The myth of the photon

Trevor W. Marshall, Dept. of Mathematics, Univ. of Manchester, Manchester M13 9PL, UK
Emilio Santos, Depto de Física Moderna, Univ. de Cantabria, 39005 Santander, Spain

Abstract
We have shown that all “single-photon” and “photon-pair” states, produced in atomic transitions, and in parametric down conversion by nonlinear optical crystals, may be represented by positive Wigner densities of the relevant sets of mode amplitudes. The light fields of all such states are represented as a real probability ensemble (not a pseudoensemble) of solutions of the unquantized Maxwell equation.

The local realist analysis of light-detection events in spatially separated detectors requires a theory of detection which goes beyond the currently fashionable single-mode “photon” theory. It also requires us to recognize that there is a payoff between detector efficiency and signal-noise discrimination. Using such a theory, we have demonstrated that all experimental data, both in atomic cascades and in parametric down conversions, have a consistent local realist explanation based on the unquantized Maxwell field.

Finally we discuss current attempts to demonstrate Schroedinger-cat-like behaviour of microwave cavities interacting with Rydberg atoms. Here also we demonstrate that there is no experimental evidence which cannot be described by the unquantized Maxwell field.

We conclude that the “photon” is an obsolete concept, and that its misuse has resulted in a mistaken recognition of “nonlocal” phenomena.

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1 Introduction

Since the beginning of modern physics, that is since the seventeenth century, there have been two views as to the nature of light. The corpuscular view has, traditionally, been supported by those most strongly attracted to formal elegance, whilst the undulatory view has been supported by those who insist on the necessity for science to give explanations of phenomena (1). It is no coincidence that one of the earliest and strongest statements against trying to explain field phenomena is in Newton’s Opticks, and that a similar statement was made, nine years later, in the Preface to his Principia.

Formal elegance, combined with Newton’s authority, dominated the eighteenth and the first part of the nineteenth century. Light corpuscles, going
near to a sharp edge, experienced, according to their ideas, instantaneous (in today’s terminology nonlocal) interactions with the edge, and that caused a phenomenon they called inflexion. Today that phenomenon is called diffraction, and the name has changed because Young, Fresnel, Faraday and Maxwell taught us that nonlocal “explanations” are not explanations at all. What they gave us instead was a consistently wave explanation of diffraction and interference, and theirs remains the only explanation of those phenomena right up to the present day.

So how does it come to pass that the strongest claims to have observed “quantum nonlocality” now come from certain opticians? We suggest that it is because certain opticians have allowed themselves to be seduced by formal elegance, just like Newton’s immediate successors. Indeed, just like their intellectual ancestors, they have been carried away by a formally elegant mechanical theory. Yes, quantum mechanics is elegant, but only as long as it applies to systems with a few degrees of freedom. Light fields have infinite degrees of freedom, and a mature treatment of them requires the considerably less elegant apparatus of quantum field theory - not only less elegant, but bristling with all sorts of problems associated with divergences and renormalizations.

Will it be necessary to abandon the quantum formalism in order to obtain a local description of “multiphoton” processes? We cannot yet give a complete answer to this question, but we do assert that, with a certain natural extension of the term “classical”, all of the light fields, including those currently classified as “nonclassical”, which have so far been produced in the laboratory are, in fact, entirely classical; they are adequately described by the unquantized Maxwell equations. We shall see that, in order to extend our notion of “classical”, it will be necessary first to escape from Hilbert space and place ourselves in classical phase space; this enables us to adopt the point of view, which has been anathematized by the Copenhagen school, that electromagnetic waves are real waves, in ordinary space and time, having both amplitude and phase. Whereas it may seem natural, as long as we are imprisoned in Hilbert space, to think of “photons”, created at one point and absorbed at another, the phase-space description we advocate keeps us entirely within the confines of classical electromagnetic theory. We believe that this step has already been taken, but not fully acknowledged, by a substantial part of the quantum-optics community. For example, three review articles (2–4) on light squeezing make extensive use of phase-space diagrams, and one of them (3) states explicitly that the photon description of the light field is not helpful in the understanding of the phenomenon. We now propose to extend this judgement to the light emitted in atomic-cascade and parametric-down-conversion processes, as well as the microwave radiation contained in cavities.
2 What is a classical light field?

