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Quantum Optical Tests of the Foundations of Physics

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77.1 INTRODUCTION: THE PHOTON HYPOTHESIS

Quantum mechanics began with the solution of the problem of blackbody radiation by Planck's quantum hypothesis: in the interaction of light with matter, energy can only be exchanged between the light in a cavity and the atoms in the walls of the cavity by the discrete amount $E = h\nu$, where h is Planck's constant and ν is the frequency of the light. Einstein, in his treatment of the photoelectric effect, reinterpreted this equation to mean that a beam of light consists of particles ("light quanta") with energy $h\nu$. The Compton effect supported this particle viewpoint of light by demonstrating that photons carried momentum, as well as energy. In this way, the wave-particle duality of quanta made its first appearance in connection with the properties of light.

It might seem that the introduction of the concept of the photon as a particle would necessarily also introduce the concept of locality into the quantum world. However, in view of observed violations of Bell's inequalities, exactly the opposite seems to be true. Here we review some recent results in quantum optics which elucidate nonlocality and other fundamental issues in physics.

In spite of the successes of quantum electrodynamics, and of the standard model in particle physics, there is still considerable resistance to the concept of the photon as a particle. Many papers have been written trying to explain all optical phenomena semiclassically, i.e., with the light viewed as a classical wave, and the atoms treated quantum mechanically [1–4]. We first present some quantum optics phenomena which exclude this semiclassical viewpoint.

In an early experiment, Taylor reduced the intensity of a thermal light source in Young's two-slit experiment, until, on the average, there was only a single photon passing through the two slits at a time. He then observed a two-slit interference pattern which was identical to that for a more intense classical beam of light. In Dirac's words, the apparent conclusion is that "each photon then interferes only with itself" [5]. However, a coherent state, no matter how strongly attenuated, always remains a coherent state (see Sect. 75.4); since a thermal light source can be modeled as a statistical ensemble of coherent states, a stochastic classical wave model yields complete agreement with Taylor's observations. The one-by-one darkening of grains of film can be explained by treating the matter alone quantum mechanically [2]; consequently, the concept of the photon need not be invoked, and the claim that this experiment demonstrates quantum interference of individual photons is unwarranted [6].

This weakness in Taylor's experiment can be removed by the use of nonclassical light sources: as discussed by Glauber [7], classical predictions diverge from quantum ones only when one considers counting statistics, or photon *correlations*. In particular, two-photon light sources, combined with coincidence detection, allow the production of single-photon ($n = 1$ Fock) states with near certainty. In the first such experiment [8], two photons, produced in an atomic cascade within nanoseconds of each other, impinged on two beam splitters, and were then detected in coincidence by means of four photomultipliers placed at all possible exit ports.

In a simplified version of this experiment [9], one of the beam splitters and its two detectors are replaced with a single detector D_1 (see Fig. 77.1). We define the anticorrelation parameter

$$\alpha \equiv N_{123}N_1/N_{12}N_{13}\,, \tag{77.1}$$

where N_{123} is the rate of triple-coincidences between detectors D_1, D_2 and D_3; N_1 is the singles rate at D_1; and N_{12} and N_{13} are double-coincidence rates. Then from Schwarz's inequality [9,6,10], $\alpha \geq 1$ for any classical wave. In essence, since the wave divides smoothly, the coincidence rate between D_2 and D_3 is never smaller than the "accidental" coincidence rate, even when measurements are conditioned on an event at D_1. (The Hanbury-Brown and Twiss experiment [11] *can* be explained classically, because the thermal fluctuations lead to "bunching," or a mean coincidence rate which is *greater* than the mean accidental rate; *cf.* Sect.75.4.3.) By contrast, the indivisibility of the photon leads to strong anticorrelations between D_2 and D_3, making α arbitrarily small. In agreement with this quantum mechanical picture, Grangier *et al.* observed a 13-standard-deviation violation of the inequality [9], corroborating the notion of the "collapse of the wave packet" as proposed by Heisenberg [12] (see also Sect. 73.5.2).

77.2 QUANTUM PROPERTIES OF LIGHT

77.2.1 Vacuum Fluctuations: Cavity QED

The above considerations necessitate the quantization of the electromagnetic field, which in turn leads to the concept of vacuum fluctuations [4] (see Sect. 75.2). Difficulties with this idea, such as the implied infinite zero-point energy of the universe, have led some reseachers to attempt to dispense with this concept altogether, along with that of the photon, in every explanation of electromagnetic interactions with matter. Of course, it is impossible to explain all phenomena, such as spontaneous emission and the Lamb shift, without some kind of fluctuating electromagnetic fields (see Chap. 75 and Sect. 76.3.3), but one can go a long way with an *ad hoc* ambient classical electromagnetic noise-field filling all of space, in conjunction with the radiation reaction [1,4].

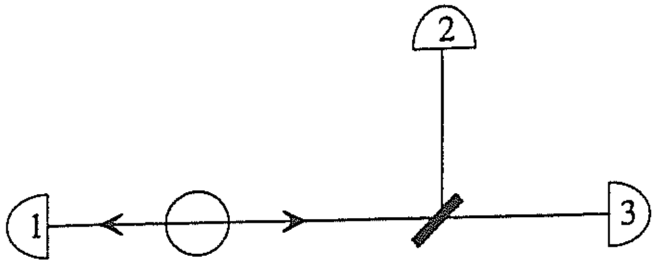

Figure 77.1. Triple-coincidence setup of Grangier *et al.* [9].

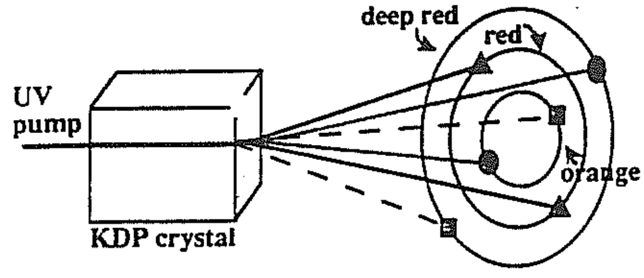

Figure 77.2. Conical emissions of down-conversion from a nonlinear crystal. Photon energy depends on the cone opening angle, and conjugate photons lie on opposite sides of the axis, e.g.. the inner "circle" orange photon is conjugate to the outer "circle" deep-red photon, etc.

In particular, even the Casimir attraction between two conducting plates (see Sect. 76.1) can be explained semiclassically in terms of dipole forces between electrons in each plate as they undergo zero-point motion and induce image charges in the other plate (see Sect. 76.3.3). Nevertheless, the effects of cavity QED (Chap. 76) [13–15], including the influence of cavity-induced boundary conditions on energy levels and spontaneous emission rates, are most easily unified via quantization of the electromagnetic field.

77.2.2 The Down-Conversion Two-Photon Light Source

The quantum aspects of electromagnetism are made more striking with a two-photon light source, in which two highly correlated photons are produced in spontaneous parametric down-conversion, or parametric fluorescence [16–18]. In this process, an ultraviolet "pump" photon produced in a laser spontaneously decays inside a crystal with a $\chi^{(2)}$ nonlinearity into two highly correlated red photons, conventionally called the "signal" and the "idler" (see Sect. 70.4.4). As shown in Fig. 77.2, a rainbow of colored cones is produced around an axis defined by the direction of the uv beam (for the case of type-I phase-matching), with the correlated down-conversion photons always emitted on opposite sides of the uv beam (see Sect. 70.3.2). Their emission times are within femtoseconds of each other, so that detection of one photon implies with near certainty that there is exactly one quantum present in the conjugate mode.

Energy and momentum are conserved in this process:

$$\hbar\omega_0 = \hbar\omega_1 + \hbar\omega_2 , \tag{77.2}$$

$$\hbar \mathbf{k}_0 \approx \hbar \mathbf{k}_1 + \hbar \mathbf{k}_2 , \tag{77.3}$$

where $\hbar\omega_0$ ($\hbar\mathbf{k}_0$) is the energy (momentum) of the parent photon, and $\hbar\omega_1$ ($\hbar\mathbf{k}_1$) and $\hbar\omega_2$ ($\hbar\mathbf{k}_2$) are the energies (momenta) of the daughter photons: $\mathbf{k}_1$ and $\mathbf{k}_2$ sum to $\mathbf{k}_0$ to within an uncertainty given by the reciprocal of the crystal length [19]. Since there are many ways of partitioning the parent photon's energy, each daughter photon may have a broad spectrum. and hence a wave packet narrow in time. However, $\omega_1 + \omega_2 = \omega_0$ is extremely well-defined, so that the *difference* in the daughter photons' arrival times, and the *sum* of their energies can be simultaneously known to high precision.

