TIME AND BELL-TYPE INEQUALITIES

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Bell-type inequalities involve coincidence counting rates and are tested in coincidence experiments. Time is ontologically intertwined with the definition of coincidence. An example is proposed in which Bell's inequality is violated by a local model, simply because the interaction between a photon and a photomultiplier is not instantaneous, but requires a finite elapse of time. Models are only illustrative examples and there is no reason why they should reflect reality. We only claim that maybe Bell-type inequalities are not as general as they are supposed to be in ruling out local hidden variable theories.

It is commonly believed that EPR's [1] feeling about the incompatibility between local realistic theories (LRT) and quantum mechanics has been put into a very precise mathematical form by Bell's theorem [2].

This note wants to throw a doubt upon the generality of this conclusion. We believe, in fact, that Bell-type inequalities lack the idea of time, and this feature might be very important, for Bell-type inequalities involve coincidence counting rates and are tested in coincidence experiments: when the description of the EPR correlated system in terms of the (hidden) variable time is compulsory, the very definition of coincidence may show traces of this and many pitfalls may occur. An example will be given to illustrate these pitfalls.

We shall deal in the following with Clauser and Horne's (CH) [3] homogeneous [‡1] version of Bell's inequality, and carefully analyze two of their footnotes. Ch's homogeneous inequality reads

$$-p_{12}(\infty, \infty) \leqslant p_{12}(a, b) - p_{12}(a, b') + p_{12}(a', b) + p_{12}(a', b') - p_{12}(a', \infty) - p_{12}(\infty, b) \leqslant 0, \tag{1}$$

where

$$p_{12}(a, b) = \int_\Lambda p_1(\lambda, a) p_2(\lambda, b) \rho(\lambda)\, d\lambda \tag{2}$$

is the average probability of a coincidence count given the polarizer settings a and b (analogously a' and b'), ∞ denotes the absence of the polarizer and λ is a shorthand notation for one or more (hidden) parameters defining the state of the physical system investigated; λ spans the space Λ with a normalized distribution function $\rho(\lambda)$. In eq. (2) the probability of detecting two photons of a correlated couple is given in terms of elementary single count probabilities $p_1(\lambda, a)$ and $p_2(\lambda, b)$ according to Bell's well-known factorability condition.

Let us focus now our attention on footnote 13 of CH's paper. It is emphasized there that by writing in (2) the distribution function as $\rho(\lambda)$ instead of $\rho(\lambda | a, b)$, LRTs are not considered such that, roughly speaking, the orientation of the analyzer can influence the state of the photon

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[‡1] It is well known that CH also derived an inhomogeneous inequality, involving *single* count probabilities at either detector. The inequality (1) of the text, which is the one tested in actual experiments, is conceptually different from the homogeneous counterpart because the so-called no-enhancement hypothesis must be made in order to derive it. We will discuss enhancement later on in this letter. We want to stress, here, that the reason why we focus our attention on eq. (1) is that we aim at pointing out some pitfalls inherent in the coincidence mechanism, and therefore we do need coincidence counting probabilities, as in (1).

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source. We will refer to the above-mentioned CH note as CH1.

Some attempts have been made in the literature directly denying this hypothesis, but they do not seem very convincing. Indeed, anybody inclined to a local philosophy would opt for distribution functions which do not depend on the setting of the polarizers, mainly because of the last Orsay experiment [4]. We wish to show that a dependence of ρ on a and b can be a direct consequence of the presence of time in the correlation function.

Before considering an explicit example, it is worth mentioning the much less showy footnote 9 of ref. [3], which we shall refer to as CH2. By defining the coincide counting rate $N_{12}(a, b)$, CH write: The practical criterion for a "coincidence count" always involves a coincidence time window ω: pairs of counts separated in time by less than ω are defined to be coincident. This procedure may appear to make the definition of $N_{12}(a, b)$ ambiguous, since in general it will depend upon the experimenter's choice for ω. However, this dependence is usually insensitive to variations in ω which satisfy $s \ll \omega \ll 1/r$, where r is the average count rate at either detector, and s is the typical time separation of "true" coincidence pairs. Thus, we tacitly require the experimental arrangement to be such that this condition obtains (suitable source strength, time separation of pairs, etc.). If a sufficiently weak source is used, the ratio of "chance" coincident counts to "true" coincident counts can be made arbitrarily small, and the corresponding dead time can also be minimized.