At the moment the accepted convention is to define as “classical”, or more precisely “Glauber-classical”, a field which is a mixture of pure coherent states. For a single mode of the field, the density matrix of such a state is

\[ \hat{\rho} = \int |\alpha\rangle P(\alpha) \langle \alpha | d^2\alpha, \]  

(1)

where \( P \) is nonnegative and

\[ |\alpha\rangle = \exp(\alpha \hat{a}^\dagger - \alpha^* \hat{a}) |0\rangle. \]  

(2)

As we have argued elsewhere (6−9), there is a strong case for extending the notion of “classical” to the set of states whose Wigner density is nonnegative. All Glauber-classical states are classical in this wider sense. Indeed the Wigner density of the above single-mode state is

\[ W(\alpha) = 2\pi^{-1} \int P(\alpha') \exp(-2 |\alpha - \alpha'|^2) d^2\alpha'. \]  

(3)

Any quantum state, defined by a density matrix, has a well defined Wigner density, but in general there is no guarantee that this density will be a nonnegative function. The vacuum state is Glauber-classical with

\[ P(\alpha) = \delta^2(\alpha), \]  

(4)

and therefore with

\[ W(\alpha) = 2\pi^{-1} \exp(-2 |\alpha|^2). \]  

(5)

According to some experts (10,11), the difference between the \( P \)- and \( W \)-representations is entirely formal, but we disagree (9). Equation (4) suggests that the vacuum is truly empty, whereas equation (5) suggests that there is a real nontrivial distribution of mode amplitudes and phases in the vacuum. Indeed such a view suggests that the word “vacuum” is (like “inflexion”) obsolete, and that we should call it something else (like “plenum”). This point of view is implicitly supported in that part of the quantum-optics community which takes phase-space diagrams seriously (2−4). More explicitly, the concept of the real zeropoint field has been central to stochastic electrodynamics since the early 1960s (12−14), and its role has been acknowledged more recently in quantum electrodynamics (15).

We consider Max Planck (16) to be the originator, in 1911, of the real zeropoint field, so we shall call light fields with nonnegative Wigner densities “Planck-classical”, and will consider a field to be nonclassical only if it is not Planck-classical. There are some Planck-classical fields which are not Glauber-classical. An outstanding example of such a field is the squeezed vacuum state

\[ |\zeta\rangle = \exp(\zeta \hat{a}^\dagger^2 - \zeta^* \hat{a}^2) |0\rangle. \]  

(6)
So, using our new definition, are there \textit{any} nonclassical states of the light field? The simple answer is, of course, Yes. Indeed most states of the Hilbert space are nonclassical, because, from a theorem of Hudson\cite{Hudson}, generalized by Soto and Claverie\cite{Soto}, no pure states other than the Gaussian subset have nonnegative Wigner density. In particular, the single-mode one-photon Fock state

\[ |1\rangle = \hat{a}^+ |0\rangle \]  

has Wigner density

\[ W_1(\alpha) = 2\pi^{-1} (4 | \alpha |^2 - 1) \exp(-2 | \alpha |^2). \]

With respect to the “nonclassical” states of the light field currently widely reported as having been observed, our response is that something approximating the squeezed vacuum, as described by equation (6), \emph{has} been observed; this, however, according to our new classification, is a \textit{classical} state, though not Glauber-classical. As for Fock states, represented, for example, by equation (7), we consider that the claims to have observed them are incorrect, and that discussions on such exotic properties (quantum nonlocality, entanglement etc.), which such states would have, if they were to exist, are misguided.

### 3 Is the “one-photon” state classical?

We have just seen that the single-mode one-photon state, represented by equation (7), is not Planck-classical. But we find it amazing that anyone\cite{remark} should try to discuss such questions as locality and causality on the basis of waves which fill the whole of space and time! The single-mode representation of a real atomic signal is clearly inadequate.

If we wish to represent the output of a \textit{real} atomic source, we must take account not only of the fact that each atomic light signal occupies a finite time interval (typically about 5ns), but also that neither the time nor the direction of the emitted radiation can be controlled. (We are advocating a return to an unambiguously wave description of light, so any signal is emitted into a range of directions. Nevertheless, the spatial distribution of the signal will be influenced, for example, by the atom’s charge distribution at the time of emission; this cannot be controlled.)

The first of these requirements leads us to a multimode description of the light signal, while the second forces us to abandon its description as a pure state (that is a vector in Hilbert space) and use, instead, a density matrix. We have shown\cite{density_matrix} that, after incorporating these two features, the density matrix is that of a chaotic state, that is its Wigner density is Gaussian and the state is classical. If, however, the signal is part of an atomic-cascade process, it is possible to use one signal in the cascade to monitor the observation time of its partner, as in the experiment of Grangier, Roger and Aspect\cite{Grangier}. In that case\cite{density_matrix} we must use
a different Wigner density - again positive, however - for the subensemble of monitored signals, and this, as we shall show in a later section, allows a purely wave explanation of what Grangier, Roger and Aspect thought was corpuscular behaviour.