In type-I phase matching the correlated photons share the same polarization, while in type-II phase matching they have orthogonal polarizations. This latter possibility has recently permitted the production of *unpolarized* down-conversion photons that nevertheless display near-perfect polarization correlations [20]; see Sect. 77.5.2.

77.2.3 Squeezed States of Light

The creation of correlated photon pairs is closely related to the process of quadrature-squeezed light production (see Sect. 75.3, and the review in Ref. [21]). For example. when the gain arising from parametric amplification in a down-conversion crystal becomes large, there is a transition from spontaneous to stimulated emission of pairs. This gain is dependent on the phase of amplified light relative to the phase of the pump light. As a result, the vacuum fluctuations are reduced ("squeezed") below the standard quantum limit in one quadrature, but increased in the other, in such a way as to preserve the minimum uncertainty-principle product [22]. This periodicity of the fluctuations at 2ω is a direct consequence of the fact that the light is a superposition of states differing in energy by $2\hbar\omega$ — the quadrature-squeezed vacuum state

$$|\xi\rangle = \exp\left(\tfrac{1}{2}\xi^* aa - \tfrac{1}{2}\xi a^\dagger a^\dagger\right)|0\rangle \tag{77.4}$$

represents a vacuum state transformed by the creation ($a^\dagger a^\dagger$) and destruction (aa) of photons two at a time.

Essentially any optical processes operating on photon pairs (e.g., four-wave mixing [23, 24]) can also produce such squeezing.

Amplitude squeezing involves preparation of states with well-defined photon number, i.e., states lacking the Poisson fluctuations of the coherent state. The possibility of producing such states (e.g., via a constant-current-driven semiconductor laser [25]) demonstrates that "shot noise" in photodetection should not be thought of as merely the result of the probabilistic (à la Fermi's Golden Rule, *cf.* Sect. 67.5) excitation of quantum mechanical atoms in a classical field, but as representing real properties of the electromagnetic field, accessible to experimental control.

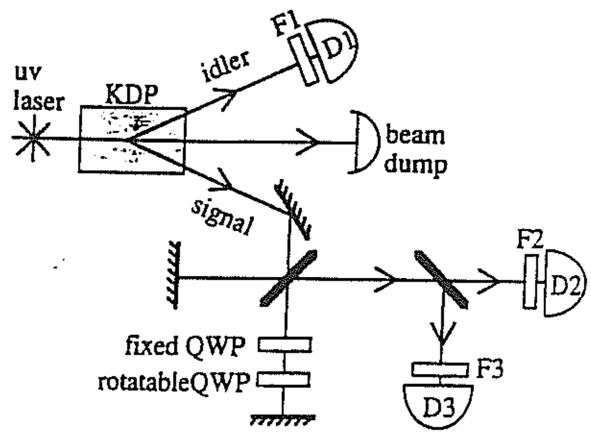

Figure 77.3. The energy-time uncertainty relation and wave function collapse were studied by investigating the effect of various filters before the detectors in a single-photon interference experiment [9, 35].

77.2.4 Quantum Nondemolition

The uncertainty principle between the number of quanta N and phase ϕ of a beam of light,

$$\Delta N \Delta \phi \geq 1/2 , \qquad (77.5)$$

implies that to know the number of photons exactly, one must give up all knowledge of the phase of the wave. In theory, a *quantum nondemolition* (QND) process is possible [26]: without annihilating any of the light quanta, one can count them. It might seem that this would make possible successive measurements on noncommuting observables of a single photon, in violation of the uncertainty principle; it is the unavoidable introduction of phase uncertainty by any number measurement which prevents this.

The most common QND schemes [27] use the intensity-dependent index of refraction arising from the optical Kerr effect (see Sect. 70.5.3) — the change in the index due to the intensity of the "signal" beam changes the optical phase shift on a "probe" beam [28–30]. Still other QND proposals use Rydberg atoms to give indirect information concerning the number of photons in a microwave cavity through which they have passed [31, 32] (see Sect. 76.5.1), or the Aharonov-Bohm effect to sense photons via the phase shift their fields induce in passing electrons [33, 34].

77.3 NONCLASSICAL INTERFERENCE

77.3.1 Single-Photon Interference and Energy-Time Uncertainty

The simplest one-photon interference experiment [9] uses the cascade source discussed in Sect. 77.1. One of the photons is directed to a "trigger" detector, while the other, thus prepared in an $n = 1$ Fock state, is sent through a Mach-Zehnder interferometer. The output photon, detected in coincidence with the trigger photon, showed fringes with a visibility > 98%. Dirac's statement that a single photon interferes with itself is thus verified.

The energy-time uncertainty principle,

$$\Delta E \Delta t \geq \hbar/2 , \qquad (77.6)$$

has recently been tested in a related interference experiment. While the canonical commutator $[x, p_x] = i\hbar$ gives rise to $\Delta x \Delta p_x \geq \hbar/2$, the corresponding commutator $[t, \mathcal{H}] = i\hbar$ does not exist. If the t operator existed, it would be the generator of translations of either sign in energy (just as $\mathcal{H}$ is the generator of translations in x), and could thus generate arbitrarily large negative energies — but unlike momentum, energy is bounded from below[1]. Thus, in the time-dependent Schrödinger equation, time is a c-number parameter [36] (as is position as well in relativistic theory), rather than an observable.

Due to energy conservation in the down-conversion process, the daughter photons of a parent photon of sharp energy E_0 are in an energy-"entangled" state, a nonfactorizable sum of product states [37]:

$$|\Psi\rangle = \int_0^{E_0} dE\, A(E) |E\rangle |E_0 - E\rangle , \qquad (77.7)$$

where $A(E)$ is the probability amplitude for the production of two photons of energies E and $E_0 - E$. A measurement of the energy of one of the photons to be E_1 causes an instantaneous "collapse" of the system to the state $|E_1\rangle|E_0 - E_1\rangle$, implying an instantaneous increase of the width of the other photon's wave packet.

[1] As can be seen from the fact that $\mathcal{H} = (N + 1/2)\hbar\omega$ and $\phi = \omega t$, these problems are closely related to the number-phase uncertainty relation and the difficulties of defining a phase operator (Sect. 75.3.4).

In a modification [35] of the previous experiment (see Fig. 77.3), an interference filter F1 is used to sharply measure the frequency of the trigger photon 1, while a Michelson interferometer is used to measure the coherence length of the conjugate photon 2. If the trigger photon passes through the filter of narrow width ΔE and is detected, then photon 2 occupies a broad wave packet of duration $\Delta t \approx \hbar/\Delta E$, whereas the absence of fringes without triggering implies a much shorter wave packet. This is a nonlocal effect in that the photons can be arbitrarily far away from each other when the collapse occurs. (The same apparatus was also used to demonstrate that Berry's phase in optics has a quantum origin [38].)

77.3.2 Two-Photon Interference

In the above experiments, interference occurs between two paths taken by a *single* photon. An early experiment to demonstrate *two*-photon interference using the down-conversion light source was performed by Ghosh and Mandel [39]. They looked at the counting rate of a detector illuminated by both of the twin beams. No interference was observed at the detector, because although the sum of the phases of the two beams emitted in parametric fluorescence is well-defined (by the phase of the pump), their *difference* is not, due to the number-phase uncertainty principle (see Sect. 75.4). However, when Ghosh and Mandel looked at the rate of *coincidence* detections between two such detectors whose separation was varied, they observed high-visibility interference fringes. Whereas in the standard two-slit experiment, interference occurs between the two paths a single photon could have taken to reach a given point on a screen, in this case it occurs between the possibility that the signal photon reached detector 1 and the idler photon detector 2, and the possibility that the reverse happened. This experiment provides a manifestation of quantum nonlocality; interference occurs between alternate global histories of a system, not between local fields. At a null of the coincidence fringes, the detection of one photon at detector 1 excludes the possibility of finding the conjugate photon at detector 2.

