to a previous account [9], we put the matrix element of Eq. (2.17) into a form characteristic of a two-level problem. The matrix element corresponding to the net emission of a photon by atom No. 1 resonantly coupled to the radiation field is

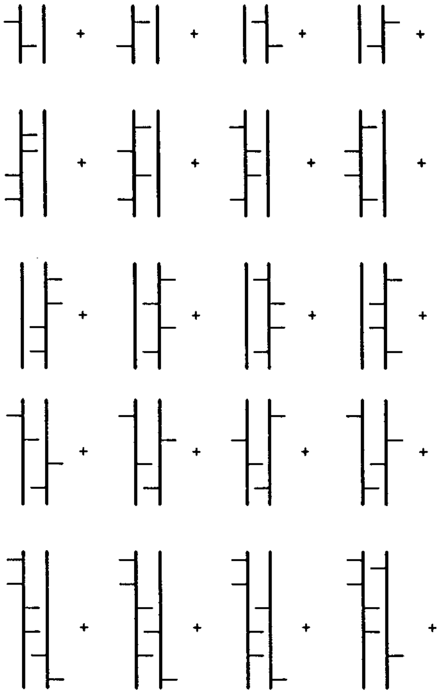

FIG. 1. Diagrammatic representation of the shift operator $R(z)$. The left- and the right-hand side fermion lines represent atom No. 1 and No. 2, respectively. They are not linked since any interaction between the atoms is excluded. The photon lines are represented by horizontal lines arranged according to the chronology of the absorption and emission processes. The last two rows correspond to diagrams inducing correlations because they account for the mixing of interactions occurring in the two atoms.

$${}^{(k+1)}G_{b_1a_2,a_1a_2}(z)=\frac{\alpha_1}{\left[z-\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-R_{a_1a_1}(z)\right]\left[z-\omega+\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-R_{a_2a_2}(z)\right]-|\alpha_1|^2}, \tag{3.1}$$

with the two atoms initially in the upper state. This formula is called the two-level formula because it has the same form as the one encountered in the theory of two-level atoms.

In Eq. (3.1), we have changed the origin of the energies by subtracting the quantity $n\omega$ everywhere. For each atom this origin is half the distance of the (naked) levels, i.e., $\omega_{0(1,2)}=|\omega_{a(1,2)}-\omega_{b(1,2)}|$. The atom-field parameters α_1 and α_2 are those of Eq. (2.4), where the subscripts and superscripts are replaced by a single subscript which distinguishes the two systems through both their atomic and field parameters. The α's are related to the intensity by the relation

$$I_{\mathrm{W/cm^2}}=\frac{I_0}{\langle|\vec{\varepsilon}\cdot\vec{r}|\rangle^2_{\mathrm{a.u.}}}\,\alpha^2_{\mathrm{a.u.}}, \tag{3.2}$$

where $I_0=14.038\times10^{16}\ \mathrm{W\,cm^{-2}}$ [9].

In the case of $1S$-$2P$ transition in hydrogen, Eq. (3.2) reduces to $I_{\mathrm{W/cm^2}}=2.53\times10^{17}\times\alpha^2_{\mathrm{a.u.}}$. The operator $R(z)$ is called the effective operator or the shift operator because it provides the diagonal contributions to the shifts of the levels a_1 and b_1. Its matrix elements are given by

$$R_{a_1a_1}(z)=\cfrac{|\alpha_1|^2}{z+\omega+\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-\cfrac{|\alpha_1|^2}{z+2\omega-\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-\cdots}-\cfrac{|\alpha_2|^2}{z+2\omega+\frac{\omega_{01}}{2}+\frac{\omega_{02}}{2}\cdots}} \tag{3.3}$$

and

$$R_{b_1b_1}(z)=\cfrac{|\alpha_1|^2}{z-2\omega-\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-\cfrac{|\alpha_1|^2}{z-3\omega+\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-\cfrac{|\alpha_1|^2}{z-4\omega-\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-\cdots}-\cfrac{|\alpha_2|^2}{z-4\omega+\frac{\omega_{01}}{2}+\frac{\omega_{02}}{2}+\cdots}}}$$