4 The role of the zeropoint field

The real zeropoint field plays a crucial role in explaining how purely wave phenomena may be misinterpreted as evidence of corpuscular behaviour. Recognition of its role would be a convincing vindication of Max Planck\(^{(16)}\), because he introduced the concept of the zeropoint field precisely in order to oppose Einstein’s Lichtquanten, which were the forerunners of photons.

We consider first the way in which a theory with a real zeropoint field views the action of a beam splitter. Such a device was used by Grangier, Roger and Aspect\(^{(21)}\) to demonstrate a phenomenon they called anticorrelation in the outgoing channels. If we consider the “vacuum” to be empty, then it seems almost unavoidable to assume that the intensity of any incoming classical signal is equal to the sum of the intensities in the outgoing channels, and also that the detection probability in both channels is proportional to the signal intensities in those channels. With such a description it is not possible to explain the anticorrelation data; these were interpreted by Grangier, Roger and Aspect as evidence that the whole “photon” goes into one or other of the outgoing channels. It is easy to see, qualitatively, how the explanatory power of a purely wave theory is increased by the recognition of a real zeropoint field. A beam splitter does not simply split an incoming wave into two parts; it mixes together two incoming waves, one of them from the “vacuum”, to give the two outgoing waves (see Figure 1).

![Figure 1: The beam splitter mixes the incoming signal (a) with the relevant modes of the zeropoint field (d) to give the signals (b) and (c) in the two outgoing channels.](image)

Something similar occurs in a nonlinear optical crystal. An intense coherent input causes two modes of the zeropoint field, initially uncorrelated, to become
both enhanced in their amplitudes and correlated (see Figure 2). This in turn causes correlated “photon” counts in the outgoing channels. The current name for what occurs in the crystal is \textit{parametric down conversion} but this is yet another example (like “inflexion”) of a bad concept - the “photon” - giving rise to a misleading name and description; it describes an incoming photon of the coherent beam as converting into two completely new photons. But all modes of the field are \textit{already present} before the intervention of the coherent beam and the nonlinear crystal. The \textit{down conversion} is, more accurately, a \textit{correlated amplification} of certain modes of the zeropoint field.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig2.png}
\caption{The coherent beam modifies the two initially uncorrelated zeropoint amplitudes $E_0$ and $E'_0$ to produce the two correlated (“signal” and “idler”) outputs.}
\end{figure}

The two outgoing signals from the beam splitter in Figure 1 are given by\textsuperscript{(22)}

\begin{align}
E_b &= TE_a + iRE_d, \\
E_c &= TE_d + iRE_a,
\end{align}

where $R$ and $T$ are real coefficients satisfying

\begin{equation}
R^2 + T^2 = 1.
\end{equation}

The two outgoing signals from the nonlinear crystal in Figure 2 are given by\textsuperscript{(22)}

\begin{align}
E_{\text{signal}} &= E_0 + gVE'_0^*, \\
E_{\text{idler}} &= E'_0 + gVE_0^*,
\end{align}

where $g$ is a coupling constant and $V$ is the analytic signal of the coherent beam. Because these outputs are linearly related to the inputs, we have been
able to deduce that the joint Wigner density of the outgoing beams is positive, both when the inputs are all Gaussian, as in Figure 2, and when one of the inputs is “single-photon”, as would be the case in the Grangier-Roger-Aspect experiment. We remark that the linearity property holds always in the beam-splitter case, but that it holds only to first-order perturbation approximation in the nonlinear-crystal case. In the full quantum formalism, higher-order effects could, possibly, give an outgoing Wigner density taking negative values, but such effects are, at present, not experimentally observable.

5 The theory of detection

We have just seen that taking account of the zeropoint field leads us to a different understanding of certain optical devices. In particular, the recognition of the previously “missing” inputs, as in Figure 1 and Figure 2, means that the sum of the intensities of the outgoing signals is not equal to the intensity of just one incoming signal. This new feature of a zeropoint field theory is sufficient to take away the mystery of enhancement at a beam splitter. Applying this idea to a polarizing beam splitter, we have been able to show that all the “nonlocal” data for polarizations of light signals from atomic cascades have a local explanation. However, it is now necessary to modify the theory of detection. All previous semiclassical theories have omitted the zeropoint field, and it has been assumed that the detection probability is proportional to the signal intensity. Since the total intensity in all the zeropoint modes is enormously greater than any signals, it must follow that all detectors are “blind”, or nearly so, to the zeropoint intensity. This must be so even when the “signal” is the light from the Sun and the detectors are our own eyes!

