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Publication Date

1973-08-01

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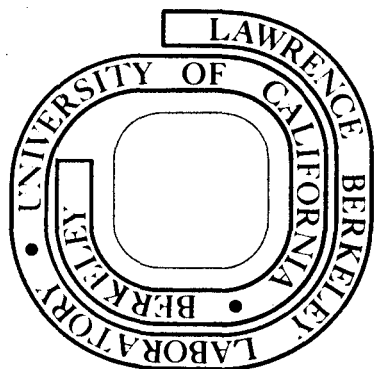
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August 2, 1973

Prepared for the U.S. Atomic Energy Commission
under Contract W-7405-ENG-48

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EXPERIMENTAL CONSEQUENCES OF OBJECTIVE LOCAL THEORIES II*

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ABSTRACT

The experimental status of objective local theories is discussed. No conclusive test of these has yet been performed. A plausible supplementary assumption is formulated, weaker than any previously exhibited. Objective local theories and local hidden-variable theories consistent with this assumption are found to be in violation of recent experimental data. It is also shown that at least some supplementary assumption is needed to demonstrate an incompatibility between these data and such theories, although experiments are possible for which this is not the case.

INTRODUCTION

The possibility that the statistical features of quantum mechanics might be described in terms of an underlying deterministic substructure has been repeatedly suggested in the literature. Such covering theories of quantum mechanics, when consistent with macrocausality are generally called local hidden-variable theories (LHVT's). Bell¹ has recently shown that no LHVT can reproduce all of the statistical predictions of quantum theory for a Gedankenexperiment of Bohm. Clauser, Horne, Shimony, and Holt² (CHSH) extended this result to a consideration of realizable experiments. They formulated a reasonable supplementary assumption concerning these theories, which allowed them to propose an actual experimental test. Freedman and Clauser³ performed the suggested experiment, and thus demonstrated the untenability of any LHVT which satisfies their supplementary condition.

Two questions were left conspicuously unanswered by the analysis of Bell and CHSH: (1) Are experiments possible which test LHVT's not constrained by the supplementary assumption of CHSH; and (2) do the experimental results of Freedman and Clauser apply to theories more general than LHVT's constrained by this supplementary CHSH assumption?

In part I of this article⁴ the first question was answered in the affirmative. That part provided a definition of objective local theories (OLT's). The deterministic LHVT's were seen to be a subclass of these. It was then shown that realizable experiments may be performed to test OLT's

with no supplementary assumptions. Obviously then no supplementary assumptions will be needed to apply such results to the more restrictive LHVT's.

In this part we answer the second question, also in the affirmative. To do this we first review the experiment of Freedman and Clauser, and show that their results unfortunately do not meet the requirements discussed in part I for a conclusive test of OLT's. Such a test will require a correlation experiment, employing highly efficient analyzers and detectors, performed upon well-collimated two-body dissociation products. No such experiment has yet been performed. Next, we state a supplementary assumption, weaker than that of CHSH, and prove that it is sufficient to guarantee an incompatibility of the predictions by OLT's and the data of Freedman & Clauser. Finally, we demonstrate the necessity of at least some supplementary assumption for this experiment by producing a counter-example which agrees with the definitions of an OLT (and of a LHVT), but reproduces exactly the quantum mechanical prediction (as well as the observed results) for this experiment. We notice, in passing, that any radiation theory employing the classical Maxwell equations satisfies our new supplementary assumption, and is summarily ruled out by experiment.

SUMMARY OF EXPERIMENTAL RESULTS

Following the suggestion by CHSH, Freedman and Clauser measured the linear polarization correlation of the successive photons emitted during a $J = 0 \rightarrow J = 1 \rightarrow J = 0$ atomic cascade. A diagram of their

apparatus is shown in Fig.1. In it the decaying atoms were viewed by two symmetrically placed optical systems, each consisting of two lenses, a wavelength filter, a rotatable and removable linear polarizer, and a single-photon detector. They measured the following quantities:

$R(\phi)$: The coincidence rate for two-photon detection, as a function of the angle ϕ between the planes of linear polarization defined by the orientation of the inserted polarizer;

R_1 : The coincidence rate with polarizer B removed;

R_2 : The coincidence rate with polarizer A removed;

R_0 : The coincidence rate with both polarizers removed.