Such interference becomes clearer in the related interferometer of Hong, Ou, and Mandel [40] (see Fig. 77.4). The identically-polarized conjugate photons from a down-conversion crystal are directed to opposite sides of a 50-50 beam splitter, such that the transmitted and reflected modes overlap. If the difference in the path lengths ΔL prior to the beam splitter is larger than the two-photon correlation length (of the order of the coherence length of the down-converted light), the photons behave independently at the beam splitter, and coincidence counts between detectors in the two output ports are observed half of the time — the other half of the time both photons travel to the same detector. However, when $\Delta L \approx 0$, such that the photon wave packets overlap at the beam splitter, the probability of coincidences is reduced, in principle to zero if $\Delta L = 0$. One can explain the coincidence null at zero path-length difference using the Feynman rules for calculating probabilities: add the probability amplitudes of *indistinguishable* processes which lead to the same final outcome, and then take the absolute square. The two indistinguishable processes here are both photons being reflected at the beam splitter (with Feynman amplitude $\mathrm{r}\cdot\mathrm{r}$) and both photons being transmitted (with Feynman amplitude $\mathrm{t}\cdot\mathrm{t}$). The probability of a coincidence detection is then

$$P_c = |\mathrm{r}\cdot\mathrm{r}+\mathrm{t}\cdot\mathrm{t}|^2 = \left|\frac{i}{\sqrt{2}}\cdot\frac{i}{\sqrt{2}}+\frac{1}{\sqrt{2}}\cdot\frac{1}{\sqrt{2}}\right|^2 = 0 \ . \tag{77.8}$$

assuming a real transmission amplitude, and where the factors of i come from the phase shift upon reflection at a beam splitter [42, 43].

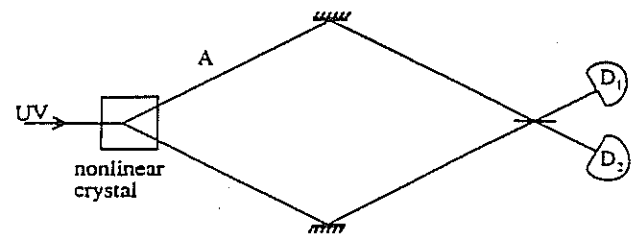

Figure 77.4. Simplified setup for a Hong-Ou-Mandel (HOM) interferometer [40]. Coincidences may result from both photons being reflected, or both being transmitted. When the path lengths to the beam splitter are equal, these processes destructively interfere, causing a null in the coincidence rate. In a modified scheme, a half waveplate in one arm of the interferometer (at "A") serves to distinguish these otherwise interfering processes, so that no null in coincidences is observed. Using polarizers before the detectors, one can "erase" the distinguishability, thereby restoring interference [41]. (See Sects. 77.4.1 and 77.5.2.)

The possibility of a perfect null at the center of the dip is indicative of a nonclassical effect. Indeed, classical field predictions allow a maximum coincidence-fringe visibility of only 50% [44]. The tendency of the photons to travel off together at the beam splitter can be thought of as a manifestation of the Bose-Einstein statistics for the photons [45]. In practice, the bandwidth of the photons, and hence the width of the null, is determined by filters and/or irises before the detectors [19]. Widths as small as 5 μm have been observed, corresponding to time delays of only 15 fs [46]. Consequently, one application is the determination of single-photon propagation times with extremely high time resolution (see Sect. 77.7).

77.4 COMPLEMENTARITY

77.4.1 Quantum Eraser

The complementary nature of wave-like and particle-like behavior is frequently interpreted as follows: due to the uncertainty principle, any attempt to measure the position (particle aspect) of a quantum leads to an uncontrollable, irreversible disturbance in its momentum, thereby washing out any interference pattern (wave aspect) [47, 48]. This picture is incomplete though; no "state reduction," or "collapse," is necessary to destroy interference, and measurements which do not involve reduction can be reversible. One must view the loss of coherence as arising from an *entanglement* of the system wave function with that of the measuring apparatus (MA) [49]. Previously interfering paths can thereby become distinguishable, such that no interference is observed. Interference may be regained, however, if one manages to "erase" the distinguishing information. This is the physical content of quantum erasure [50, 51]. The primary lesson is that one must consider the total physical state, including any MA with which the interfering quantum has become entangled. If the coherence of the MA is maintained, then interference may be recovered.

One demonstration of a quantum eraser is based on the interferometer in Fig. 77.4 [41]. A half waveplate inserted into one of the paths before the beam splitter serves to rotate the polarization of light in that path. In the extreme case, the polarization is made orthogonal to that in the other arm, and the r·r and t·t processes become distinguishable: hence, the destructive interference which led to a coincidence null does not occur. The distinguishability can be erased, however, by using polarizers just before the detectors. In particular, if the initial polarization of the photons is horizontal, and the waveplate rotates one of the photon polarizations to vertical, then polarizers at 45° before both detectors restore the original interference dip. If one polarizer is at 45° and the other at −45°, interference is once again seen, but now in the form of a peak instead of a dip (see Fig. 77.5). There are four basic measurements possible on the MA (here the polarization) — two of which yield which-path information, one of which recovers the initial interference fringes (here the coincidence dip), and one of which yields interference anti-fringes (the peak instead of the dip). In some implementations, the decision to measure wave-like or particle-like behavior may even be delayed until *after detection* of the original quantum, an irreversible process[2] [55, 56]. But in all cases, one must correlate the

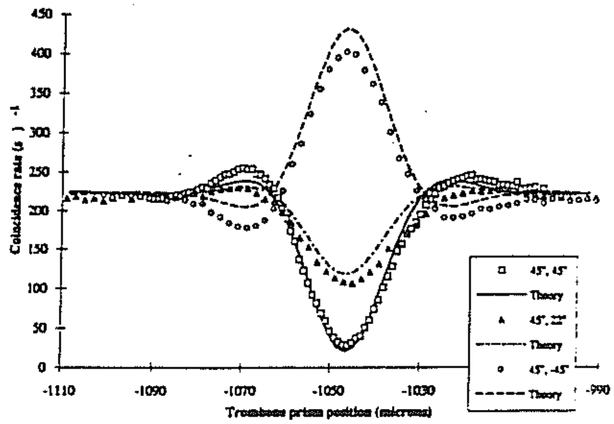

Figure 77.5. Experimental data and scaled theoretical curves (adjusted to fit observed visibility of 91%) with polarizer 1 at 45° and polarizer 2 at various angles. Far from the dip, there is no interference and the angle is irrelevant [41].

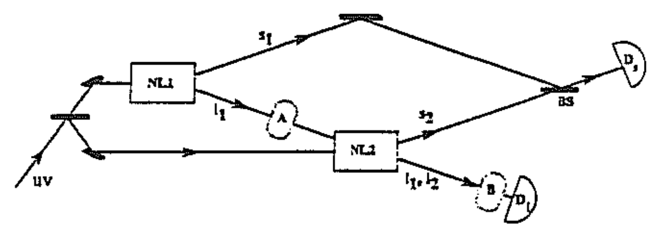

Figure 77.6. Schematic of setup used in [58]. The idler photons from the two crystals are indistinguishable; consequently, interference fringes may be observed in the signal singles rate at detector D_s. Additional elements at A and B can be used to make a quantum eraser.

results of measurements on the MA with the detection of the originally interfering system.

77.4.2 Vacuum-Induced Coherence

A somewhat different demonstration [57, 58] of complementarity involves two down-conversion crystals, NL1 and NL2, aligned such that the trajectories of the idler photons from each crystal overlap (see Fig. 77.6). A beam splitter acts to mix the signal modes. If the path lengths are adjusted correctly, and the idler beams overlap precisely, there is no way to tell, even in principle, from which crystal a photon detected at D_s originated. Interference appears in the signal *singles* rate at D_s, as any of the path lengths is varied. If the idler beam from crystal NL1 is prevented from entering crystal NL2, then the interference vanishes, because the presence or absence of

[2]This is an extension of the original delayed-choice discussion by Wheeler [52], and the experiments by Hellmuth *et al.* and Alley *et al.* [53, 54], in which the decision to display wave-like or particle-like aspects in a light beam may be delayed until after the beam has been split by the appropriate optics.