$$+\cfrac{|\alpha_2|^2}{z+\frac{\omega_{01}}{2}+\frac{\omega_{02}}{2}-\cfrac{|\alpha_2|^2}{z+\omega+\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-\cfrac{|\alpha_1|^2}{z+2\omega-\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}-\cdots}-\cfrac{|\alpha_2|^2}{z+2\omega+\frac{\omega_{01}}{2}+\frac{\omega_{02}}{2}+\cdots}}}$$

$$+\cfrac{|\alpha_1|^2}{z-2\omega+\frac{\omega_{01}}{2}+\frac{\omega_{02}}{2}-\cfrac{|\alpha_1|^2}{z-3\omega-\frac{\omega_{01}}{2}+\frac{\omega_{02}}{2}-\cfrac{|\alpha_1|^2}{z-4\omega+\frac{\omega_{01}}{2}+\frac{\omega_{02}}{2}-\cdots}-\cfrac{|\alpha_2|^2}{z-4\omega-\frac{\omega_{01}}{2}-\frac{\omega_{02}}{2}+\cdots}}} \tag{3.4}$$

One must note that these matrix elements of $R(z)$ contain the whole physics of the process we are concerned with. Thus, the correlations come from crossed processes where one or several photon emissions or absorptions by one atom is followed by an equivalent number of absorptions or emissions by the other atom. More precisely, the atoms are correlated by the mixing of the two species of atom-field quantities in the denominators of the continued fractions.

B. Diagrammatical representation

The expansion of the operator $R(z)$ is represented in Fig. 1 by Feynman diagrams of increasing orders inspired by those encountered in the many-body theory [10]. For each diagram, the vertical line on the left stands for atom No. 1 and that on the right stands for atom No. 2. The photon absorptions and emissions are represented by horizontal lines located on the left-hand side and the right-hand side of each fermion line, respectively. In contrast to many-body diagrams, there is no coulombic line because we have assumed that the atomic wave functions do not overlap (no interaction). Therefore, one would obtain an equivalent representation by superimposing everywhere the two atomic lines and by labelling the photons arriving or leaving the single electron lines. Concerning the topology, to each order there are diagrams containing (i) the absorption and the emission operators of atom No. 1, (ii) the absorption and the emission operators of atom No. 2, and (iii) the mixing of absorption and emission operators coming from the two atoms. This last class of crossed diagrams is responsible for the correlations since they mix the events corresponding to atom No. 1 and atom No. 2. The operator expressions are obtained from the diagrams of Fig. 1 by applying the following rules: (i) the operators are written in the order they appear when the diagrams are scanned upward, (ii) a photon arriving at (leaving) the atomic line lying on the left-hand side of the diagram is represented by the operator $A(B)$, (iii) a photon arriving at (leaving) the atomic line located at the right-hand side of the diagram is be represented by the operator $X(Y)$. For example, the first diagram of the fifth array of Fig. 1 stands for $YBBAXA$, a six-order contribution to $R(z)$.

C. Numerical results

To calculate the integral of Eq. (2.12) by the technique of residues, one needs the poles of ${}^{(k+1)}G_{b_1a_2,a_1a_2}(z)$. They are determined to any desired accuracy by a method which consists of searching the position of the divergences which these poles induce. In Figs. 2(a)–2(c), we have plotted the energy as a function of ω_{01}, the energy-level separation of atom No. 1. The field frequency ω, whose value is unity in the calculation, is used as a scaling parameter. These curves correspond to atoms of equal energy-level separations ($\omega_{01}=\omega_{02}$). They have been drawn by breaking the symmetry between the two atom-field systems. This break of the symmetry is produced by changing the ratio of the two atom-field parameters. For atom No. 1 $\alpha_1=0.3$ a.u., which according to Eq. (3.2) corresponds to an intensity of 2.27×10^{16} W/cm^2, while that of atom No. 2 is chosen in such a way that α_2/α_1 varies from 0.2 to 0.8. This ratio accounts for all the possible causes that can make the two systems different from each other (intensity, polarization, phase, etc.). Five iterations in the continued fractions are enough to ensure a good stability of the solution. We observe that the poles are symmetrically distributed on both sides of the line $z=\omega$. Even with the simplified formulas of Eqs. (2.18)–(2.22) and truncated continued fractions, there exists a great number of poles in the upper and the lower parts of the energy plane. These result from the higher-order iterations. In spite of their abundance, the contributions of these poles are much smaller than those provided by the three principal poles lying in the energy range [$z=0$, $z=2\omega$]. These curves confirm what we have previously observed in Eqs. (2.11), i.e., a couple of corre-

