Photoelectron antibunching and absorber theory

1980 J. Phys. A: Math. Gen. 13 1389

(http://iopscience.iop.org/0305-4470/13/4/028)

J. Phys. A: Math. Gen. **13** (1980) 1389–1394. Printed in Great Britain

Photoelectron antibunching and absorber theory

D T Pegg
School of Science, Griffith University, Brisbane 4111, Australia

Received 26 June 1979, in final form 12 September 1979

Abstract. The recently detected photoelectron antibunching effect is considered to be evidence for the quantised electromagnetic field, i.e. for the existence of photons. Direct-action quantum absorber theory, on the other hand, has been developed on the basis that the quantised field is illusory, with quantisation being required only for atoms. In this paper it is shown that photoelectron antibunching is readily explicable in terms of absorber theory and in fact is directly attributable to the quantum nature of the emitting and detecting atoms alone.

The physical nature of the reduction of the wavepacket associated with the detection process is briefly discussed in terms of absorber theory.

1. Introduction

There has been much interest in the recently detected phenomenon of photon antibunching in the light emitted from single atoms acted on by an external source (Dagenais and Mandel 1978, Kimble *et al* 1977, Carmichael and Walls 1976a, b, Cohen–Tannoudji 1977, see also Knight 1977). The importance of the effect is that it is taken as direct evidence for the intrinsic quantum nature of the electromagnetic field, i.e. for the existence of photons. Although there is a connection between this effect and the dynamic Stark effect which itself has been regarded as evidence for the quantised field (Kimble and Mandel 1976), it would seem that photon antibunching is a more rigorous and convincing test, in that it appears to be more directly attributable to the commutation relations of the field operators.

In direct-action quantum absorber theory, however, there are no quantised fields or field commutation relations and all quantum effects must result from quantisation of the atoms alone (Hoyle and Narlikar 1969). It has been shown recently (Pegg 1979) that semiclassical absorber theory, without field quantisation, is sufficient to account for the dynamic Stark effect. The question thus arises: can absorber theory, without photons, stand up to the more rigorous and direct test of photon antibunching?

In order to answer this question, it must be remembered, of course, that the name 'photon antibunching' is only applicable within the paradigm of quantum electrodynamics and the actual effect observed can be more directly termed 'photoelectron antibunching'. In this paper it is shown that photoelectron antibunching is predicted by absorber theory, and in this theory it is a quantum effect arising directly from the quantisation of the emitter and detector atoms alone.

0305-4470/80/041389+06$01.50

2. Quantum absorber theory

In the classical absorber theory of Wheeler and Feynman (1945) the radiative reaction force on an accelerated charge is not due to self-action but to the advanced response of the absorber particles to the retarded action of the accelerated charge. In their direct-action formulation of classical absorber theory Wheeler and Feynman (1949) showed that the concept of the field as a separate entity is not necessary. As a consequence, in quantum absorber theory spontaneous emission and the Lamb shift do not arise from a quantised self-field or from quantum fluctuations of the electromagnetic field, as they appear to in quantum electrodynamics, but instead from the direct action with the absorber atoms, i.e. all the other atoms in the universe (Hoyle and Narlikar 1969). This led Hoyle and Narlikar to suggest that the quantised field is merely an illusion, with quantisation being required for the atoms only. To account for the dynamic Stark effect in terms of semiclassical non-relativistic quantum absorber theory, Pegg (1979) used a formalism of the theory which involved two simplifications. The first was the representation of the absorber by a large spherical shell, centred on the emitting atom, composed of just sufficient two-level atoms to absorb all radiation at all frequencies incident on it. This was called the minimal perfect absorber. The second simplification involved being able to use a Hamiltonian approach by employing a time variable representing the retarded time instead of the usual time variable. The retarded time is the usual time minus r/c where r is the distance of the particular atom from the emitting atom. The second simplification is possible because the response of an absorber (or detector) atom b to a change in the emitter atom a arrives back at a simultaneously with that change. Similarly the reaction of a to the responsive change in b is felt by b simultaneously with that responsive change. This is because a acts on b via a retarded action but b acts on a via an advanced action. The use of the retarded time variable t can thus transform the situation to one of instantaneous, in terms of t, interaction.