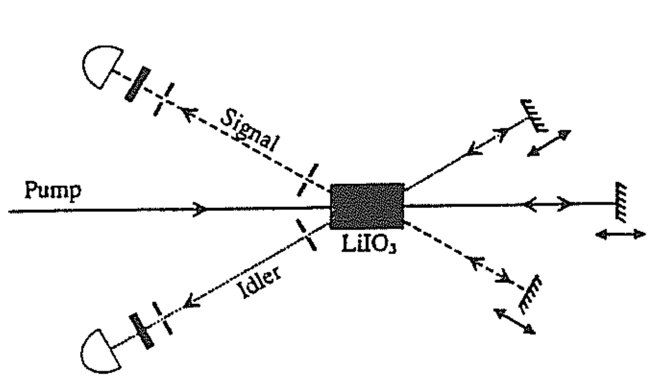

Figure 77.7. Schematic of the experiment to demonstrate enhancement and suppression of spontaneous down-conversion [59].

an idler photon at D_i then "labels" the parent crystal. Thus one explanation for the effect of blocking this path is that coherence is established by the idler-mode *vacuum* field seen by both crystals.

77.4.3 Suppression of Spontaneous Parametric Down-Conversion

A modification [59] of this two-crystal experiment uses only a single nonlinear crystal (see Fig. 77.7). A given pump photon may down-convert in its initial rightward passage through the crystal, or in its left-going return trip (or not at all, the most likely outcome). As in the previous experiment, the idler modes from these two processes are made to overlap; moreover, the signal modes are also aligned to overlap. Thus, the left-going and right-going production processes are indistinguishable and interfere. The result is that fringes are observed in all of the counting rates (i.e., the coincidence rate and both singles rates) as any of the mirrors is translated. One interpretation is as a change in the spontaneous emission of the down-converted photons. In contrast to the cavity QED demonstrations discussed in Sects. 77.2.1 and 76.1, here the distances to the mirrors are much longer than the coherence lengths of the spontaneously emitted photons.

With the inclusion of waveplates to label the photons' paths, and polarizers to erase this information, an improved quantum eraser experiment was recently completed [56]. The arrangement has also shown that in a sense, there are always photons between the down-conversion crystal and the mirrors, even in the case of complete suppression of the spontaneous emission process [60].

77.5 THE EPR PARADOX AND BELL'S INEQUALITIES

77.5.1 Generalities

Nowhere is the nonlocal character of the quantum mechanical entangled state as evident as in the "paradox" of Einstein, Podolsky, and Rosen (EPR) [61], the version of Bohm [62], and the related inequalities by Bell [63, 64]. Consider two photons traveling off back-to-back, described by the entangled singlet-like state

$$|\psi\rangle = (|H_1, V_2\rangle - |V_1, H_2\rangle)/\sqrt{2}, \tag{77.9}$$

where the letters denote horizontal (H) or vertical (V) polarization, and the subscripts denote photon propagation direction. This state is isotropic — it has the same form regardless of what basis is used to describe it. Measurement of any polarization component for one of the particles will yield a count with 50% probability; individually, each particle is unpolarized. Nevertheless, if one measures the polarization component of particle 1 in any basis, one can predict with certainty the polarization of particle 2 in the same basis, seemingly without disturbing it, since it may be arbitrarily remote. Therefore, according to EPR, to avoid any nonlocal influences one should ascribe an "element of reality" to every component of polarization. A quantum mechanical state cannot specify that much information, and is consequently an incomplete description, according to the EPR argument. The intuitive explanation implied by EPR is that the particles leave the source with definite, correlated properties, determined by some local "hidden variables" not present in quantum mechanics (QM).

For two entangled particles, a local hidden variable (LHV) theory can be made which correctly describes perfect correlations or anti-correlations (i.e., measurements made in the same polarization basis). The choice of an LHV theory versus QM is then a philosophical decision, not a physical one. In 1964 John Bell [63] discovered that QM gives different statistical predictions than does *any* LHV theory, for situations of nonperfect correlations (i.e., analyzers at intermediate angles). His inequality constraining various joint probabilities given by any local realistic theory was later generalized to include any model incorporating locality [65, 66], and also extended to apply to real experimental situations [67, 68]; the Clauser-Horne (CH) form relates the directly observable singles rates S_1 and S_2 and the coincidence rate C_{12}, rather than "inferred" probabilities, by

$$C_{12}(a, b) + C_{12}(a, b') + C_{12}(a', b) - C_{12}(a', b') \leq S_1(a) + S_2(b), \tag{77.10}$$

where a and a' are any pair of analyzer (e.g., polarizer) settings at detector 1, and b and b' a pair of settings at detector 2. For certain choices of a, a', b, and b', QM predicts a violation of this inequality.

In fact, inequality (77.10) has not yet been violated in any experiment because the minimum detector efficiency required is 67% [69]. Experiments thus far have employed an additional assumption, roughly equivalent to the "fair-sampling" assumption that the fraction of pairs detected is representative of the entire ensemble (see [70] for a fuller discussion). Moreover, no experiment has fulfilled the additional requirement that the analyzer settings be switched randomly and rapidly. Two efforts are now underway to attempt a direct violation of Eq. (77.10) with no auxiliary assumptions. One utilizes down-converted photons [71] and newly developed high-efficiency (> 85%) single-photon detectors [72, 73]; the other uses dissociated mercury dimers as the correlated particles, the advantage being detection efficiencies of 95% by photoionizing the atoms and detecting the photoelectrons [74].

77.5.2 Polarization-Based Tests

With the caveat of supplementary assumptions, Bell's inequalities have now been tested many times, and the vast majority of experiments have violated these inequalities, in support of QM. The first such test was performed with pairs of photons produced via an atomic cascade, and a later version incorporated rapid (albeit periodic) switching of the analyzers [75, 76]. Unfortunately, the angular correlation of the cascade photons is not very strong. In contrast, the strong correlations of the down-converted photons make them ideal for such tests, the first of which were performed using setups essentially identical to that already discussed in connection with quantum erasure [77, 78], (see Fig. 77.4). With the condition of equal path lengths, and with the half wave-plate oriented to rotate the polarization from horizontal to vertical, the state of light after the 50-50 beam splitter is

$$|\psi\rangle \approx \tfrac{1}{2}\left[|H_1V_2\rangle - |V_1H_2\rangle + i|V_1H_1\rangle + i|V_2H_2\rangle\right] , \qquad (77.11)$$

which reduces to Eq. (77.9) if one considers only the terms which can lead to coincidence detections (*viz.*, the first two). The probability of coincidence detection depends nonlocally on the difference in the orientation angles of the two polarizers: $P_c \propto \sin^2(\theta_2 - \theta_1)$. A more recent experiment [79] using the orthogonally-polarized photons from collinear type-II down-conversion (see Sect. 70.3.2) has produced a violation of a Bell's inequality by 22 standard deviations. Yet more recently, a *non*-collinear type-II down-conversion experiment has produced true polarization-entangled states (i.e., without the need to discard half of the terms) and produced a 100 standard deviation violation of a Bell's inequality.

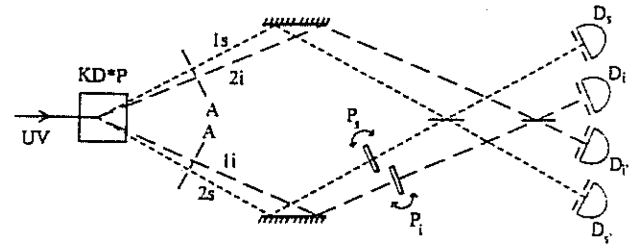

Figure 77.8. Outline of apparatus used to demonstrate a violation of a Bell's inequality based on momentum entanglement [80].

77.5.3 Nonpolarization Tests

The advent of parametric down-conversion has also led to the appearance of several nonpolarization-based tests of Bell's inequalities, using, for example, an entanglement of the photon momenta (see Fig. 77.8) [80]. By use of small irises (labeled 'A' in the figure), only four down-conversion modes are examined: 1s, 1i, 2s, and 2i. Beams 1s and 1i correspond to one pair of conjugate photons; beams 2s and 2i correspond to a different pair. Photons in beams 1s and 2s have the same wavelength, as do photons in beams 1i and 2i. With proper alignment, after the beam splitters there is no way to tell whether a pair of photons came from the 1s-1i or the 2s-2i paths. Consequently, the coincidence rates display interference, although the singles rates at the four detectors indicated in Fig. 77.8 remain constant. This interference depends on the difference of phase shifts induced by rotatable glass plates P_i and P_s in paths 1i and 2s, respectively, and is formally equivalent to the polarization case considered above, in which it is the difference of polarization-analyzer angles that is relevant. By measuring the coincidence rates for two values for each of the phase shifters — a total of four combinations — the experimenters were able to violate an appropriate Bell's inequality. One interpretation is that the paths taken by a given pair of photons are not elements of reality.