We will now show an example in which a local realistic model can violate the CH inequality (1). We will then analyze the causes for this violation in the light of CH1 and CH2. We stress that the model we are going to propose must be considered as an example, and nothing more than that. No claim is made about its physical soundness and reliability. Our only purpose is to show that *maybe* Bell-type inequalities are less general than they are supposed to be, because local realistic models in which the time must be explicitly taken into account might not be that easy to rule out.

In atomic cascade tests of Bell's inequality (refs. [4–6] and references therein), a source emits pairs of correlated photons. If the first photon is emitted at the time $x = 0$, the probability for the second photon of the couple to be emitted between x and $x + \mathrm{d}x$ is

$$\rho_c(x)\,\mathrm{d}x = \mathrm{e}^{-x/T}\frac{\mathrm{d}x}{T}, \tag{3}$$

where T is the lifetime of the middle level of the involved cascade [‡2]. Let us now assume the existence of a hidden vector λ, perpendicular to the velocity of the photon, such that the interaction between the photon and the photomultiplier follows an exponential law; if the photon reaches the photomultiplier at the time $t = 0$, it leaves it (in the form of an electrical pulse) between t and $t + \mathrm{d}t$ with probability

$$\rho(t)\,\mathrm{d}t = \mathrm{e}^{-t/\tau(\lambda-l)}\frac{\mathrm{d}t}{\tau(\lambda-l)}, \tag{4}$$

The "lifetime" τ of the photon–photonmultiplier interaction is, therefore, a function of the angle between the vector λ and the photon polarization l [‡3]. One could think of an interaction mechanism, for instance, in which $\lambda - l \approx 0$ photons interact very rapidly with the electrons in the photomultiplier photocathode, whereas $|\lambda - l| \approx \pi/2$ ones need a certain elapse of time to interact, due to the "mismatch" of the two vectors. It is remarkable that taking into account real photomultipliers does not spoil any reasoning, as far as eq. (4) is concerned. If the quantum efficiency of the photomultiplier is η ($\eta < 1$), it is enough to multiply both sides of (4) by η, and continue similarly. We will not make any hypothesis about the mathematical expression of the photon transmission probability through a polarizer set at an angle a.

‡2 In the Stirling experiment [6] the correlated-photon emissions are simultaneous (second-order decay process). It suffices, in that case, to take the limit $T \to 0$ in the final formula.

‡3 We stress two points: (i) the model just considered is compatible with any wave–particle picture, dualistic or not, of the photon; (ii) we might have assumed eq. (4) to determine the interaction between the photon and the polarizer. In that case, l would have been the orientation of the polarizer and λ *could* have been the photon polarization before the interaction. The reasoning which follows would still hold, in outline, in this latter case. We chose the former one for reasons that will be clear later on.

We assume, for the sake of simplicity, that two correlated photons leave the source with the same polarization l, uniformly distributed, and with hidden vectors λ, λ'. For the sake of generality, let the photon transmission probability be a function of $(l-a)$ as well as $(\lambda - a)$:

$$p_1 = p_1(l-a, \lambda - a) \tag{5}$$

(analogously $p_2 = p_2(l-b, \lambda' - b)$); we keep the indices 1, 2 to label the photon transmission probability only for convenience; of course, the mathematical expression of p_1 and p_2 is exactly the same. In the notations

$$b - a = \alpha, \tag{6}$$

$$\lambda - a = \phi, \quad \lambda' - b = \phi', \tag{6a}$$

$$l - a = \theta, \quad l - b = \theta - \alpha, \tag{6b}$$

the probability of a coincidence count is

$$p_{12}(\alpha) = \int_0^\pi \frac{d\theta}{\pi} \int_0^\pi d\phi' \int_0^\pi d\phi\, \sigma(\phi, \phi') p_1(\theta, \phi) \times p_2(\theta - \alpha, \phi') \int_0^\infty dt\, \rho_1(t) \int_0^\infty dx\, \rho_e(x) \times \int_{t-x}^{t+\omega-x} dt'\, \theta(t') \rho_2(t'), \tag{7}$$

where $\sigma(\phi, \phi')$ is the (normalized) distribution function of ϕ and ϕ', p_1 and p_2 are defined in eq. (5), $\rho_e(x)$ in eq. (3), and ρ_i $(i = 1, 2)$ (eq. (4)) refer to the ith photon of the atomic cascade:

$$\rho_1(t) = e^{-t/\tau(\phi)}/\tau(\phi), \quad \rho_2(t) = e^{-t'/\tau(\phi')}/\tau(\phi')$$

(obviously, after having crossed the polarizers, the photon polarizations are a and b, respectively). The constraints

$$t \leqslant t' + x \leqslant t + \omega \tag{8}$$

are responsible for the "anomalous" limits in the last integral in (7). This is due to the experimental definition of coincidence. The second photon of a correlated pair must be turned into an electrical pulse after the first one ($t' + x \geqslant t$) but within the coincidence window ω ($t' + x \leqslant t + \omega$). The presence of the unit step function θ in the last integral is required because the integration limits may become negative.