Because of the simplicity achieved we shall again use the Hamiltonian approach in terms of the retarded time t and the minimal perfect absorber here. For a single atom a, acted on by some external driving agency, which can be classical if desired, and surrounded by the minimal perfect absorber, the Hamiltonian is

$$\mathcal{H}_F = \mathcal{H}_a + \mathcal{H}_{sc} + \mathcal{H}_A + \mathcal{H}_{aA} \tag{1}$$

where $\mathcal{H}_a$ is the Hamiltonian of atom a, $\mathcal{H}_{sc}$ represents the external perturbation acting on a, $\mathcal{H}_A$ is the Hamiltonian of the absorber and $\mathcal{H}_{aA}$ is the direct interaction Hamiltonian between the atom a and the absorber.

3. Photoelectron antibunching

Because the usual correlation functions of quantum electrodynamics are not relevant to quantum absorber theory because they involve field operators, it is necessary to work entirely in terms of photoelectrons produced by the direct action of the emitter atom a on the detector. The experiment we consider is as follows.

In addition to the system comprising the two-level atom a surrounded by its minimal perfect absorber and acted on by the external agency, we also consider two detector atoms b and c. These are two-level atoms in their ground states with the same energy gap ω (in units $\hbar = 1$) as atom a. When b or c makes an upward transition we assume a

photoelectron is produced. At time $t = 0$ atom c is placed inside the minimal perfect absorber, i.e. it is directly exposed to a. At time $t = \Delta$ atom c is removed, i.e. after this time it is shielded from a. Similarly at a later time $t = T$, atom b is exposed to a until time $t = T + \Delta$. We wish to find the probability that both c and b produce photoelectrons.

The complete Hamiltonian for the system is now

$$\mathcal{H} = \mathcal{H}_a + \mathcal{H}_{sc} + \mathcal{H}_A + \mathcal{H}_{aA} + \mathcal{H}_b + \mathcal{H}_c + \mathcal{H}_{ab} + \mathcal{H}_{ac} \tag{2}$$

where $\mathcal{H}_b$ and $\mathcal{H}_c$ are the Hamiltonians of atoms b and c, and $\mathcal{H}_{ab}$ and $\mathcal{H}_{ac}$ are their direct interaction Hamiltonians with a. $\mathcal{H}_{ac}$ is a function of time in that it is on at $t = 0$ and off, i.e. zero, at $t = \Delta$. Similarly $\mathcal{H}_{ab} = 0$ except between T and $T + \Delta$. $\mathcal{H}_{aA}$ takes its full value between $t = \Delta$ and $t = T$ but is diminished slightly by the shielding effect of c or b between $t = 0$ and $t = \Delta$ and between T and $T + \Delta$. $\mathcal{H}_{sc}$ causes transitions in a in a characteristic time Ω^{-1}. For example, Ω would be the Rabi frequency in the case of interaction with an oscillating electric field. $\mathcal{H}_{aA}$ causes spontaneous decay at a rate Γ. We choose the interval Δ such that $\Delta \gg \omega^{-1}$ but $\Delta \ll \Omega^{-1}, \Gamma^{-1}$. The problem is to find the probability for both atom c to make an upward transition $|0_c\rangle \rightarrow |1_c\rangle$ and for b to make the upward transition $|0_b\rangle \rightarrow |1_b\rangle$ in the time between $t = 0$ and $t = T + \Delta$.

The unitary time displacement operator for the complete system can be factorised as

$$U(T + \Delta, 0) = U(T + \Delta, T)U(T, \Delta)U(\Delta, 0). \tag{3}$$

The factor $U(T, \Delta)$ just corresponds to the Hamiltonian

$$\mathcal{H}_a + \mathcal{H}_{sc} + \mathcal{H}_A + \mathcal{H}_{aA} + \mathcal{H}_b + \mathcal{H}_c$$

because $\mathcal{H}_{ab} = \mathcal{H}_{ac} = 0$ during this period. We can thus write:

$$U(T, \Delta) = U_F(T, \Delta) \exp(-i(\mathcal{H}_b + \mathcal{H}_c)(T - \Delta)]$$

where $U_F(T, \Delta)$ is the time displacement operator for the Hamiltonian $\mathcal{H}_F$ in (1), because $\mathcal{H}_b$ and $\mathcal{H}_c$ commute with $\mathcal{H}_F$.