Several groups [81, 82] have implemented a different Bell's inequality proposal [83], based on energy-time entanglement of the photons. One member of each down-converted pair is directed into an unbalanced Mach-Zehnder-like interferometer, allowing both a short and long path to the final beam splitter; the other photon is directed into a separate but similar interferometer. There arises interference between the indistinguishable processes ("short-short" and "long-long") which could lead to coincidence detection. Using fast detectors to select out only these processes, the reduced state [*cf.* Eq. (77.9)] is

$$|\psi\rangle = \frac{1}{2}\left(|S_1, S_2\rangle - e^{i\phi}|L_1, L_2\rangle\right) . \qquad (77.12)$$

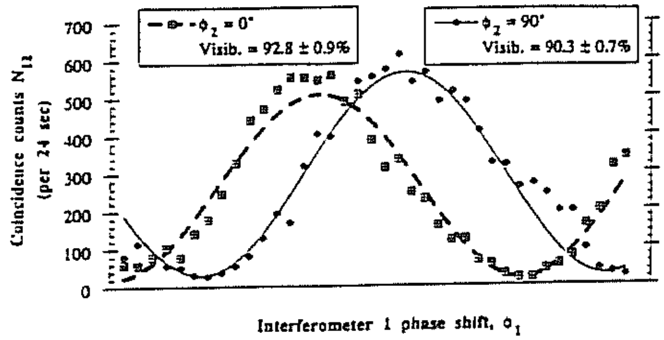

Figure 77.9. High-visibility coincidence fringes in a Franson dual-interferometer experiment [83] for two values of the phase in interferometer 2 as the phase in interferometer 1 is slowly varied. The curves are sinusoidal fits.

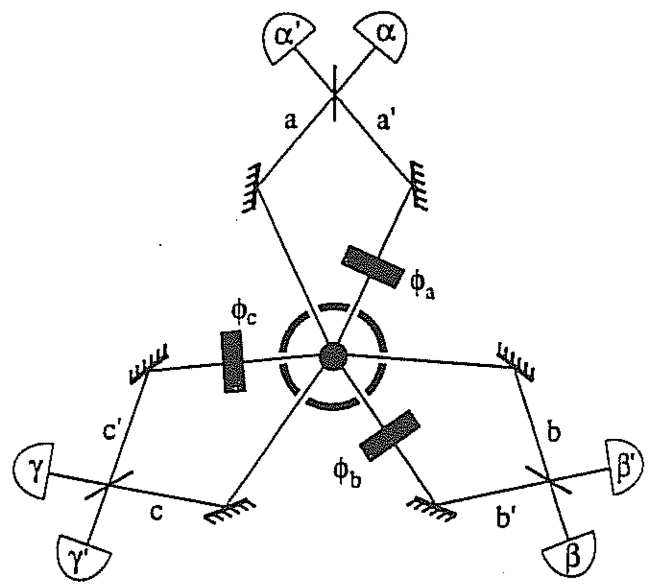

Figure 77.10. A three-particle gedanken experiment to demonstrate the inconsistency of quantum mechanics and any local realistic theory. All beam splitters are 50-50 [57].

where the letters indicate the short or long path, and the phase is the sum of the relative phases in each interferometer. Although no fringes are seen in any of the singles count rates, the high-visibility coincidence fringes (Fig. 77.9) lead to a violation of an appropriate Bell's inequality. One conclusion is that it is incorrect to ascribe to the photons a definite time of emission from the crystal, or even a definite energy, unless these observables are explicitly measured. A more general interpretation, applicable to all violations of Bell's inequalities, is that the predictions of QM cannot be reproduced by any completely local theory. It must be that the results of measurements on one of the photons depend on the results for the other, and these correlations are not merely due to a common cause at their creation [84, 85].

77.6 RELATED ISSUES

77.6.1 Nonlocality Without Inequalities

In the above experiments for testing nonlocality, the disagreement between quantum predictions and Bell's constraints on local realistic theories are only statistical. Recently, Greenberger, Horne, and Zeilinger (GHZ) pointed out that in some systems involving *three* or more entangled particles [86, 87], a contradiction could arise even at the level of perfect correlations. A schematic of one version of the GHZ *Gedankenexperiment* is shown in Fig. 77.10. The source at the center is posited to emit trios of correlated particles. Just as the Rarity-Tapster experiment selected two pairs of photons (*cf.* Fig. 77.8), the GHZ source selects two trios of photons: these are denoted by abc and $a'b'c'$. Hence, the state coming from the source may be written

$$|\psi\rangle = (|abc\rangle + |a'b'c'\rangle) / \sqrt{2} . \qquad (77.13)$$

After passing through a variable phase shifter (e.g., ϕ_a), each primed beam is recombined with the corresponding unprimed beam at a 50-50 beam splitter. Detectors (denoted by Greek letters) at the output ports signal the occurrence of triple coincidences. The following simplified argument conveys the spirit of the GHZ result.

Given the state (77.13), one can calculate from standard QM the probability of a triple coincidence as a function of the three phase shifts:

$$P(\phi_1, \phi_2, \phi_3) = \tfrac{1}{8}\left[1 \pm \sin(\phi_a + \phi_b + \phi_c)\right] , \qquad (77.14)$$

where the plus sign applies for coincidences between all unprimed detectors, and the minus sign for coincidences between all primed detectors. For the case in which all phases are 0, it will occasionally happen (1/8th of the time) that there will be a triple coincidence of all primed detectors. Using a "contrafactual" approach, we ask what would have happened if ϕ_a had been $\pi/2$ instead. By the locality assumption, this would not change the state from the source, nor the fact that detectors β' and γ' went off. But from (77.14) the probability of a triple coincidence for primed detectors is zero in this case; therefore, we can conclude that detector α would have "clicked" if ϕ_a had been $\pi/2$. Similarly, if ϕ_b or ϕ_c had been $\pi/2$, then detectors β or γ would have clicked. Consequently, if all the phases had been equal to $\pi/2$, we would have seen a triple coincidence between unprimed detectors. But according to Eq. (77.14) this is impossible: the probability of triple coincidences between unprimed detectors when all three phases are equal to $\pi/2$ is strictly zero! Hence, if one believes the quantum mechanical predictions for these cases of perfect correlations, it is not possible to have a consistent local realistic model. By similar arguments, one can even demonstrate nonlocality without inequalities, at least in a *gedanken* experiment, using just two particles [88, 89]. When the arguments are suitably modified to deal with real experiments [90], inequalities once again result. Nevertheless, these new schemes underscore the vast disagreement between quantum theory and locality.

77.6.2 Information Content of a Quantum

The inherent nonlocality of particles in an entangled state cannot be used to transmit superluminal messages. If A and B receive an EPR pair of photons, which A then collapses in a certain basis by performing a polarization measurement. B can only extract one bit of information from a measurement on his photon — this bit corresponds not to A's choice of basis, but to the (random) *outcome* of A's measurement. However, instantaneous communication *would* be possible if one could make copies ("clones") of a single photon in an unknown polarization state: by performing measurements on n copies of his photon, B could determine its polarization to a resolution of n bits, thereby accurately determining A's choice of basis. Taking into account quantum fluctuations [91], one finds that no amplifier can make a sufficiently faithful copy for such a scheme to work — it is impossible to clone photons [92].

This makes the recent work by Bennett *et al.* on "quantum teleportation" and related effects all the more remarkable. For example, a single photon can be used to transfer two bits of information, when it is part of an entangled EPR pair [93]. Again consider A and B, each possessing one photon of such a pair. By manipulating only her photon (via a polarization rotator and a phase shifter), A can convert the joint state into any of the four two-particle "Bell states" $[|\psi^{\pm}_{AB}\rangle = (|\mathrm{H}_A, \mathrm{V}_B\rangle \pm |\mathrm{V}_A, \mathrm{H}_B\rangle)/\sqrt{2},\ |\phi^{\pm}_{AB}\rangle = (|\mathrm{H}_A, \mathrm{H}_B\rangle \pm |\mathrm{V}_A, \mathrm{V}_B\rangle)/\sqrt{2}]$, (such states are analogous to measuring J and m for a pair of spin-1/2 particles, rather than the individual spin projections) and then send her photon to B. By making a suitable measurement on both photons, B can then in principle determine which of the four states A produced [94]: A's single photon carried two bits.