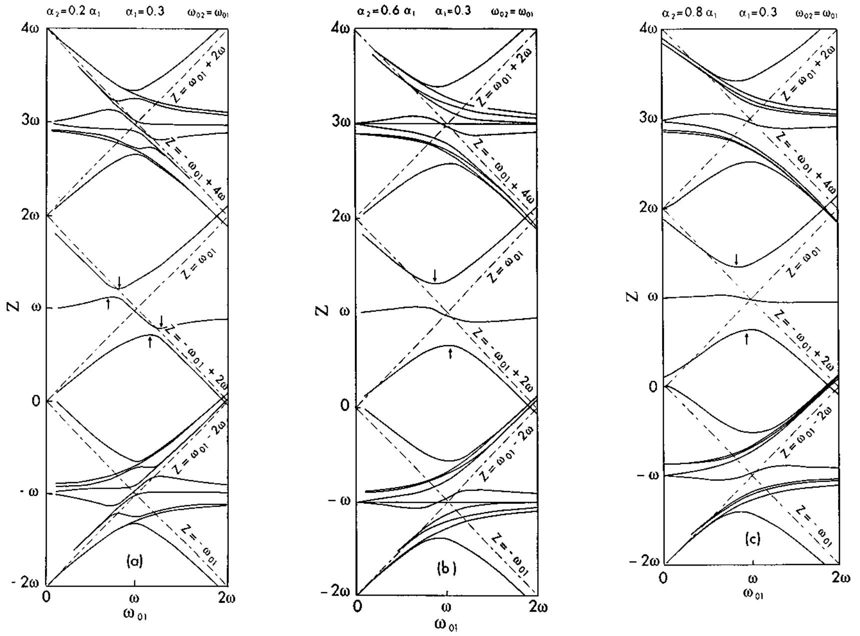

FIG. 2. Energy of a system, consisting of two correlated two-level atoms resonantly coupled with a laser field, plotted against the energy separation of the atomic levels. The dashed lines are the energy curves in the absence of level shifts. (a), (b), and (c) are obtained in the case where $\omega_{01}=\omega_{02}$ and $\alpha_1/\alpha_2=0.2$, 0.6, and 0.8 a.u., respectively. The value of α_1 is chosen to be 0.3 a.u. The curves lying within the range $0\leq z\leq 2\omega$ correspond to the principal poles whose extrema are identified by the arrows. They tend to be identical to the ones of a three-level system when the energy parameters of the atoms become equal.

lated atoms behaves like a three-level atom. They show that a modulation of α_2 generates a modulation of the perturbation which atom No. 1 suffers from the presence of atom No. 2. The energy of the two correlated atoms is represented by three curves whose asymptotes are governed by the equations $z=\omega_{01}$, $z=\omega$ (not represented on the curves), *and* $z=-\omega_{01}+2\omega$. In the fully correlated case (absence of asymmetry between the two atom-field systems), the former and the latter are symmetrically distributed with respect to the remaining one, which reduces to a straight line. It is remarkable that the distortions of the curves displayed in Figs. 2(a)–2(c) account for the disturbance which atom No. 1 suffers from the presence of atom No. 2. Thus, the more distorted the curves are, the less symmetric are the atoms. The most important departures of the curves with respect to the idealistic ones occur when the two atom-field systems are completely uncorrelated. These distortions are fully exploited in the effect we have previously called STOC.