The quantum mechanical predictions for these rates are given by

$$R(\phi) = [\frac{1}{4} Q_A Q_B p(\theta) q(\theta) S(\epsilon_+^A \epsilon_+^B + \epsilon_-^A \epsilon_-^B F(\theta) \cos 2 \phi)] R_e \quad (1)$$

$$R_1 = [\frac{1}{2} Q_A Q_B p(\theta) q(\theta) S \epsilon_+^A] R_e, \quad (2)$$

$$R_2 = [\frac{1}{2} Q_A Q_B p(\theta) q(\theta) S \epsilon_+^B] R_e, \quad (3)$$

$$R_0 = [Q_A Q_B p(\theta) q(\theta) S] R_e, \quad (4)$$

where R_e is the rate at which the source lamp emits cascade pairs, Q_A and Q_B are the effective quantum efficiencies of the detectors,⁵ θ is the half-angle of the cones subtended by the primary lenses, and where $\epsilon_+^i = \epsilon_M^i + \epsilon_m^i$ and $\epsilon_-^i = \epsilon_M^i - \epsilon_m^i$ ($i = A, B$). Here ϵ_M^i (ϵ_m^i) is the efficiency of the i^{th} polarizer for light polarized parallel (perpendicular) to the polarizer axis. The function $p(\theta)$ is the probability that the first emission of a cascade pair enters the

primary lense of optical assembly A;

$$p(\theta) = \frac{1}{2} (1 - \cos \theta). \quad (5)$$

The function $q(\theta)$ is the conditional probability (angular correlation factor) that the second emission will enter assembly B, given that the first emission enters assembly A;

$$q(\theta) = \frac{3}{8} \frac{[(G_2(\theta))^2 + \frac{1}{2}[G_3(\theta)]^2]}{(1 - \cos \theta)}. \quad (6)$$

The function,

$$F(\theta) = \frac{2 [G_1(\theta)]^2}{[(G_2(\theta))^2 + \frac{1}{2}[G_3(\theta)]^2]} \quad (7)$$

represents a depolarization due to the noncollinearity of the two photons, and approaches unity for infinitesimal detector solid angles (i.e. as $\theta \rightarrow 0$). The functions G_1 , G_2 , and G_3 are given in Ref. 2. The quantity S (less than or equal to one) represents a loss of the second photon of the cascade, given that the first was detected. For example, the loss may be due to imperfect simultaneous focusing of the two optical systems, the existence of alternative de-excitation routes to other final states, or perhaps other causes. The quantum mechanical predictions for the singles rates at the A and B detectors in this experiment are easily shown to be bounded by $r_A \geq \frac{1}{2} p(\theta) \epsilon_+^A R_e$ and $r_B \geq \frac{1}{2} p(\theta) \epsilon_+^B R_e$, respectively. (In practice, detector dark currents, the existence of alternative final state de-excitation routes, stray light, imperfect focusing,

etc, made the observed singles rates significantly exceed these lower bounds.)

Previous work has considered the experimental consequences of the class of theories implied by the following assumptions:

- A. The two photons propagate as separate localized particles (particle assumption).
- B. A deterministic or quasideterministic selection process, transmission or no-transmission, occurs for each photon at a polarizer (determinism assumption).⁶
- C. Locality requires that this selection not depend upon the orientation of the distant polarizer (locality assumption).
- D. All photons incident on a detector have a detection probability that is independent of whether or not the photon has passed through the associated polarizer (detector efficiency assumption).

Bell and CHSH examined the consequences of assumptions A-C, and CHSH added condition D to allow the experiment to be performed with current technology. These assumptions are sufficient to derive the experimental prediction^{2,3}

$$-1 \leq \Delta(\theta) \leq 0 \quad (8)$$

where

$$\Delta(\phi) = \frac{3R(\phi)}{R_0} - \frac{R(3\phi)}{R_0} - \frac{R_1 + R_2}{R_0} \quad (9)$$

The quantum mechanical prediction for $\Delta(\phi)$ obtained from Eqs. (1)-(7) and (9) has a maximum at $\phi = 22.5^\circ$ and a minimum at $\phi = 67.5^\circ$. Furthermore, these predicted extrema violate inequality (8) provided

$$\epsilon_+ \geq \frac{2}{\sqrt{2}(\epsilon_-/\epsilon_+)^2 F(\theta) + 1} \quad (10)$$