In the even more striking *quantum teleportation* effect [95], an unknown polarization state f (with its in-principle infinite amount of information) can be "teleported" from B to A, if each already possesses one photon of an EPR pair (in the singlet-state ψ^{-}_{AB}). First, B jointly measures his EPR photon and the photon F (whose state f is to be teleported) in the basis defined by the Bell states of these two photons. Via a mere two bits of classical information, B then informs A which of the four (equally probable) Bell states he actually measured. With this information, A can transform the state of her EPR particle into f.

If B found the singlet-state ψ^{-}_{BF}, then the polarization of his EPR photon must have been orthogonal to that of F (because the polarizations of particles in a singlet-state are always perfectly anticorrelated, regardless of the quantization basis). But because the two EPR photons were initially also in a singlet-state, *their* polarizations must also be orthogonal, so A's EPR photon is already in state f. If instead B found ϕ^{-}_{BF}, for example, then A simply makes the same transformation (with a polarization rotator) that would have changed ψ^{-}_{AB} into ϕ^{-}_{AB}, again leaving her photon in the state f. Thus, although one may only extract one bit of (normally useless) information from an EPR particle, the perfect correlations may be used to transfer an infinite amount of information, i.e., precise specification of a point in the state space of the particle (the Poincaré sphere for a photon or the Bloch sphere for an electron). The "no cloning" theorem is not violated, since B irrevocably alters the state of F by the measurement he performs, leaving only one particle in f.

77.6.3 Cryptography

Although EPR schemes cannot send signals superluminally, they have other potential applications in cryptography. In the "one-time pad" of classical cryptography [96], two collaborators share a secret "key" (a random string of binary digits) in order to encode and decode a message. Such a key may provide an absolutely unbreakable code, provided that it is unknown to an eavesdropper. The problem arises in key distribution: any classical distribution scheme is subject to noninvasive eavesdropping, e.g., using a fiber-coupler to tap the line, without disturbing the transmitted classical signal. In quantum cryptography proposals, security is guaranteed by using single-photon states [97, 98], some of the schemes employing particles prepared in an EPR-entangled state. Each collaborator receives one member of each correlated pair, and measures the polarization in a random basis. After repeating the process many times, the two then discuss publicly which bases were used for each measurement, but not the actual measurement results. The cases where different bases were chosen are not used for conveying the key, and may be discarded, along with instances where one party detected no photon. In cases where the same bases were used, however, the participants will now have correlated information, from which a random, shared key can be generated. As long as single photons are used, any attempt at eavesdropping, even one relying on QND, will necessarily introduce errors due to the uncertainty principle. If the eavesdropper uses the wrong basis to study a photon before sending it on to the real recipient, the very act of measuring will disturb the original state.

77.6.4 Issues in Causality

Outstanding causal paradoxes in optics include the paradox of Barton and Scharnhorst [99], closely related to the Casimir effect (recently given indirect confirmation [14]; see Sect. 76.3.3). In this case, the amplitude for light-by-light scattering is modified by the presence of closely-spaced parallel plates. It appears that this can lead to propagation of light in vacuum (albeit a vacuum "colored" by the presence of the Casimir plates) faster than c. Similar paradoxes have been pointed out in connection with localization of any particle in a quantum field theory [100] and with the Glauber theory of photodetection (see Sect. 75.5) [101]. It may be that the high momentum/energy components introduced by localizing a detector to within a wavelength lead to spurious "dark" detections [102], which can be mistaken for causality-violating events [103].

77.7 THE SINGLE-PHOTON TUNNELING TIME

77.7.1 An Application of EPR Correlations to Time Measurements

In this section we discuss experiments involving the quantum propagation of light in matter. Due to the sharp time correlations of the paired photons from spontaneous down-conversion, one can use the HOM interferometer (see Sect. 77.3.2, and Fig. 77.4) to measure very short relative propagation delay times for the signal and idler photons. One early application was therefore to confirm that single photons in glass travel at the group velocity [46]. At least until recently, the only quantum theory of light in dispersive media was an *ad hoc* one [104]. The shift of the interference dip resulting from a medium introduced into one of the interferometer arms can be accurately measured by determining how much the path lengths must be changed to compensate the shift and recover the dip. This result suggests that when looking for a microscopic description of dielectrics, it is unnecessary to consider the medium as being polarized by an essentially classical electric field due to the collective action of all photons present, and reradiating accordingly. Linear dielectric response is not a collective effect in this sense — each photon interferes only with itself (as per Dirac's dictum) as it is partially scattered from the atoms in the medium. The single-photon group velocity thus demonstrates "wave-particle unity."

The standard limitation for measurements of short-time phenomena is that to have high time-resolution, one needs short pulses (or at least short coherence lengths), but these in turn require broad bandwidths and are therefore very susceptible to dispersive broadening. It is a remarkable consequence of the EPR energy correlations of the down-conversion photons (Sect. 77.3.1) that time measurements made with the HOM interferometer are essentially immune to such broadening [105, 106, 46]. In effect, the measurement is sensitive to the *difference* in emission times while the broadening is sensitive to the *sum* of the frequencies. While frequency and time cannot both be specified for a given pulse, the crucial feature of EPR correlations such as those exhibited by down-converted photons is that this difference and this sum correspond to commuting observables, and both may be arbitrarily well defined. The photon which reaches detector 1 could either have traversed the dispersive medium and been transmitted, or traversed the empty path and been reflected, leaving its twin to traverse the medium. The medium thus samples *both* of the (anticorrelated) frequencies, leading to an automatic cancellation of any first-order (and in fact, all odd-order) dispersive broadening. Measurements can be more than 5 orders of magnitude more precise than would be possible via electronic timing of direct detection events, and in principle better than those performed with nonlinear autocorrelators (which rely on the same fundamental physics as down-conversion, but do not benefit from a cancellation of dispersive broadening).

77.7.2 Superluminal Tunneling Times

Another well-known problem in the theory of quantum propagation is the delay experienced by a particle as it tunnels. There are difficulties associated with calculating the "duration" of the tunneling process, since evanescent waves do not accumulate any phase [107–109]. First, the kinetic energy in the barrier region is negative, so the momentum is imaginary. Second, the transit time of a wavepacket peak through the barrier, defined in the stationary phase approximation by

$$\tau^{(\phi)} \equiv \partial[\arg t(\omega)e^{ikd}]/\partial\omega \ , \qquad (77.15)$$

tends to a constant as the barrier thickness diverges, in seeming violation of relativistic causality[3]. For example,

[3] Actually, it is shown in [43] that such saturation of the delay time is a natural consequence of time-reversal symmetry, and in [110] that one can deduce from the principle of causality itself that *every* system possesses a superluminal group delay, at least at the frequency where its transmission is a minimum.

the transmission function for a rectangular barrier,

$$t(k,\kappa) = \frac{e^{-ikd}}{\cosh \kappa d + i\frac{\kappa^2-k^2}{2k\kappa}\sinh \kappa d}, \tag{77.16}$$

leads in the opaque limit ($\kappa d \gg 1$) to a traversal time of $2m/\hbar k\kappa$, independent of the barrier width d. The same result applies to photons undergoing frustrated total internal reflection [43], when the mass m is replaced by $n^2\hbar\omega/c^2$, and similar results apply to other forms of tunneling.

Some researchers have therefore searched for some more meaningful "interaction time" for tunneling, which might accord better with relativistic intuitions and perhaps have implications for the ultimate speed of devices relying on tunneling [111]. The "semiclassical time" corresponds to treating the *magnitude* of the (imaginary) momentum as a real momentum. This time is of interest mainly because it also arises in Büttiker and Landauer's calculation of the critical timescale in problems involving oscillating barriers, which they take to imply that it is a better measure of the duration of the interaction than is the group delay [107]. The Larmor time [112] is one of the early efforts to attach a "clock" to a tunneling particle, in the form of a spin aligned perpendicular to a small magnetic field confined to the barrier region. The basic idea is that the amount of Larmor precession experienced by a transmitted particle is a measure of the time spent by that particle in the barrier. This clock turns out to contain components corresponding both to the distance-independent "dwell time" and the linear-in-distance semiclassical time. Curiously, the most common theories for tunneling times become superluminal in certain cases anyway, whether or not they deal with the motion of wave packets.