During the interval from $t = 0$ to $t = \Delta$ we have $\mathcal{H}_{ab} = 0$. Also, because $\Delta \ll \Omega^{-1}, \Gamma^{-1}$ the interaction terms $\mathcal{H}_{sc}$ and $\mathcal{H}_{aA}$ will have only a small effect during this period. Thus, to a first approximation we can write

$$U(\Delta, 0) \approx U_c(\Delta, 0)$$

where $U_c(\Delta, 0)$ corresponds to the Hamiltonian $\mathcal{H}_o + \mathcal{H}_{ac}$ where

$$\mathcal{H}_o = \mathcal{H}_a + \mathcal{H}_A + \mathcal{H}_b + \mathcal{H}_c$$

and

$$\mathcal{H}_{ac} = -\frac{\mu_0 e^2 \boldsymbol{v}_a \cdot \boldsymbol{v}_c}{4\pi r}$$

(Pegg 1979) where $\boldsymbol{v}$ is the (Schrödinger) velocity operator. Perturbation theory yields

$$U_c(\Delta, 0) = U^0(\Delta, 0) - ik \int_0^\Delta U^0(\Delta, \tau_2) v_a v_c U^0(\tau_2, 0)\, d\tau_2 + \text{second-order terms in } v_c + \ldots$$

where the constant k includes the factor r^{-1} and the cosine of the angle between $\boldsymbol{v}_a$ and $\boldsymbol{v}_c$, and

$$U^0(t_1, t_2) = \exp[-i\mathcal{H}_0(t_1 - t_2)].$$

A similar expression is obtainable for $U_b(T + \Delta, T)$ in terms of v_b.

Substitution of these results into (3) gives an expression which contains one zero-order term in v_b and v_c, two first-order terms, three second-order terms plus higher-order terms. Because of the off-diagonal nature of v_b and v_c the transition amplitude

$$\langle A_i|\langle a_j|\langle 1_c|\langle 1_b|U(T+\Delta, 0)|0_b\rangle|0_c\rangle|a(0)\rangle|A_0\rangle$$

will, as far as second order, have only one non-zero term, that which contains both v_b and v_c. Here $|a_i\rangle$ and $|A_i\rangle$ are eigenstates of a and the absorber, $|a(0)\rangle$ is the state of a at $t=0$ and $|A_0\rangle$ is the initial state of the absorber, in which all absorber atoms are in their ground states. This non-zero term is proportional to

$$\langle A_i|\langle a_j|\langle 1_c|\langle 1_b|\int_T^{T+\Delta} U^0(T+\Delta, \tau_1)v_a v_b U^0(\tau_1, T)\,d\tau_1 U_F(T, \Delta)\exp[-i(\mathscr{H}_b+\mathscr{H}_c)$$

$$\times(T-\Delta)]\int_0^{\Delta} U^0(\Delta, \tau_2)v_a v_c U^0(\tau_2, 0)\,d\tau_2|0_b\rangle|0_c\rangle|a(0)\rangle|A_0\rangle. \tag{4}$$

For now we write $|a(0)\rangle = m|e\rangle + n|g\rangle$ where $|e\rangle$ and $|g\rangle$ are the excited and ground states of a, and shall discuss the state $|a(0)\rangle$ in more detail later. Working through from the left, we find for the latter part of expression (4):

$$\ldots U_F(T, \Delta)|0_b\rangle|1_c\rangle|A_0\rangle \exp[-i\omega(T-\Delta)]v_{10}[m|g\rangle \exp(-i\omega\Delta)v_{ge}\Delta + n|e\rangle v_{eg}$$

$$\times\int_0^{\Delta}\exp(-2i\omega\Delta+2i\omega\tau_2)\,d\tau_2]$$

where $v_{ge}=\langle g|v_a|e\rangle$ and $v_{10}=\langle 1|v_c|0\rangle$. The Δ in the term involving $m|g\rangle$ comes from a factor