Here we have assumed for simplicity that $\epsilon_+^A = \epsilon_+^B \equiv \epsilon_+$ and $\epsilon_-^A = \epsilon_-^B \equiv \epsilon_-$. Note that the detector efficiencies, Q_A and Q_B , the probability $p(\theta)$, and the angular correlation factor $q(\theta)$, do not enter this expression. The conditions in the Freedman-Clauser experiment, $\epsilon_+ \approx 1.00$, $\epsilon_- \approx 0.94$, $\theta \approx 30^\circ$ and $F(30^\circ) \approx 0.99$, satisfy Ineq.(10). The experimental results $\Delta(22.5^\circ) = 0.104 \pm 0.026$ and $\Delta(67.5^\circ) = -1.097 \pm 0.018$, are consistent with the quantum mechanical prediction but in clear violation of Ineq.(8) and thus A-D. We shall see in the next two sections, however, that they do not violate the predictions given in part I for a general OLT or even some LHVT's which violate assumption D.

EXPERIMENTAL STATUS OF OBJECTIVE LOCAL THEORIES

In this section we summarize very briefly the results of part I. In that paper we introduced $p(a, \lambda_A)$ and $p(b, \lambda_B)$ as the locally defined probabilities of a count in a given interval of time. We postulated that these depended only upon the local objective properties of nature, and the adjustable analyzer orientations a and b , as indicated by the respective arguments. (For linear polarization of photons, a and b define planes of polarization. The angles between these planes, and some reference plane including the apparatus axis, are denoted by a and b .) Thus, assumptions A and B were effectively replaced by the more general assumption of the objectivity of the properties λ_A and λ_B . Assumption C was weakened to include any processes at the detectors and analyzers, consistent with locality. For fixed a , λ_A , b and λ_B we argued that p_A and p_B must be independent probabilities. Simply from the definability of these quantities and the associated independence, the following inequality was derived which constrains observable "overlap coincidence rates":

$$R(a,b) - R(a,b') + R(a',b) + R(a',b') - r_A(a') - r_B(b) \leq 0. \quad (11)$$

In this correspondence, $R(a,b)$ is the coincidence rate as a function of a and b ; r_A and r_B are respectively the singles rates at the A and B photon detectors.

Let us now compare the quantum mechanical predictions for this experiment with those by OLT's. It should first be noted that the rather involved sequence outlined in Part I for the precise, local definition of a coincidence was not followed by Freedman and Clauser. Also the apparatus orientations were not readjusted while the photons were in flight as was specified by part I.⁷ The first difference will change the data analysis slightly, but it is highly doubtful that such a difference might alter the conclusion. The second difference is probably more open for criticism. A realization of this experimental requirement is indeed within the capability of present technology, and is perhaps worth-while.

Assuming irrelevant, for the moment, the differences between the experimental procedures, we compare the predictions (1)-(4) with the constraint imposed by Ineq. (11). For simplicity we further assume that $Q_A = Q_B = Q$, $\epsilon_-^A = \epsilon_-^B = \epsilon_-$ and $\epsilon_+^A = \epsilon_+^B = \epsilon_+$. We find the condition for violation of Ineq. (11) to be

$$S \cdot Q \cdot q(\theta) \cdot \epsilon_+ \geq \frac{2}{\sqrt{2} (\epsilon_- / \epsilon_+)^2 F(\theta) + 1} \quad (12)$$

Because of the relatively small value of $q(\theta)$, this condition is not satisfied for any value of the detectors' half-angle θ , even when the polarization analyzer and detector efficiencies are ideal ($Q = \epsilon_+ = \epsilon_- = S = 1$). Therefore at least for cascade photon experiments, the quantum mechanical predictions and observed results are compatible with the predictions

for a general OLT, even in the domain of ideal apparatus.