Here, we will restrict ourselves to discussing the time of appearance of a peak of a single-photon wave packet. While other tunneling-time experiments have been performed in the past [113], optical tests offer certain unique advantages [114], including the ease of construction of a barrier with no dissipation, very little energy-dependence, and a superluminal group delay. The transmitted wave packets suffer little distortion, and are essentially indistinguishable from the incident wave packets. At a theoretical level, the fact that photons are described by Maxwell's (fully relativistic) equations is an important argument against interpreting superluminal tunneling predictions as a mere artifact of the nonrelativistic Schrödinger equation. Also, one is denied the recourse suggested by some workers [115] of interpreting the superluminal appearance of transmitted peaks to mean that only the high-energy components (which, for matter waves, traveled faster even before reaching the barrier [116]) were transmitted.

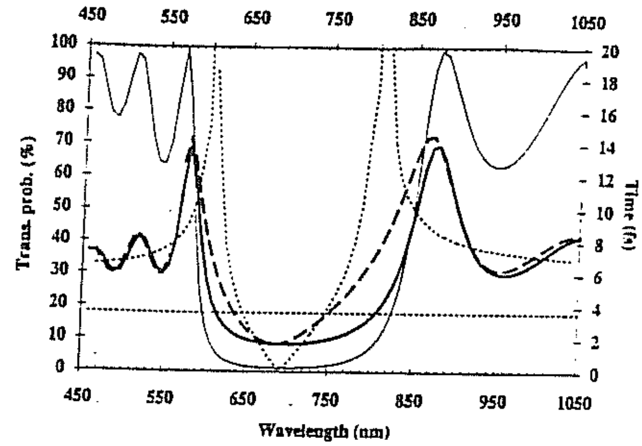

Figure 77.11. Transmission probability for the tunnel barrier used in [117] (light solid curve); heavy, dashed, and dotted curves show group delay, Larmor time, and semiclassical time. $d/c = 3.6$ fs is shown for comparison.

77.7.3 Tunneling Delay in a Multilayer Dielectric Mirror

A suitable optical tunnel barrier can be a standard multilayer dielectic mirror. The alternating layers of low and high index material, each one quarter-wave thick at the design frequency of the mirror, lead to a photonic bandgap [118] analogous to that in the Kronig-Penney model of solid state physics (see Sect. 76.3.5). The gap represents a forbidden range of energies, in which the multiple reflections will interfere constructively so as to exponentially damp any incident wave. The analogy with tunneling in nonrelativistic quantum mechanics[4] arises because of the exponential decay of the field *envelope* within the periodic structure, i.e., the imaginary value of the quasimomentum. The same qualitative features arise for the transmission time: for thick barriers, it should saturate at a constant value, as was verified in a recent experiment employing short *classical* pulses [125].

The phenomenon was investigated at the single-photon level by using the high time-resolution techniques discussed in Sect. 77.7.1 to measure the relative delay experienced by down-conversion photons [117] when such a tunnel barrier (consisting of 11 layers) was introduced

[4] A more direct analogy, that of waveguides beyond cutoff, yielded similar results in a classical microwave experiment [119], while another paper has reported superluminal effects related to the penetration of diffracted or "leaky" microwaves into a shadow region [120]. All these experiments involve very small detection probabilities, just as in Chu and Wong's pioneering experiment on propagation within an absorption line [121] (see Sect. 67.5). However, it has been predicted that superluminal propagation could occur *without* high loss or reflection [122–124], by operating outside the resonance line of an *inverted* medium (see Sect. 68.2). One can understand the effect as off-resonance "virtual amplification" of the leading edge of a pulse.

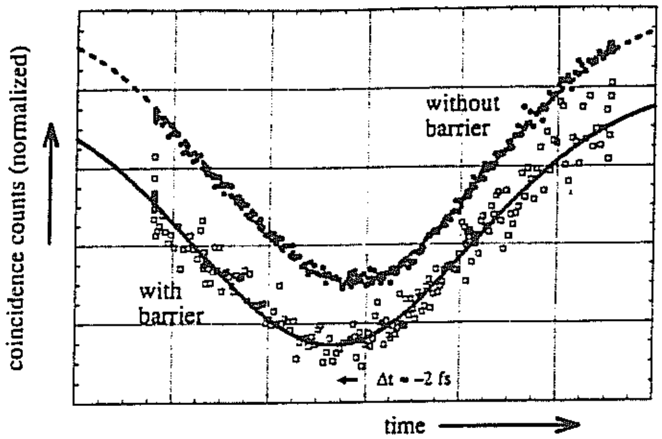

Figure 77.12. Coincidence profiles with and without the tunnel barrier map out the single-photon wave packets. The lower profile shows the coincidences with the barrier; this profile is shifted by ≈ 2 fs to *negative* times relative to the one with no barrier (upper curve): the average particle which tunnels arrives *earlier* than the one which travels the same distance in air.

into one arm of a HOM interferometer. The transmissivity of the barrier was relatively flat throughout the bandgap (extending from 600 nm to 800 nm; see Fig. 77.11), with a value of 1% at the gap center (700 nm), where the experiment was performed. The HOM coincidence dip was measured both with the barrier (and its substrate) and with the substrate alone (see Fig. 77.12). Each dip was subsequently fitted to a Gaussian, and the difference between their centers was calculated. When several such runs were combined, it was found that the tunneling peak arrived 1.47 ± 0.21 fs *earlier* than the one traveling through air, in reasonable agreement with the theoretical prediction of approximately 1.9 fs. Taking into account the 1.1-μm thickness of the barrier, this implies an effective photon tunneling velocity of $1.7c$. The results exclude the "semiclassical" time but are consistent with the group delay. A recent investigation of the energy-dependence of the tunneling time has been performed by angling the dielectric mirror, thus shifting its bandgap [126]. The data confirm the group delay in this limit as well, and rule out identification of Büttiker's Larmor time with a peak propagation time.

77.7.4 Interpretation of the Tunneling Time

Even though a wave packet peak may appear on the far side of a barrier sooner than it would under allowed propagation, it is important to stress that no information is transmitted faster than c, nor on average is any energy. These effects occur in the limit of low transmission, where the transmitted wavepacket can be considered as a "reshaped" version of the leading edge of the incident pulse [127, 121, 128]. At a physical level, the reflection from a multilayer dielectric is due to destructive interference among coherent multiple reflections between the different layers. At times before the field inside the structure reaches its steady-state value, there is little interference, and a non-negligible fraction of the wave is transmitted. This preferential treatment of the leading edge engenders a sort of "optical illusion," shifting the transmitted peak earlier in time. A signal, such as a sharp onset, relies on high-frequency components, which would not benefit from this illusion, but instead travel slower than c. Even for a smooth wave packet, no energy travels faster than light; most is simply reflected by the barrier. Only if one considers the Copenhagen interpretation of quantum mechanics, with its instantaneous collapse, does one find superluminal propagation of *those particles* which happen to be transmitted. This leads to the question of whether it is possible to ask which part of a wave packet a given particle comes from.

One paper argued that transmitted particles do in fact stem only from the leading edge of the wave packet [129]. While it is true that the transmission only depends on causally connected portions of the incident wave packet, further analysis revealed that *simultaneous* discussion of such particle-like questions and the wave nature of tunneling ran afoul of the complementarity principle [130]. In essence, labeling the initial positions of a tunneling particle destroys the careful interference by which the reshaping occurs (as in the quantum eraser). However, one picture in which the transmitted particles really *do* originate earlier is the Bohm–de Broglie model of quantum mechanics [131, 132]. This theory considers Ψ to be a real field (residing however in configuration space, thus incorporating nonlocality) which guides pointlike particles in a deterministic manner. It reproduces all the predictions of quantum mechanics without incorporating any randomness; the probabilistic predictions of QM arise from a range of initial conditions. Bohm's equation of motion has the form of a fluid-flow equation, $\mathbf{v}(\mathbf{x}) = \hbar\nabla \arg \Psi(\mathbf{x})/m$, implying that particle trajectories may never cross, as velocity is a single-valued function of position. Consequently, all transmitted particles originate earlier in the ensemble than all reflected particles [133, 130]. This approach yields trajectories with well-defined (and generally subluminal) dwell times in the barrier region. However, the fact that the mean tunneling delay of Bohm particles diverges as the incident bandwidth becomes small, along with other interpretational issues [134, 135], leaves open the question of whether time scales as defined by the Bohm model have any physical meaning.