$$\int_0^{\Delta}\exp\{i(\omega_c-\omega)\tau_2\}\,d\tau_2 \tag{5}$$

where ω_c is the energy gap of atom c, which we have chosen to equal ω. Upon integration, with $\Delta \gg \omega^{-1}$ the term involving $n|e\rangle$ goes to zero, ignoring a small imaginary part, leaving expression (4) as

$$m\exp[-i\omega(T+\Delta)]v_{10}v_{ge}\Delta\langle A_i|\langle a_j|\langle 1_b|\int_T^{T+\Delta} U^0(T+\Delta, \tau_1)v_a v_b U^0(\tau_1, T)\,d\tau_1$$

$$\times|0_b\rangle(|g\rangle\langle g|+|e\rangle\langle e|)U_F(T, \Delta)|g\rangle|A_0\rangle \tag{6}$$

where we have inserted the closure expression for unity. Again, just as before, of the two terms in (6) only that involving $|e\rangle\langle e|$ is non-zero upon integration, and we obtain, neglecting the phase factors which will not affect the probability and omitting the constant factors v_{10}, v_{ge}, Δ,

$$m\langle a_j|g\rangle\langle A_i|\langle e|U_F(T, \Delta)|g\rangle|A_0\rangle.$$

The probability that c and b make upward transitions is found by multiplying this amplitude by its complex conjugate and summing over all final states of the atom and the absorber, i.e. over i and j. In the summation over j only the term with $\langle a_j| = \langle g|$ is non-zero and the probability is thus proportional to

$$|m|^2\sum_i|\langle A_i|\langle e|U_F(T, \Delta)|g\rangle|A_0\rangle|^2. \tag{7}$$

The factor $|m|^2$ is the quantum-mechanical probability for the atom a to be in its excited state at $t = 0$. Expression (7) has been obtained by writing $|a(0)\rangle = m|e\rangle + n|g\rangle$. In general, however, $|a(0)\rangle$ would not be known so precisely and is more appropriately described by a statistical mixture which involves a range of values of $|m|^2$ (see, for example, Glauber 1965). For such an initial-state expression (7) will be replaced by its average over all possible values of $|m|^2$. Thus, in place of $|m|^2$ we have the ensemble average $\overline{|m|^2}$ which represents the combined statistical and quantum-mechanical probability for the atom to be in its excited state at $t = 0$.

The second factor in (7) is just the probability for an upward transition of a from $|g\rangle$ to $|e\rangle$ in the interval between the detection times, under the combined action of the external agency and spontaneous emission (Pegg 1979). Thus, as the time interval $T - \Delta$ approaches zero the probability of two-photoelectron detection approaches zero also, producing the antibunching effect.

4. Detectors

The two-level atoms b and c were chosen as detectors for simplicity in directly illustrating that quantum absorber theory does give the photoelectron antibunching effect. To derive the dynamic Stark spectrum from absorber theory such a two-level atom was used as a detector finely tuned to its Bohr frequency (Pegg 1979). Because in the present paper two-level atoms are also being used to detect what is effectively a frequency-integrated signal from the atom a, it is worth discussing briefly the means by which the same type of atom can accomplish these two very different functions. Basically this is due to the different lengths of observation time. In both cases the probability for an upward transition in a detector atom is, from (5), proportional to the expression

$$\int_0^{\Delta} \exp[\mathrm{i}(\omega_c - \omega)\tau]\,\mathrm{d}\tau \times \mathrm{CC} \tag{8}$$

where ω_c and ω are the energy gaps of the detector atom and the atom a. This expression significantly differs from zero only for $|\omega_c - \omega| < \Delta^{-1}$. For the dynamic Stark effect the observation time Δ is very long, which makes the detector atom very frequency-selective. For the antibunching effect we chose, for convenience, an atom with $\omega_c = \omega$; however, the result obtained from (4) is not appreciably altered if we let ω_c differ from ω in (5) by amounts up to Δ^{-1}, which is quite large for the small observation time in this experiment, i.e. $\Delta \ll \Omega^{-1}, \Gamma^{-1}$. Thus the occurrence of an upward transition in the detector specifies the frequency of the signal from a only to within this order of uncertainty. This, of course, is just in accord with Heisenberg's uncertainty principle. In effect, the short observation time destroys the selectivity, i.e. broadens the frequency response of the detecting atom to give a bandwidth of order Δ^{-1}.