The essential difficulty is the smallness of $q(\theta)$. This is due to the fact that an atomic cascade is a three-body decay, the ground-state atom being the third body. However, for a correlated two-body decay, quantum mechanics predicts an incompatibility. For example, the annihilation of ground state positronium into two $1/2$ - MeV γ rays, or the dissociation of a spin-zero molecule into two spin- $1/2$ particles produces correlations analogous to those for cascade photons. For these we have $g(\theta) \approx 1$, even for small θ , (provided the center-of-mass velocity of the decaying object is sufficiently small). However the use of these elements in an experimental test of OLT's (and LHVT's) still requires the rather high efficiencies of the detectors and analyzers discussed in part I.⁸ Fortunately, there appears to be no a priori reason why such experimental conditions cannot be achieved in practice.

EXPERIMENTAL STATUS OF OBJECTIVE LOCAL THEORIES WITH A SUPPLEMENTARY ASSUMPTION.

We now present a new supplementary condition, weaker than D, and prove that it is sufficient to yield an incompatibility of any OLT with the existing experimental results. The assumption is that for the objective properties λ_A (or λ_B), the probability of a count at A (or B) in a given interval of time with a polarizer in place is less than the corresponding probability with the polarizer simply removed. We denote the condition in which the

polarizer is removed by the symbol ∞ . Thus, we assume that for all λ_A and λ_B the following inequalities hold:

$$\begin{aligned} 0 \leq p_A(a' \text{ or } a, \lambda_A) &\leq p_A(\infty, \lambda_A) \leq 1 \\ 0 \leq p_B(b' \text{ or } b, \lambda_B) &\leq p_B(\infty, \lambda_B) \leq 1 \end{aligned} \quad (13)$$

We call Ineq.(13) the "no-enhancement" assumption. In other words, with collimators, etc. fixed, the action of a polarizer is to attenuate and not amplify anything passing through it. This assumption is observably true for the averages over λ_A and λ_B . In fact, if the quantum statistical predictions are to hold, we have the much stronger conditions $\langle p_A(a, \lambda_A) \rangle \leq \langle p_A(\infty, \lambda_A) \rangle$, etc. But our assumption is made for every λ_A and λ_B , and is not directly testable. It will be discussed further in a later section.

Inequality (11) was derived by the use of a theorem, the proof of which may be found in the Appendix of part I. The theorem states that given any six numbers $x_1, x_2, y_1, y_2, X,$ and Y , such that $0 \leq x_1 \leq X, 0 \leq x_2 \leq X, 0 \leq y_1 \leq Y, 0 \leq y_2 \leq Y$, then they are further constrained by the inequality

$$-XY \leq x_1y_1 - x_1y_2 + x_2y_1 + x_2y_2 - Yx_2 - Xy_1 \leq 0. \quad (14)$$

Taking $X = p_A(\infty, \lambda_A), Y = p_B(\infty, \lambda_B), x_1 = p_A(a, \lambda_A),$

$$x_2 = p_A(a', \lambda_A), y_1 = p_B(b, \lambda_B), \text{ and } y_2 = p_B(b', \lambda_B),$$

we have

$$\begin{aligned} p_A(\infty, \lambda_A) p_B(\infty, \lambda_B) &\leq p_A(a, \lambda_A) p_B(b, \lambda_B) - p_A(a, \lambda_A) p_B(b', \lambda_B) \\ &+ p_A(a', \lambda_A) p_B(b, \lambda_B) + p_B(a', \lambda_A) p_B(b', \lambda_B) - p_A(a', \lambda_A) p_B(\infty, \lambda_B) \\ &- p_A(\infty, \lambda_A) p_B(b, \lambda_B) \leq 0. \end{aligned} \quad (15)$$

Following part I, we multiply Ineq.(15) by $\rho(\lambda_A, \lambda_B)$ (the ensemble probability distribution of λ_A and λ_B), integrate over the full domain of λ_A and λ_B , and divide by the detector clock pulse length τ . We then identify the quantities (see part I):

$$R(a,b) = \frac{1}{\tau} \iint P_A(a, \lambda_A) P_B(b, \lambda_B) \rho(\lambda_A, \lambda_B) d\lambda_A d\lambda_B \quad (16)$$

$$R_1(a) = R(a, \infty) \quad (17)$$

$$R_2(b) = R(\infty, b) \quad (18)$$

$$R_0 = R(\infty, \infty) \quad (19)$$

as the overlap coincidence rates defined in part I, corresponding to the measured rates defined by Eqs. (1)-(4). Using these definitions, we obtain

$$R_0 \leq R(a,b) - R(a,b') + R(a',b) + R(a',b') - R_1(a') - R_2(b) \leq 0. \quad (20)$$