Similar features are also observed on the probability curves of Fig. 3, which are calculated from the technique of residues. The plots are the resonant single-photon emission probability versus ω_{01} for the values of the ratio $\alpha_2/\alpha_1=0.2$, 0.4, 0.6, 0.8, 1. The value of this ratio increases when the figure is read downward. For values of α_2/α_1 less than 0.2 a.u. the resonance curves are identical to that of a single atom. The amplitudes are near unity and the maxima are shifted with respect to ω, which is the resonant frequency in the absence of Bloch-Siegert shift. When the ratio increases, the maximum is no longer unity, but decreases, with the greatest flattening occurring when $\alpha_1=\alpha_2$, i.e., when the two systems are identical. The decrease of the curve's amplitude is around 60%, a ratio which makes possible the detection of the signal variations.

In Figs. 4(a)–4(c), we have plotted the energy of a two-atom system against ω_{01} by keeping constant the atom-field parameters which have the common value $\alpha_1=\alpha_2=0.3$ a.u. The symmetry is broken by changing the ratio ω_{02}/ω_{01} which takes the values 0.25, 0.5, and 0.75 a.u. in Figs. 4(a)–4(c) respectively. Here again, we observe important distortions of the curves representing the principal poles which show the same features as before. In particular, the shape of the curves of a three-level system is practically restored if the ratio ω_{02}/ω_{01} is greater than 0.75 a.u., i.e., when the two atom-field systems tend to be identical. The slight difference is that the curves of Figs. 4(a)–4(c) are smooth compared to the curves of Figs. 3(a)–3(c). The consequences are shown in Fig. 5 where we observe important oscillations of the probability. By remembering that the probability is inversely proportional to the distance of the curves involved in an avoided crossing, we can interpret these fluctuations as being due to the competition of several avoided crossings which, in contrast to the preceding case, results in comparable contributions to the probability. Notice that the number of oscillations increases when the ratio ω_{02}/ω_{01} tends to unity. In this case the avoided crossings take place in a region of the energy plane which becomes increasingly smaller and favors

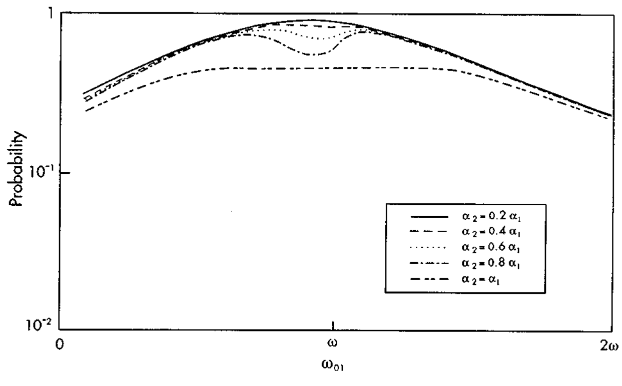

FIG. 3. Resonance curves corresponding to the net emission of a photon by atom No. 1 in the presence of atom No. 2 for different values of the ratio α_2/α_1 ($\alpha_1=0.3$) and for $\omega_{01}=\omega_{02}$. For values of the ratio less than 0.2 the curves tend to be identical to those characterizing a single atom. While the ratio increases, the maximum is no more unity, it decreases and the greatest flattening occurs when the two systems are identical. In this case, the lowering of the amplitude can be about 60%.