Another point, which can be noticed from (6), is that the upward transition probability in one particular detector atom in time Δ is proportional to Δ^2, and consequently although the two-level detector does show the antibunching effect it is not an *ideal* detector in the sense described by Glauber (1965), i.e. it does not give rise to a Δ-independent counting rate. We could obtain such an ideal detector by replacing the single excited state $|1_c\rangle$ in the detector atom c by a large number of evenly and closely spaced states $|k_c\rangle$ which have $\langle k_c|v|0\rangle$ reasonably constant, and similarly replace atom b. This in fact would correspond more closely to detectors actually used in practice. The

result for such a system is obtainable by replacing $\langle 1_c|$ and $\langle 1_b|$ in (4) by $\langle k_c|$ and $\langle z_b|$ and summing, or integrating, the final probability obtained over k and z. Expression (8) with ω_c replaced by ω_k would no longer equal Δ^2, and it is not difficult to show that upon integration over ω_k we obtain a factor proportional to Δ instead, provided the range of values of $(\omega - \omega_k)$ is centred around zero and has a breadth greater than Δ^{-1}. The summation over k and z which accompanies the use of such more realistic detecting atoms also naturally gives a larger signal than that produced by two-level detectors, but importantly leaves unaffected the factor (7) in the final result for the probability, which is the characteristic of the antibunching effect. Thus absorber theory predicts photoelectron antibunching from both two-level and ideal detecting atoms.

5. Conclusions

The simplicity with which direct action quantum absorber theory displays the photoelectron antibunching effect allows us to interpret this effect as a quantum effect, *not* of the electromagnetic field but of the source and detector atoms alone. In fact, the antibunching effect, like the dynamic Stark effect, can not be taken as evidence even for the existence of the field. This lends further support to the suggestion of Hoyle and Narlikar (1969) concerning the illusory nature of the quantised field.

An interesting aspect of the absorber theory approach is the light it sheds on the nature of the reduction of the wavepacket pointed out by Cohen–Tannoudji (1977). In absorber theory an upward transition in the detector atom c, which constitutes the measurement process, can be interpreted as the actual physical *cause* of the corresponding reversion of atom a to its ground state despite the fact that the transition in c occurs *after* a has reverted to its ground state. This is because a physically stimulates c by retarded action, eliciting a response from c which acts physically on a by means of advanced potential action. This self-consistent cycle by which a transition in a causes a transition in c which causes the original transition in a has been discussed by Hoyle and Narlikar (1968).

References

Carmichael H J and Walls D F 1976a *J. Phys. B: Atom. Molec. Phys.* **9** L43
—— 1976b *J. Phys. B: Atom. Molec. Phys.* **9** 1199–219
Cohen-Tannoudji C 1977 *Proc. 2nd Laser Conf., Megeve* (Berlin: Springer-Verlag)
Dagenais M and Mandel L 1978 *Phys. Rev.* A **18** 2217–28
Glauber R J 1965 *Quantum Optics and Electronics, Les Houches* ed C deWitt, A Blandin and C Cohen-Tannoudji (New York: Gordon and Breach) pp 64–108
Hoyle F and Narlikar J V 1968 *Nature* **219** 340–1
—— 1969 *Ann. Phys., NY* **54** 207–39
Kimble H J, Dagenais M and Mandel L 1977 *Phys. Rev. Lett.* **39** 691
Kimble H J and Mandel L 1976 *Phys. Rev* A **13** 2123–44
Knight P 1977 *Nature* **269** 647
Pegg D T 1979 *Ann. Phys., NY* **118** 1–17
Wheeler J A and Feynman R P 1945 *Rev. Mod. Phys.* **17** 157–81
—— 1949 *Rev. Mod. Phys.* **21** 425–33