It may seem unbelievable, but the constraint (8) can cause a violation of the CH inequality (1).

Eq. (7) is very difficult to tackle in its general form. Let us analyze, therefore, a particular case:

$$\sigma(\phi, \phi') = \delta(\phi - \phi')/\pi. \tag{9}$$

Eq. (7) becomes

$$p_{12}(\alpha) = \int_0^\pi \frac{d\theta}{\pi} \int_0^\pi \frac{d\phi}{\pi} p_1(\theta, \phi) p_2(\theta - \alpha, \phi - \alpha) \times \int_0^\infty dt\, \rho_1(t) \int_0^\infty dx\, \rho_e(x) \times \int_{t-x}^{t+\omega-x} dt'\, \theta(t') \rho_2(t'), \tag{10}$$

with

$$\rho_1(t) = e^{-t/\tau(\phi)}/\tau(\phi),$$

$$\rho_2(t') = e^{-t'/\tau(\phi-\alpha)}/\tau(\phi - \alpha),$$

and after a straightforward but rather lengthy calculation one arrives at

$$p_{12}(\alpha) = \int_0^\pi \frac{d\theta}{\pi} \int_0^\pi \frac{d\phi}{\pi} p_1(\theta, \phi) p_2(\theta - \alpha, \phi - \alpha) \times f(\phi, \phi - a), \tag{11}$$

where

$$f(\phi, \phi - \alpha) = (1 - e^{-\omega/\tau'}) \frac{\tau'^2}{(\tau + \tau')(\tau' - T)} + (1 - e^{-\omega/T}) \frac{T^2}{(\tau + T)(T - \tau')} \tag{12}$$

($\tau = \tau(\phi)$ and $\tau' = \tau(\phi - \alpha)$). Let us now compare eqs. (2) and (11), and let us consider the factor $\pi^{-2} f(\phi, \phi - \alpha)$ as a distribution function ρ [‡4]: CH1 is violated, f depending on the polarizer settings a and b! In a previous paper [7], we pointed out that in Bell-inequality tests, the coincidence counting electronics could act as a "source of nonlocality" in the correlation function. In this case, the conclusion is somehow stronger. It is certainly true that nonlocal effects in the correlation function arise from the coincidence counting

[‡4] f is normalized to some constant $c < 1$ (from its definition and/or eqs. (7), (8)), and not to 1; but this does not alter any conclusion.

electronics, but furthermore, even if $\omega \to \infty$, f does not approach unity. This strange effect is to be ascribed to the fact that the electronics always "knows" which photon arrived first (the lower limit in the last integral in (7) and (10) does not vanish even if $\omega \to \infty$), and the coincidence probability shows traces of that. Is it possible to violate the CH inequality (1) by means of a correlation function of the type (11)? In atomic cascade experiments, the Freedman inequality [8] is usually tested:

$$|p_{12}(\pi/8) - p_{12}(3\pi/8)|\,p_0 \leqslant \tfrac{1}{4}, \tag{13}$$

where p_0 is the coincidence counting probability in absence of polarizers; (13) follows directly from (1) after a short algebraic manipulation. Let us choose, for instance,

$$p_1(\theta, \phi) = p_2(\theta - \alpha, \phi - \alpha) = \tfrac{1}{2} \tag{14}$$

in eq. (11). This choice corresponds to considering circularly polarized photons which may not even be spin-correlated at all. Then (11) becomes

$$p_{12}(\alpha) = \tfrac{1}{4}\int_0^{\pi} \frac{d\phi}{\pi} f(\phi, \phi - \alpha). \tag{15}$$

Even for the oversimplified version (15) of eq. (7), a complete discussion is a hard mathematical task due to the high nonlinearity of the function τ in (12). But since we make no pretentions for our rough illustrative model to be considered as a sound physical theory, an exhaustive analysis of (15) is out of place in this context. We have therefore limited ourselves to a very partial analysis of (15), by computing numerically the integral (15) for different choices of the function τ.