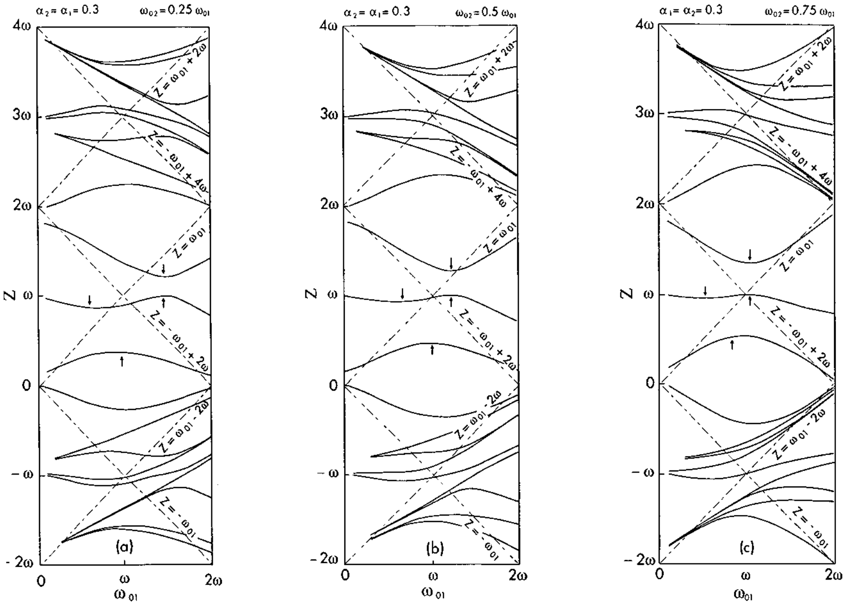

FIG. 4. Same as Figs. 2(a)–2(c) except that the atom-field parameters of the two atoms are both equal to 0.3 a.u. while the ratio ω_{02}/ω_{01} takes the values 0.25, 0.5, and 0.75 a.u. As in Figs. 2, the arrows indicate the extrema of the curves corresponding to the three principal poles.

the interference between the different contributions.

From the curves displayed in Figs. 3 and 5 it appears that the effect is more easily detectable in the case when the break of the symmetry is produced by the modulation of the ratio α_2/α_1 instead of ω_{02}/ω_{01}. Anyway, in both cases, the distortions of the resonance curves of atom No. 1 are due to the presence of atom No. 2. The behavior of these curves is in complete agreement with what it is found in the theory of dispersive forces. In both cases the presence of neighboring atoms produces a lowering of the decay probability of any excited atom and thus increases the lifetime of the excited state. This is true whether the atoms interact with each other or not and/or whether the decay proceeds via spontaneous or stimulated emission. The problem is to find the most favorable conditions in order to improve the signal-to-noise ratio and thus to increase the probability for observing the perturbation suffered by atom No. 1.

The next step of our foregoing studies will be to evaluate to what extent the amplitude of the signal prevails over the noise when the principal causes of decorrelation are taken into account. By generalizing the calculation to the more complicated system of several atoms, one can expect an enhancement of the effect. This is due to the presence of additional terms appearing in the denominators of the continued fractions, which reinforce the correlations and thus, contribute to magnify the amplitude fluctuations observed with two atoms.

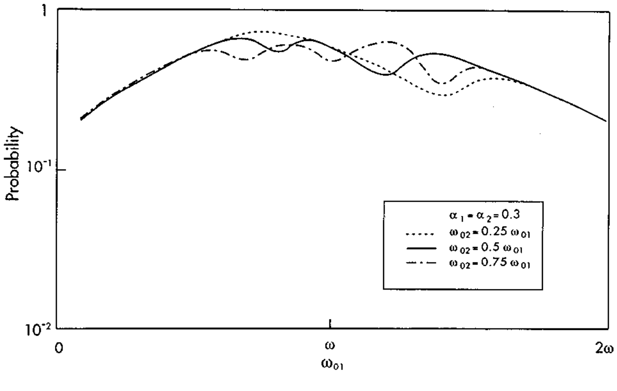

FIG. 5. Same as Fig. 3 except that $\alpha_1=\alpha_2=0.3$ a.u. and $\omega_{02}/\omega_{01}=0.25$, 0.5, and 0.75 a.u. The frequency of the oscillations observed on the resonance curves increases when the value of the ratio ω_{02}/ω_{01} tends to unity while their amplitudes decrease to give the lowest curve displayed in Fig. 3.