R_1 and R_2 were found experimentally to be independent of a and b , and R was likewise found to depend only upon the angle between the analyzer planes, $\phi = a - b$. Thus Eq.(20) can be written in the simpler form

$$-1 \leq \frac{3R(\phi)}{R_0} - \frac{R(3\phi)}{R_0} - \frac{R_1 + R_2}{R_0} \leq 0. \quad (21)$$

Inequality (21) is identical to the previous Ineq (8) which was violated by existing experimental results.

NECESSITY OF AN ADDITIONAL ASSUMPTION

In this section we demonstrate the necessity of at least some additional assumption for the Freedman-Clauser experiment by

producing an explicit model which satisfies the definitions given in part I for an OLT (and for a LHVT), but reproduces the quantum mechanical predictions (1)-(4) (and the results of that experiment). For simplicity, we shall exhibit the model only for the idealized case in which the detectors subtend infinitesimal solid angles [$\theta \rightarrow 0$, $F(\theta) \rightarrow 1$]. The additional complexity required for an extension of the model to the finite solid angles of the actual experiment [$\theta = 30^\circ$, $F(\theta) = 0.99$] introduces nothing, but obscures the essential point. The predictions to reproduce are then summarized by

$$R(\phi)/R_0 = \frac{1}{4} [\epsilon_+^A \epsilon_+^B + \epsilon_-^A \epsilon_-^B \cos(2\phi)] \quad (22)$$

$$R_1/R_0 = \epsilon_+^A/2 \quad (23)$$

$$R_2/R_0 = \epsilon_+^B/2. \quad (24)$$

In our model a source atom emits a pair of "particle-like" photons. Each emission pair leading to a coincident count thus consists of one particle traveling along the +z axis to detector A, and the other particle traveling along -z axis to detector B. Both members of either pair possess a common state variable λ , which is simply an azimuthal angle; that is, it specifies a direction perpendicular to its flight axis from the reference axis (see Fig. 2). The complete ensemble of emission pairs is characterized by a normalized probability density which is isotropic,

$$\rho(\lambda) d\lambda = \frac{d\lambda}{2\pi}, \text{ for } 0 \leq \lambda < 2\pi. \quad (25)$$

With the polarizers removed, we specify the conditional probability of a count, given that particles enter the A and B lenses to be a constant independent of λ . Thus we specify

$$p_A(\lambda, \infty) = \alpha_A, p_B(\lambda, \infty) = \alpha_B; \text{ with } 0 \leq \alpha_{A,B} \leq 1. \quad (26)$$

With the polarizers in place, the conditional probability of a count at detector B, given that a particle enters the B lens system, is prescribed to be

$$p_B(\lambda, b) = \alpha_B \left[\frac{1}{2} \epsilon_+^B + \frac{1}{2} \epsilon_-^B \cos 2(\lambda - b) \right]. \quad (27)$$

The corresponding probability of a count at A is more pathological, thus⁹

$$p_A(\lambda, a) = \begin{cases} \frac{\alpha_A \pi \epsilon_+^A}{2\delta} & \text{for } a - \frac{1}{2}\delta + N\pi \leq \lambda \leq a + \frac{1}{2}\delta + N\pi, \\ 0 & \text{otherwise.} \end{cases} \quad (28)$$

where N is any integer, and δ depends upon the ratio $\epsilon_-^A / \epsilon_+^A$ and is defined by the equation

$$\frac{\sin \delta}{\delta} = \frac{\epsilon_-^A}{\epsilon_+^A}. \quad (29)$$

These prescriptions clearly satisfy the definitions of a LHVT.