Let us set our minds on calcium atomic cascade experiments, and in particular on the Orsay ones (ref. [4] and references therein). In that case, $T \approx 5$ ns and $\omega \approx 19$ ns. Functions such as

$$\begin{aligned} &\tau(\phi) = A + B(\phi - k\pi)^2, \quad A \geqslant 0,\ B \geqslant 30 \text{ ns},\\ &\tau(\phi) = A + B|\phi - k\pi|^3, \quad A \geqslant 0,\ B \geqslant 24 \text{ ns},\\ &\tau(\phi) = A + B(\phi - k\pi)^4, \quad A \geqslant 0,\ B \geqslant 20 \text{ ns},\\ &\dots\\ &\text{if } |\phi - k\pi| < \tfrac{1}{2}\pi, \quad k = 0, \pm 1, \pm 2, \dots, \end{aligned} \tag{16}$$

lead to violations of (13). This is quite remarkable, if we think that the correlation function in the hypothesis (14) would be

$$p_{12}(\alpha) = \tfrac{1}{4} \tag{17}$$

if retardation effects had not been taken into account. In other words, the correlation function may be quite flat, like in (17), and a violation of (13) can be *entirely* due to delay effects. We do hope that this is a convincing proof that Bell-type inequalities do not fully take into account the pitfalls that may occur from the explicit presence of time in the correlation function.

Why is (13) violated by a local realistic model? The reason for this can be found in CH2. In our example, the strong inequality $s \ll \omega$ does not hold. Indeed, from eq. (16),

$$\tau(\phi) = A + B|\phi - k\pi|^n \quad \text{if } |\phi - k\pi| < \tfrac{1}{2}\pi \quad (k \in \mathbb{Z},\ n \in \mathbb{N})$$

it follows that

$$\langle \tau(\phi) \rangle_\phi = A + B\frac{\pi^n}{2^n(n+1)} \approx B$$

(if $A \ll B$ and n is not too big) and B must be of the order of ω, if we want the functions in (16) to violate (13). This means that the CH inequality can be violated if photons stop somewhere for a time of the order of the coincidence window. In this case, the "typical time separation of true coincidence pairs" s becomes of the order of ω and CH2 does not hold anymore. As we have seen before, effects of this kind can even provoke the failure of the "milestone" CH1.

Two other authors [9] have pointed out coincidence-window effects in the correlation function. We have shown that in their papers a violation of CH2 exists [10] and criticized their models because the coincidence window is chosen ad hoc to violate Bell's inequality. But we realize now that some of their conclusions are similar to ours. Moreover, attention to the importance of time delays has also been drawn by Santos [11]. In a certain sense, this paper is the consequence of a careful reconsideration of refs. [7,9–11]. Apparently nonlocal effects, coincidence-window effects and violations of CH1 and CH2 can be due to the lack of time in the Bell-type inequality analysis of the EPR problem.

Some more words about (7). Set

$$\sigma(\phi, \phi') = \pi^{-1}\delta(\phi - \phi' + \tfrac{1}{8}\pi),$$

take

$$p_1 = p_1(\theta), \quad p_2 = p_2(\theta - \alpha),$$

then a straightforward calculation leads to the factorized expression

$$p_{12}(\alpha) = \int_0^\pi \frac{d\theta}{\pi} p_1(\theta) p_2(\theta - \alpha) \times \int_0^\pi \frac{d\phi}{\pi} f(\phi, \phi - \alpha + \tfrac{1}{8}\pi). \quad (18)$$

Choose

$$\int_0^\pi \frac{d\theta}{\pi} p_1(\theta) p_2(\theta - \alpha) = \tfrac{1}{2} - \alpha/\pi,$$

which is the "spin 1 version" of Bell's linear correlation function [2]. The quantum mechanical prediction for the calcium atomic cascade correlation is

$$p_{12}^{QM}(\alpha)/p_0^{QM} = \tfrac{1}{4}(1 + \cos 2\alpha).$$

In our case, p_0 may be considered as a mere normalization factor. We can think, for instance, that photons which have crossed no polarizer interact with the photomultiplier with a lifetime τ_0 = constant. In that case,

$$p_0 = (1 - e^{-\omega/\tau_0}) \frac{\tau_0}{2(\tau_0 - T)} + (1 - e^{-\omega/T}) \frac{T^2}{T^2 - \tau_0^2}, \quad (19)$$

which is a monotonically decreasing function of τ_0 (the discontinuity in $\tau_0 = T$ is eliminable). We can, therefore, use the parameter τ_0 in (19) to suitably normalize $p_{12}(\alpha)$ in (18). The question is: is it possible to find any function τ such that