The subtraction of the zeropoint noise is, we claim, already a feature of the standard theory, in which light detectors are considered to be normal-ordering devices. The probability of joint detection in the two outgoing channels of Figure 1 is given by

\[
\text{Pr}[\text{joint detection}] = \eta_b \eta_c \int_0^T dt \int_0^T dt' \langle N[\hat{I}_b(t)\hat{I}_c(t')] \rangle,
\]

where \( N \) denotes that the field amplitudes in \( \hat{I}_b \) and \( \hat{I}_c \) are normally ordered, and \( \eta_b, \eta_c \) are the detector efficiencies. We have shown that an equivalent expression is

\[
\text{Pr}[\text{joint detection}] = \eta_b \eta_c \int_0^T dt \int_0^T dt' \langle S[\{\hat{I}_b(t) - I_0\} \{\hat{I}_c(t') - I_0\}] \rangle,
\]

Footnote 1: More accurately, the linearity property holds in the approximation where we may neglect changes in the laser intensity - so-called depletion - resulting from absorption in the crystal. Within this approximation, the linearity property holds to all orders of perturbation theory. This footnote was added after submittal of the article to the editors of the Vigier Conference Proceedings.
where S denotes a symmetric ordering of the field amplitudes, and $I_0$ is the zeropoint intensity in the relevant modes. This enabled us to replace the whole expression by an integral\(^5\), over the classical phase space, involving the (positive) Wigner density, and hence give a purely wave explanation of the anticorrelation of the photoelectron counts in these channels. We have made a similar analysis\(^22\) of the correlated signals in Figure 2, and hence have been able to give a purely wave explanation of the experiments (some of which the authors describe as “mind boggling”) described in Reference\(^19\).

It remains a problem to explain how this formal subtraction of the zeropoint is actually achieved by the detectors. There must exist some positive functional of the field amplitudes whose average value, weighted by the Wigner density, is very small when only the zeropoint is present. Note that real detectors all have a finite dark rate, so the zeropoint will always give some detection events; hence the theory of detection we are demanding will actually explain more than the current theory.

In the absence of such a theory we constructed a simpler model theory\(^23,24\) in order to illustrate how the noise subtraction, in combination with the enhancement mechanism described in the previous section, gave rise to certain “particle-like” counting statistics in the two channels.

The explanation of how detectors are able to extract signals from the very large zeropoint background is a very difficult problem which we have not yet managed to solve. The day that theoretical physicists begin seriously to confront it is when they will begin, perhaps, to recover the respect of the rest of the scientific community. Despite our failure to resolve it, we state our conclusions, namely that the photon is obsolete, that light is nothing but waves, and that all wave fields fluctuate (see the opening sentence of Reference\(^4\)). The next section is nothing but a postscript to this conclusion.

6 The microwave field in a cavity

Proposals have been made for the construction of experimental situations resembling the Schroedinger cat\(^27,28\), in which two quite large objects, namely a Rydberg atom and a microwave cavity, are put in a superposition state, so that certain of their properties are “entangled”.

Since it is not, at present, possible to observe directly the state of the cavity, this entanglement (if it really existed!) could not be demonstrated so readily as would be the case, with “perfect” detectors, for entangled light signals. Hence, any serious attempt to construct a Schroedinger cat must either seek to entangle two successive Rydberg atoms passing through the same cavity\(^29\) or make some plausible additional assumption about the single Rydberg atom. We have taken the latter course\(^30\), since we are sure that it leads to experimental requirements which are easily achievable with current technology. The additional hypothesis we propose is that the two-state stochastic process represented by the Rydberg
atom be stationary. With this condition the Rydberg atom may be treated as a macroscopic system - it is bigger than a protein molecule - and the inequalities for such systems, deduced by Leggett and Garg\(^{(31)}\) (also called temporal Bell inequalities) should apply. With the very high Q-values claimed by experimenters for the cavity, it should be possible, according to current theory, to demonstrate a violation of the Leggett-Garg inequality, but our analysis of the data\(^{(32)}\) so far available shows no such violation. In the light of our experience with atomic cascades, one should be modest about the conclusions one draws. Hypotheses which seem plausible before doing an experiment should, properly, often be rejected in the light of the new evidence. This was our experience with Clauser and Horne’s\(^{(33)}\) hypothesis of no enhancement at a beam splitter. For the moment our inclination is to persist with the stationarity hypothesis for the Rydberg atom. It seems to us highly probable that the Q-values currently claimed for supercooled cavities may not take fully into account all the possible relaxation mechanisms for the radiation in the cavity, and it would not take much relaxation to convert the “ideal” quantum electrodynamic autocorrelation of the process into one satisfying the Leggett-Garg inequalities.