When both polarizers are in place, the coincidence rate is equal to the emission rate of particle pairs, multiplied by the probability that both detectors will count, given that a pair was emitted into the lens system. Thus we calculate

$$R(a, b) \equiv R_e W(a, b) = \frac{R_e}{\pi} \int_{a - \frac{1}{2}\delta}^{a + \frac{1}{2}\delta} \frac{\alpha_A \pi \epsilon_+^A}{2\delta} \alpha_B \left[\frac{1}{2} \epsilon_+^B + \frac{1}{2} \epsilon_-^B \cos 2(\lambda - b) \right] d\lambda$$

$$\begin{aligned}
 &= R_{e_A \alpha_A \alpha_B} \left[\frac{1}{4} \epsilon_+^A \epsilon_+^B + \frac{1}{4} \epsilon_+^A \epsilon_-^B \frac{\sin \delta}{\delta} \cos 2(a-b) \right], \\
 &= R_{e_A \alpha_A \alpha_B} \frac{1}{4} \left[\epsilon_+^A \epsilon_+^B + \epsilon_-^A \epsilon_-^B \cos 2(a-b) \right]. \quad (30)
 \end{aligned}$$

The final expression here results from an application of Eq. (29). In similar manner the coincidence rates with one or both polarizers removed maybe calculated from Eq. (26)-(28). We find

$$R_1 = \frac{1}{2} R_{e_A \alpha_A \alpha_B} \epsilon_+^A, \quad (31)$$

$$R_2 = \frac{1}{2} R_{e_A \alpha_A \alpha_B} \epsilon_+^B, \quad (32)$$

$$R_0 = R_{e_A \alpha_A \alpha_B}. \quad (33)$$

Dividing Eq. (30)-(32) by (33) we obtain exactly the desired predictions given by Eqs.(22)-(24).

Finally, we must check that for the values $\epsilon_+ \approx 1.00$ and $\epsilon_-/\epsilon_+ \approx 0.94$ the resulting enhancement does not make $p_A(\lambda, a)$, via Eq. (28), exceed the sensible probability limit $p_A \leq 1$. This will be true as long as we have

$$\alpha_A \leq 2\delta / (\pi \epsilon_+) \leq 0.38. \quad (34)$$

The detector efficiencies in the actual experiment (excluding the loss due to solid angle limitation) were $\alpha_A \lesssim 0.026$ leaving (34) satisfied by a rather wide margin.

APPLICATION TO SEMICLASSICAL RADIATION THEORIES

The various forms of semiclassical radiation theory currently under discussion¹⁰ violate both assumptions A and D. Thus, an earlier discussion by Clauser¹¹ employed other plausible assumptions to show an incompatibility between these theories and experiment. Such theories are clearly OLT's. If they are to be consistent with Maxwell's equations, they must also satisfy the no-enhancement assumption. Hence the assumptions employed there are seen to be unnecessary. These theories are thus explicitly ruled out by the experimental results of Ref. 3.

It is discomfoting, however, to find existing theories that violate assumption D in agreement with much experimental data. After all, the no-enhancement assumption is simply a weaker form of assumption D. Without resorting to the use of a specific model, it is evidently impossible to test this assumption directly. The violation of D by these theories may be associated with their simultaneous violation of assumption A. Thus additional tests of A and D will lend reassurance as to the validity of the no-enhancement assumption. Such tests have been performed at this laboratory and will be reported elsewhere.

CONCLUSION

We see that any objective local theory (and hence any local hidden-variable theory) which employs the no-enhancement assumption is in direct violation of experimental results. Semiclassical radiation theories are to be included in this class. In view of the difficulty of performing fully conclusive experiments which do not require supplementary assumptions (e.g., the no-enhancement assumption), these results appear to be the best obtainable at the present time. However, as quantum mechanical correlation experiments improve, the strength of required additional assumptions should reasonably diminish.

FOOTNOTE AND REFERENCES

* Work supported by U. S. Atomic Energy Commission.

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² J. F. Clauser, M. A. Horne, A. Shimony, R. A. Holt, *Phys. Rev. Letters* 23, 880 (1969); J. F. Clauser, *Bull. Am. Phys. Soc.* 14, 578 (1969); A. Shimony, in Foundations of Quantum Mechanics, p.182, (see Ref.1); M. A. Horne, Ph.D. Thesis (Boston University, 1970) (unpublished).

³ S. J. Freedman and J. F. Clauser, *Phys. Rev. Letters* 28, 938 (1972); S. J. Freedman, Ph.D. Thesis (University of California, Berkeley, 1972) Lawrence Berkeley Laboratory Report LBL 391 (unpublished).