$$\frac{1}{p_0}\left(\frac{1}{2} - \frac{\alpha}{\pi}\right) \int_0^\pi \frac{d\phi}{\pi} f(\phi, \phi - \alpha + \tfrac{1}{8}\pi) \approx \tfrac{1}{4}(1 + \cos 2\alpha), \quad (20)$$

so that the quantum mechanical curve can be closely approached? Fig. 1 answers this question. We stress that no claim is being made about the physical reliability of the model proposed. The unnatural assumption about $\sigma(\phi, \phi')$ and the absence of any physical mechanism that could justify the hypothesis (4) make it a purely ad hoc model, devised in order to reproduce the quantum mechanical predictions. Our only aim is to show that there are local realistic models which can violate the CH inequality, and to point out some limits of inequalities of this kind. It is enough to devise a mechanism that will accomplish delay effects in photon detection so as to get rid of some and not others, due to the finiteness of the coincidence window (that is what eqs. (4) and the last integral in (7) actually do), and Bell-type inequalities do not hold anymore. It is maybe worth emphasizing that a variation of the coincidence window ω within reasonable limits does not point out the presence of delay effects of the kind proposed. We varied ω within the range 10–40 ns (experimental value: ω = 19 ns) ad evaluated numerically eqs. (15), (20). The coincidence probabilities underwent variations of the order of 2–5% only! This means that delay effects due to hidden variables, if existing at all, might be very difficult to bring to light by varying the coincidence window. A change in ω of more than 100% would affect the experimental data in a negligible manner.

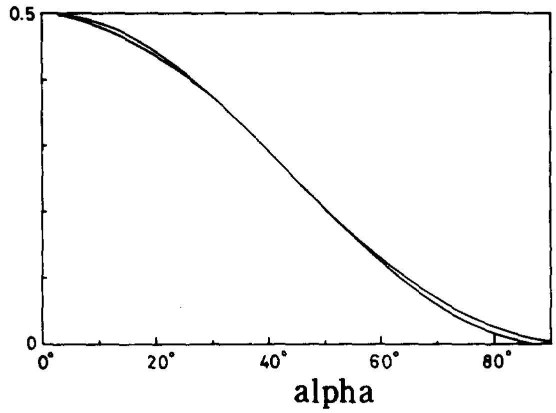

Fig. 1. Comparison between the left- and right-hand sides of eq. (20) for $\tau(\phi) = 0.02 + 22|\phi|^3$ (τ in ns, ϕ in rad). The quantum mechanical value is bigger for $0° < \alpha < 7°$ and $31° < \alpha < 45°$.

We would like to end off this note by pointing out some quite unexpected connections between the proposed model and the so-called no-enhancement hypothesis (NEH) [3,6]. Let λ denote the state of the photon. If $p_i(\lambda, \infty)$ denotes the probability of a count from detector i ($i = 1, 2$) when the polarizer is absent and $p_i(\lambda, a)$ the same probability when the polarizer setting is a, then the NEH is:

$$p_i(\lambda, a) \leqslant p_i(\lambda, \infty), \quad \forall\lambda. \tag{21}$$

The NEH has recently been studied by several authors [12–14], and some remarkable results have been achieved; many people consider it to be the last loophole in favour of LRT.

Consider now eq. (4), which we rewrite here in the notation (6a):

$$\rho(t)\ \mathrm{d}t = \mathrm{e}^{-t/\tau(\phi)} \frac{\mathrm{d}t}{\tau(\phi)},$$

and assume that a photon that has crossed no polarizer interacts with the photomultiplier with a "lifetime" τ_0:

$$\rho_0(t)\ \mathrm{d}t = \mathrm{e}^{-t/\tau_0} \frac{\mathrm{d}t}{\tau_0}.$$

If a coincidence window allows photon detections between $t = 0$ and $t = \omega$, then the probability for a photon to be detected is

$$p(\phi) = \tfrac{1}{2} \int_0^\omega \rho(t)\ \mathrm{d}t = \tfrac{1}{2}(1 - \mathrm{e}^{-\omega/\tau(\phi)}), \tag{22}$$

if the photon has crossed a polarizer set at a ($\phi = |\lambda - a|$), and

$$p_0 = \int_0^\omega \rho_0(t)\ \mathrm{d}t = 1 - \mathrm{e}^{-\omega/\tau_0} \tag{23}$$

if it has not crossed any. It is then very easy to show that if

$$\tau_0 > \omega/\ln 2$$

then, from (22), (23), it is possible to find functions $\tau(\phi)$ such that

$$p(\phi_0) > p_0 \quad \text{for some } \phi_0. \tag{24}$$

Choose for instance a function $\tau(\phi)$ which vanishes (or becomes very small) for some $\phi = \phi_0$; then ϕ_0 satisfies (24). Eqs. (21) and (24) are in manifest disagreement. The conclusion is that *enhancement* in the photon detection probability *can arise from delay effects in the photon interactions.*