We wish Jean-Pierre Vigier a very happy birthday.

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References
1. T. W. Marshall, Found. Phys. 22, 363 (1992).
2. R. E. Slusher and B. Yurke, Sci. Amer. 258, 32 (1988).
3. E. Giacobino, C. Fabre, A. Heidmann and S. Reynaud, La Recherche 21, 170 (1990).
4. R. Loudon and P. L. Knight, J. Mod. Opt. 34, 709 (1987).
5. J. Peřina, Quantum Statistics of Linear and Nonlinear Optical Phenomena (Reidel, Dordrecht, 1984).
6. T. W. Marshall and E. Santos, Found. Phys. Lett. 5, 573 (1992).
7. T. W. Marshall, E. Santos and A. Vidiella-Barranco, Proc. 3rd Int. Workshop on Squeezed States and Uncertainty Relations (D. Han, Y. S. Kim, N. H. Rubin, Y. Shih, W. W. Zachary eds.) (NASA Conf. Series, No. 3270, 1994) page 581.
8. T. W. Marshall and E. Santos, Phys. Rev. A 41, 1582 (1990).
9. T. W. Marshall, Phys. Rev. A 44, 7854 (1991).
10. P. Kinsler and P. D. Drummond, Phys. Rev. A 44, 7848 (1991).
11. P. Milonni, The Quantum Vacuum (Academic, San Diego, 1993) page 142.
12. T. H. Boyer, in Foundations of Radiation Theory and Quantum Electrody-
13. L. de la Peña, in *Stochastic Processes Applied to Physics and Other Related Fields* (B. Gomez, S. M. Moore, A. M. Rodriguez-Vargas and A. Rueda, eds.) (World Scientific, Singapore, 1983).

14. T. H. Boyer, Sci. Amer. August 1985.

15. J. Dalibard, J. Dupont-Roc and C. Cohen-Tannoudji, J. Phys. (Paris) **43**, 1617 (1982).

16. M. Planck, *Theory of Heat Radiation* (Dover, New York, 1959).

17. R. L. Hudson, Rep. Math. Phys. **6**, 249 (1974).

18. F. Soto and P. Claverie, J. Math. Phys. **24**, 97 (1983).

19. D. M. Greenberger, M. A. Horne and A. Zeilinger, Phys. Today **46**, 22 (1993).

20. T. W. Marshall, in *Fundamental Problems in Quantum Physics* (M. Ferrero and A. van der Merwe, eds.) (Kluwer, Dordrecht, 1995) page 187.

21. P. Grangier, G. Roger and A. Aspect, Europhys. Lett. **1**, 173 (1986).

22. A. Casado, T. W. Marshall and E. Santos, preprint Univ. de Cantabria FMESC 3 (1995).

23. T. W. Marshall and E. Santos, Found. Phys. **18**, 185 (1988).

24. T. W. Marshall and E. Santos, Phys. Rev. A **39**, 6271 (1989).

25. A. J. Duncan and H. Kleinpoppen, in *Quantum Mechanics versus Local Realism* (F. Selleri, ed.) (Plenum, New York, 1988).

26. L. Mandel, Prog. in Optics **13**, 27 (1976).

27. E. Schroedinger, Proc. Camb. Phil. Soc. **31**, 555 (1935).

28. S. Haroche, M. Brune, J. M. Raimond and L. Davidovich, in *Fundamentals in Quantum Optics* (F. Ehlotzki, ed.) (Springer, Berlin, 1993).

29. S. J. D. Phoenix and S. M. Barnett, J. Mod. Opt. **40**, 979 (1993).

30. S. F. Huelga, T. W. Marshall and E. Santos, Phys. Rev. A **52**, R2497-R2500 (1995).

31. A. J. Leggett and A. Garg, Phys. Rev. Lett. **54**, 587 (1985).

32. G. Rempe, H. Walther and N. Klein, Phys. Rev. Lett. **58**, 353 (1987).

33. J. F. Clauser and M. A. Horne, Phys. Rev. D **10**, 526 (1974).