⁴ J. F. Clauser, Part I of this two-part article. *Phys. Rev.*

⁵ The numbers include losses in the optics, filters, electron collection, and other effects, but not the solid-angle loss which is included in the functions p and q. Deadtime effects are neglected.

⁶ Assumption A has normally been stated together with B as a single assumption. See Ref.4 for further discussion.

⁷ The idea of readjusting the analyzers while the particles are in flight was originally due to Y. Aharonov and D. Bohm [Phys. Rev. 108, 1070 (1957)].

⁸ The experiment by L. Kasday, J. Ulman, and C. S. Wu, [Bull. Am. Phys. Soc. 15, 586 (1970)]; see also L. R. Kasday in Foundations of Quantum Mechanics, p.195 (see Ref.1)] is noteworthy in that the prepared quantum state (positronium annihilation quanta) met the requirements stated in part I, but the polarization analyzers and detectors did not.

⁹ Photons resulting from a $J = 0 \rightarrow 1$ transition (channel A) can perhaps be tagged by some additional variable to distinguish them from those produced in a $J = 1 \rightarrow 0$ transition (channel B).

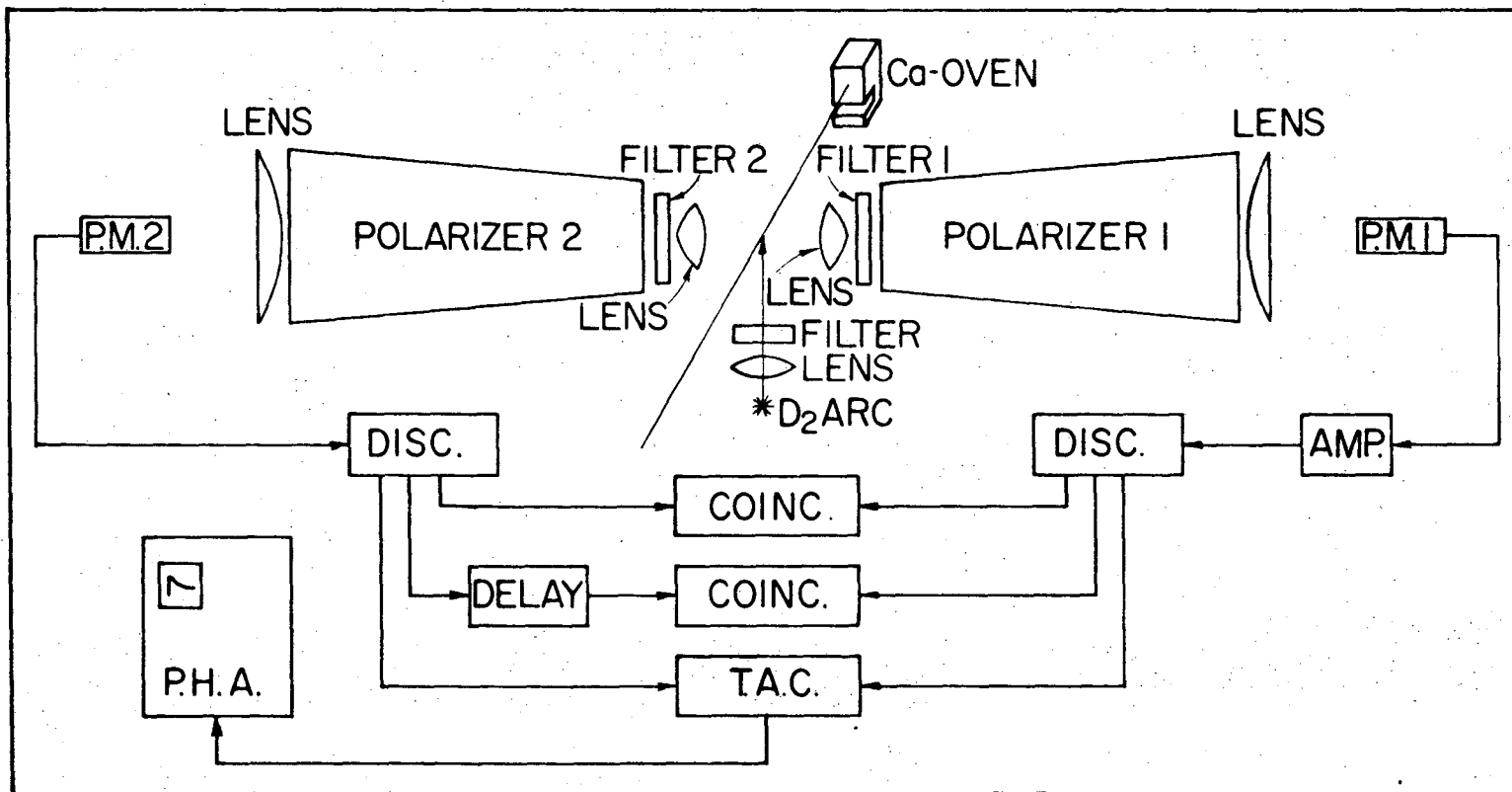
¹⁰ For a review of these see E. T. Jaynes in Coherence and Quantum Optics, Proceedings of the Third Rochester Conference on Coherence and Quantum Optics, edited by L. Mandel and E. Wolf, (Plenum Press, N.Y. 1973), p.35.