The "weak measurement" approach of Aharonov *et al.* [136, 137] or equivalently, complex conditional probabilities obeying Bayes's theorem [138, 139], can be used to address the question of tunneling interaction times in an experimentally unambiguous way. The real part of the resulting complex times determine the effect a tunneling particle would have on a "clock" to which it coupled, while the imaginary part indicates the clock's back-action on the particle. They unify various approaches such as the Larmor and Büttiker-Landauer times, as well as Feynman-path methods. In addition, they allow one to discuss separately the *histories* of particles which have been transmitted or reflected by a barrier, rather than discussing only the wave function as a whole. Interestingly, these calculations do *not* support the assertion that transmitted particles originate in the leading edge of a wave packet.

77.8 THE REALITY OF THE WAVE FUNCTION

77.8.1 Status of the Wave Function in Measurement

The odd behavior of tunneling wave packets pushes us to reexamine our interpretation of wave functions. We customarily use the wave function only as a calculational tool, but we have also learned that it is in some sense physical, and should not be regarded merely as some distribution from classical statistics. A recent proposal [140, 141] suggests that the wave function of a single particle should be regarded as a real entity. When a state is "protected" from change, e.g. by an energy gap, and measurements are performed sufficiently "gently", one should be able to determine not just the expectation value of position, but the wave function at many different positions, without altering the state of the particle[5]. No violation of the uncertainty principle or of the no-cloning theorem (see Sect. 77.6.2) arises from this, as the ability to "protect" a state relies on some preexisting knowledge about the state; but it assigns a deeper significance to the wave function, one Aharonov terms "ontological," as opposed to merely epistemological (but see also [143]).

77.8.2 "Interaction-Free" Measurements

The real (but undetectable) field of the Bohm model, reminiscent of the "vacuum fluctuations" of quantum field theory, can produce observable effects. This can be seen particularly well in considering a recently proposed *gedanken* experiment [144], although it can be described perfectly well by standard QM. Consider a simple single-particle interferometer, with particles injected one at a time. The path lengths are adjusted so that all the particles leave a given output port (A), and never the other (B). Now suppose that a nontransmitting object is inserted into one of the interferometer's two arms — to emphasize the result, we consider an infinitely sensitive "bomb", such that interaction with even a single photon will cause it to explode. By classical intuition, any attempt to check for the presence of the bomb involves interacting with it in some way, and by hypothesis this will inevitably set it off.

Quantum mechanics, however, allows one to be certain some fraction of the time that the bomb is in place, *without* setting it off. After the first beam splitter of the interferometer, a photon has a 50% chance of heading towards the bomb, and thus exploding it. On the other hand, if the photon takes the path without the bomb, there is no more interference, since the *non*explosion of the bomb provides *welcher Weg* ("which way") information (see Sect. 77.4). Thus the photon reaches the final beam splitter and chooses randomly between the two exit ports. Some of the time (25%), it leaves by output port B, something which never happened in the absence of the bomb. This immediately implies that the bomb (or some object) is in place — even though (since the bomb is unexploded) it has not interacted with any photon. It is the possibility that the bomb *could* have interacted with a photon which destroys interference. An experimental implementation of these ideas [145] used down-conversion to prepare the single photon states (see Sect. 77.1), and a single-photon detector as the "bomb".

How can the final state depend on what is in an arm where no particle has passed? According to the Bohm theory the real Ψ-field *does* interact with the bomb, but in the absence of the particle is incapable of exploding it. At the final detector, Ψ, which now depends on whether or not the bomb is present, "pilots" the particle accordingly. Thus the "empty" Bohm field here plays a role very similar to that of the vacuum fluctuations in the two-crystal *welcher Weg* experiment described in Sect. 77.4.2. In such a picture for this experiment, it is the vacuum (incapable of triggering the bomb) which interferes with the fields at the beam splitter [146, 147], influencing which port a photon exits. The placement of a bomb in the path of the field destroys the coherence on which interference relies[6]. The choice the photons make at the final beam splitter is thus affected not by direct action of the bomb on any photons, but by its influence on the fluctuating fields permeating the vacuum.

An improved method has recently been discovered with which one can make the fraction of interaction-free measurements arbitrarily close to 1 [145]. For example, consider a photon initially in cavity #1 of two identical cavities coupled by a lossless beam splitter

[5]This idea of measuring the entire wave function of a single particle is often erroneously conflated with Raymer *et al.*'s fascinating work on the reconstruction of the quantum state of a light field by repeated sampling of a large ensemble: see [142] and Sect. 75.5.1.

[6]Even if the bomb's trigger is of the QND sort, the number-phase uncertainty relation guarantees this.

whose reflectivity $R = cos^2(\pi/2N)$. If the photon's herence length is shorter than the cavity length, after N cycles the photon will with certainty be located in cavity #2, due to an interference effect. However, if cavity #2 instead contains an absorbing object (e.g., the ultra-sensitive bomb), at each cycle there is only a small chance ($= 1 - R$) that the photon will be absorbed; otherwise, the *non*absorption projects the photon wave packet entirely back onto cavity #1. After all N cycles, the total probability for the photon to be absorbed by the object is $1 - R^N$, which goes to 0 as N becomes large. The photon effectively becomes trapped in cavity #1, thus indicating umambiguously the presence of the object in cavity #2.

77.9 GRAVITATIONAL ANTENNAS

According to general relativity, gravitational radiation can be produced and detected by moving mass distributions [148–150]. However, gravitational radiation is coupled only to time-varying mass *quadrupole* moments in lowest order, since the mass dipole moment is $\sum_j m_j \mathbf{r}_j = M\mathbf{R}_{cm}$ and $\mathbf{R}_{cm}$ for a closed system can only exhibit uniform rectilinear motion. Current efforts focus on detecting gravitational waves (typically at 100 Hz to kHz) from astrophysical sources, such as supernovae or collapsing binary stellar systems. For example, it is expected that in the nearby Virgo cluster of galaxies, several such events should occur per year, each yielding a fractional strain ($\Delta L/L$) of 10^{-21} on Earth. However, there are large uncertainties in this estimate.

There are at present two main efforts toward gravitational radiation detection[7]. A detector of the first type is known as a resonant-mass detector, or "Weber bar" after its inventor. It utilizes a cylindrical mass, whose fundamental mode of acoustical oscillation is excited by time-varying tidal forces produced by the passage of a gravitational wave. The induced motions are typically detected by piezoelectric crystals, or by SQUIDs (superconducting quantum interference devices) [152], yielding strain sensitivities better than $10^{-18}/\sqrt{\text{Hz}}$. Such detectors have been in operation since before 1977 [150], but to date no incontrovertible detections have been reported. Attempts to improve the signal-to-noise ratio here led to the consideration of back-action-evading sensors [a special case of QND measurements; see (77.2.4)] to circumvent the standard quantum noise limit [153].

[7]Like radio antennas, gravitational antennas must obey the principle of reciprocity, so they could in principle also be used as sources, with equal efficiencies. Currently this is infeasible, due to the small value of the coupling constant between mass and field (Newton's constant G), and the resulting miniscule efficiencies of the detection systems considered here. If a practical system overcoming these limitations were found [151], it would allow one to transmit and receive gravitational radiation, and perform a Hertz-type experiment.

The second kind of detector utilizes a Michelson interferometer. A passing gravitational wave alters the relative path length in the interferometer arms, thereby slightly shifting the output fringes. Although the gravitational mass of the light is much smaller than that of the Weber bar, very long interferometer arms (planned as 4 km, with a Fabry-Perot cavity in each arm to increase the effective length) more than make up for this disadvantage. The US initiative, called LIGO (Laser Interferometer Gravitational-Wave Observatory) and the European version, VIRGO, are expected to have initial sensitivities of $10^{-21}/\sqrt{\text{Hz}}$ with an ultimate goal of $10^{-23}/\sqrt{\text{Hz}}$. The signal-to-noise ratio for the detection of a fringe shift depends on the power of the light; an additional external mirror will be used to recirculate the unmeasured light, thus increasing the stored light power up to 10 kW. Because the standard quantum noise limit of these detectors is ultimately determined by the vacuum fluctuations incident on the unused input ports of the interferometers, it is in principle possible to achieve reduced noise levels by using *squeezed* vacuum instead [146, 150] (Sect. 77.2.3). This has been demonstrated experimentally [154]. At present, however, serious problems involving seismic and thermal isolation, absorption and heating, intrinsic thermal noise and optical quality are being addressed. Coincidence operation between two LIGO interferometers located respectively in Hanford, Washington and Livingston, Louisiana is expected around the year 2000 [155].

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