IV. EXPERIMENT

The STOC effect predicted along the preceding lines deserves experimental investigation since it is the only way to convince oneself of the existence of atom-atom correlations. Without anticipating future experimental devices, we can sketch the principles of a *gedankenexperiment.*

Since it is assumed that the photons interacting with the atoms are identical, the phase of the field must be the same at the position of the atoms. This condition is more easily fulfilled if the photons come from the same radiation field. For this reason, as it is shown in Fig. 6, we consider an intense laser beam which interacts with atoms contained in different cells. According to the preceding discussion, the experimental device is such that the field is the same inside the cells. To produce the ‘‘reciprocal’’ of the volumes confined in the cells, we assume that they lie around two points separated from each other either by a distance much smaller than the wavelength of the field or by an exact multiple of this wavelength.

To surmount the difficulty of observing the emission of photons at the frequency of the field, we could imagine producing a photon cascade by introducing a third level. The strength of the emission line between the two lower levels is directly related to the emission rate (involving the two upper levels) we want to measure.

As is shown in Fig. 6(a), in the absence of any perturbation on the sample *A*, the signals seen by the two observers are the same. Now, we break the symmetry between the two samples [Fig. 6(b)] by detuning the resonance between the two upper levels. For example, such a detuning can be pro-

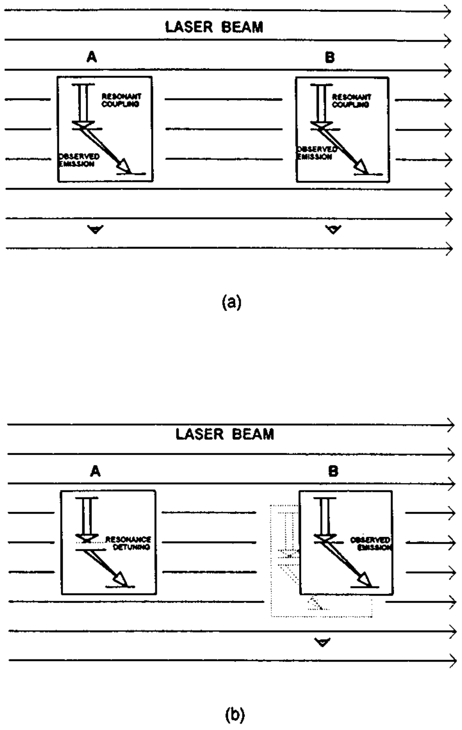

FIG. 6. Schematic of a *gedankenexperiment* for the observation of STOC. A laser beam interacts resonantly (coupling between the two upper levels) with the atoms contained in two cells located at two point where the field is the same (reciprocal areas). In the absence of any perturbation (a), the signals seen by the two observers are the same (emission line between the two lower levels). When the resonance in sample A is detuned (b), the symmetry is destroyed and the detuning can be detected in the sample B by observing the frequency shift of the emission line (dotted sketch) or by measuring the attenuation of the resonance peak.

duced by superimposing on the laser field an external field which induces a Stark shift on the atomic levels of the atoms contained in sample A. Our theory predicts that this perturbation has a non-negligible probability to be detected by an observer facing sample B. So, a detector located at point B picks up two signals, one of them comes from the unperturbed emission of the atoms contained in sample B (signal B), and the other one is due to the shifted line generated in sample A (signal A) and shared by some atoms contained in sample B by means of atom-atom correlations.

V. CONCLUDING REMARKS

The present work is the first step of forthcoming studies concerning the effect we have called STOC, whose main interest lies in its ability to transmit informations between two systems interacting resonantly and separately with a radiation field. It is produced by the coupling of two identical atoms with the same radiation field. Such an identity makes the two atom-field systems indistinguishable and thus enables the study of their behavior within a space-independent model. The quantitative analysis of Sec. IV shows that the correlations strongly depend on the states of the two systems. Mathematically, the correlations are displayed by the nonvanishing commutators involving destruction and creation operators coming from the two atom-field systems.