¹¹ J. F. Clauser, Phys. Rev. A 6, 49 (1972).

FIGURE CAPTIONS

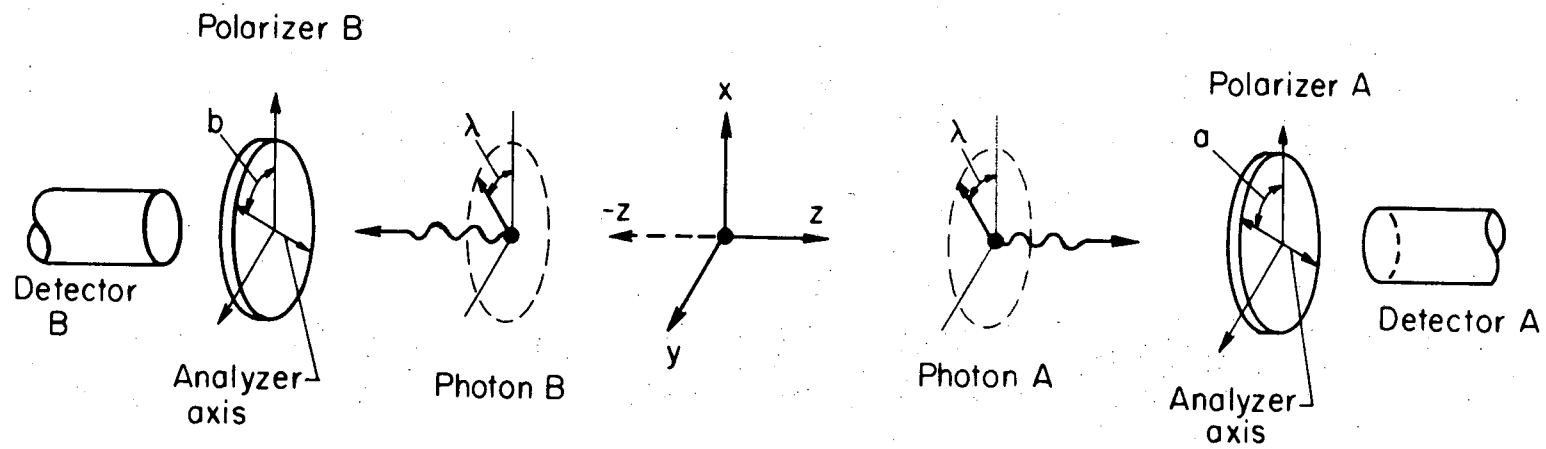
Fig.1. Schematic diagram of apparatus and associated electronics of experiment by Freedman and Clauser. Scalers (not shown) monitored the outputs of the discriminators and coincidence circuits.

Fig.2. Coordinate system for hidden-variable counter example. Photon particles A and B carry the same azimuthal direction λ , which, along with the analyzer orientation a or b , determines their probability for passage through the associated polarizer. The probability of a detector response also depends upon whether the photon has passed through the polarizer.



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Fig. 1



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Fig. 2

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