But there is even more to that. A theorem has been recently proven [14] stating that, *if analyzers work symmetrically*, even LRTs that admit enhancement in the photon detection cannot be equivalent to quantum mechanics. A glance at formulae (11), (12), (as well as (7)) shows that delay effects in photon interactions can cause quite an asymmetry in the correlated photons detection. Therefore, the main hypothesis of Caser's theorem fails and one could even aspire to reproduce *identically* the quantum mechanical predictions for the joint detection probabilities. Home and Marshall [13] also suggested a "symmetrization procedure" which is able to provide a counterexample to Caser's theorem. It is noteworthy, however, that in the example we propose, one need not symmetrize the model, the asymmetry (and the violation of Caser's hypothesis) being ontological in a coincidence detection process, when one or more (hidden) variables depend on time.

We wish to conclude this letter by summarizing what has been achieved. We propose that the CH inequality (and, in general, Bell-type inequalities) does not fully take into account the possibility that some (hidden) parameters can cause delays in photon interactions. If this does actually happen in nature, then the hypotheses CH1 and CH2, necessary to derive the inequality (1), are not general enough, and the CH inequality is not a necessary constraint for *all* the LRTs. Moreover, the possibility of an unexpected connection seems to exist between retardation effects and enhancement in the photon detection probability. Furthermore, the main hypothesis necessary to derive a theorem [14] which discriminates quantum mechanics and LRTs which admit enhancement does not necessarily hold if delay phenomena take place. We believe that further investigation is required to analyze the role of time in Bell-type inequalities. In particular, a recent paper by Franson [15] seems very interesting to us, because the possibility is considered that the outcome of an event is not determined until some time after its occurrence ("delayed determinism theories"). Many authors think there is no problem in connection with EPR. Others believe there has never been one. In our opinion, the EPR puzzle is far from being solved.

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References

[1] A. Einstein, B. Podolsky and N. Rosen, Phys. Rev. 47 (1935) 777.
[2] J.S. Bell, Physics 1 (1964) 195.
[3] J.F. Clauser and M.A. Horne, Phys. Rev. D 10 (1974) 526.
[4] A. Aspect, J. Dalibard and G. Roger, Phys. Rev. Lett. 49 (1982) 1804.
[5] J.F. Clauser and A. Shimony, Rep. Prog. Phys. 41 (1978) 1681.
[6] W. Perrie, A.J. Duncan, H.J. Beyer and H. Kleinpoppen, Phys. Rev. Lett. 54 (1985) 1790.
[7] S. Pascazio, Phys. Lett. A 111 (1985) 339; in: Microphysical reality and quantum formalism, eds. G. Tarozzi and A. van der Merwe (Reidel, Dordrecht), to be published.
[8] S.J. Freedman, Lawrence Berkeley Lab. Rep. LBL 391 (1972).
[9] G.C. Scalera, Lett. Nuovo Cimento 38 (1983) 16; 40 (1984) 353;
S. Notarrigo, Nuovo Cimento B 83 (1984) 173.
[10] S. Pascazio, in: Italian studies in the foundations of quantum mechanics, eds. G. Tarozzi and A. van der Merwe (Reidel, Dordrecht), to be published.
[11] E. Santos, Phys. Lett. A 101 (1984) 379.
[12] T.W. Marshall, E. Santos and F. Selleri, Phys. Lett. A 98 (1983) 5;
T. W. Marshall, Phys. Lett. A 99 (1983) 163; 100 (1984) 225;
A. Garuccio and F. Selleri, Phys. Lett. A 103 (1984) 99;
T.W. Marshall and E. Santos, Phys. Lett. A 107 (1985) 164;
F. Selleri, Phys. Lett. A 108 (1985) 197;
M. Ferrero and E. Santos, Phys. Lett. A 108 (1985) 373.
[13] D. Home and T.W. Marshall, Phys. Lett. A 113 (1985) 103.
[14] S. Caser, Phys. Lett. A 102 (1984) 152.
[15] J.D. Franson, Phys. Rev. D 31 (1985) 2529.