The idea that is underlying in this account is not completely disconnected from the reality of the observed processes. In particular, our starting point has been a microscopic interpretation of cooperative emission, which is the consequence of correlations between atoms confined to regions of space much smaller than the wavelength of the field. Since in this case the correlations require that the atoms are exposed to the same field, we have generalized this concept by assuming that such correlations hold between atoms which see the same field in regions of space separated by many wavelengths (reciprocal areas). In the present work, the problem has been reduced to the simple case of a pair of two-level atoms resonantly coupled to an intense radiation field. We have calculated the emission probability of a single photon by one of the two atoms (say atom No. 1), which were both initially in their upper state. The occurrence of the effect is not restricted to a special preparation of the system since the other atom (say atom No. 2) reveals its presence in every cases. Thus, the situation where atom No. 1 is in the upper state and atom No. 2 in the lower state leads to the same conclusions. It is the ability for the two atoms to exchange photons with the field which originates the correlations. By examining the energy and the resonance curves, we see that a break of the symmetry coming from a change of the value of any parameter (modulation) characterizing one of the two atoms can be observed on the other atom. An experiment based on the principles previously discussed is proposed. It puts our ideas in concrete form and illustrates how future experiments must proceed in order to create reciprocal areas, (i) by removing the two samples far away from each other, and (ii) by inserting a screen inside an atomic vapor confined in a cell. In all cases, the measurement will be made in regions of space where the field has the same intensity, the same phase and the same polarization. These conditions are expected to be more easily fulfilled by using a laser beam instead of two radiation fields.

Our calculation does not say anything about how the information is transmitted, but the time spent for the photons to be absorbed can be invoked in order to preserve the Einstein's causality principle. We limit ourselves to detecting on atom No. 1 the perturbation suffered by atom No. 2, without concern about the instant at which the perturbation was initiated and the time spent by the signal to reach the experimental device. In this respect, we adopt the point of view that correlated measurements on the two systems require only the presence of two observers exchanging information at a velocity less than that of the light [11]. Anyway, only experimental tests of such an effect, if observable, will provide reliable arguments concerning the way in which the signal chooses to propagate. The objections concerning this effect essentially deal with the causes which can lessen or occultate it.

ACKNOWLEDGMENTS

The author would like to thank Professor M. L. Ter-Mikaelian, Professor J. H. Eberly, and Dr. P. Agostini for stimulating discussions. He is also indebted to Dr. L. Di Mauro for reviewing the presentation of the manuscript.

[1] M. J. Stephen, J. Chem. Phys. **40**, 669 (1964); D. A. Huchinson and H. F. Hameka, *ibid*. **41**, 2006 (1964).
[2] R. H. Dicke, Phys. Rev. **93**, 99 (1954).
[3] D. Bohm, *Quantum Theory* (Prentice-Hall, New York, 1951), Chap. 23; D. Bohm and Y. Aharonov, Phys. Rev. A **108**, 1070 (1957).
[4] A. Einstein, N. Rosen, and B. Podolsky, Phys. Rev. **47**, 777 (1935).
[5] E. R. Buley and F. W. Cummings, Phys. Rev. A **134**, 1454 (1964).
[6] Y. Gontier and M. Trahin, Phys. Rev. A **7**, 1899 (1973).
[7] Y. Gontier, N. K. Rahman, and M. Trahin, Phys. Rev. A **14**, 2109 (1976).
[8] Y. Gontier, Phys. Rev. A **31**, 279 (1985).
[9] Y. Gontier and M. Trahin, Phys. Rev. A **40**, 1351 (1989).
[10] I. Lindgren and J. Morrison, *Atomic Many-body Theory*, Chemical Physics Vol. 13 (Springer Verlag, Berlin, 1982).
[11] A. K. Biswas, G. Compagno, G. M. Palma, R. Passante, and R. Persico, Phys. Rev. A **42**, 4291 (1990).