Single-particle Nonlocality: A Proposed Experimental Test

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We show that controlled interference of a particle’s wavefunction can be used to perform a quantum mechanical measurement in an incomplete basis. This happens because the measurement projects the particle into a lower dimensional subspace of the Hilbert space of the incoming wave. It allows a sender (Alice) to signal the receiver (Bob) nonlocally, by Alice’s choosing to measure in a complete or incomplete basis (in general: bases of differing incompleteness). If experimentally confirmed, it furnishes a new quantum communication act: nonlocal transmission of a bit without concomitant causal communication. However, the question of its compatibility with special relativity remains unresolved.

I. INTRODUCTION

It has been known for a long time that two distant particles, when related by the property of quantum entanglement, can display remarkable nonlocal behaviour, eg., the violation of Bell type inequalities \cite{1}. Interestingly, a formally similar phenomenon, in which the entangled twin is replaced by the vacuum, has been predicted for single particles, too \cite{2}. In contrast, the single-particle nonlocality we talk about in the present article does not involve any entanglement. Instead, it refers to the familiar sudden localization of a particle wavefunction following a position measurement. As a result, the nonlocal effect we envisage is not of a two-party correlation type, but involving probability modulation in a quantum informationally non-trivial way.

The article is arranged as follows. The experimental set-up for and the basic idea behind the test are presented in the next section. The main result is derived more rigorously in Section II, and physically interpreted in Section IV, where we show that it owes its origin to measurement via in an incomplete basis via controlled interferometry. We then conclude with a brief final section.

II. EXPERIMENT

The proposed experimental test is quite simple. It involves a light beam (preferably laser, in view of its low divergence), of width $w$, from a spatially coherent source being split up into three parallel beams ($a$, $b$ and $h$) by means of two 50/50 beams splitters (BS1 and BS2), as shown in Figure 1. Beam $a$ ($b$) is deflected by mirrors M0, MA45 and MA60 (MB45 and MB60) so that both beams illuminate the horizontal region $k$. Mirrors M0, MA45 and MB45 are inclined to the horizontal at $\pm 45^\circ$, mirrors MA60 and MB60 at $\pm 60^\circ$. The size of region $k$ is therefore $w' \equiv w \sec 30^\circ$.

Beams $a$ and $b$ are observed by an observer, denoted Alice, and beam $h$ by another, denoted Bob. Alice is equipped with a movable vertically oriented detector system that can be positioned either at the position DB2, which covers region $k$, or DB1, which lies ahead (Figure 1). Path length compensator PLC is provided on beam $b$, so that both beams are in phase when incident on MA45 and MB45. After reflection from mirror MA60 (MB60), beam $a$ ($b$) has a vertical wavefront (at angle $120^\circ$, rather than $90^\circ$, to the direction of propagation). In this way, we can ensure that both beams $a$ and $b$ interfere constructively at all points in region $k$. Careful orientation of the mirrors and positioning of detector DA2 is crucial to the experiment.

The action of the beam splitters and the mirrors leaves the input photon as it approaches the detectors in the superposition state

$$|\psi\rangle = \int \left( \frac{\lambda'}{\sqrt{2}} |a(z)\rangle + |b(z)\rangle \right) dz,$$

(1)
where $|X(z)\rangle$ ($X = a, b, h$) labels a ray state at height $z$ in beam $X$, and $\lambda$ and $\lambda'$ are the linear amplitude densities of the beams.

Suppose Alice positions her detector at position DB1. If she detects a photon, it will be either at region $l$ or $m$. Because of the property of the beam splitters, the probability that Alice finds a photon in the arm $a$ or $b$, $P(a) = P(b) = 1/4$, whilst the probability $P(h)$ that Bob finds a photon in the arm $h$ is $1/2$. Writing

$$P(a) = P(b) = \int_w \frac{|X|^2}{2} dz = 1/4; \quad P(h) = \int_w \frac{|\lambda/\sqrt{2}|^2}{2} dz = 1/2.$$  \hspace{1cm} (2)

we find $\lambda = 1/\sqrt{w}$ and $\lambda' = 1/\sqrt{w'}$. The photon count rate in each arm is given by $C(a) = C(b) = (1/2)C(h) = C/4$ where $C$ is the photon count rate for the light entering BS1 from the source.

On the other hand, suppose she positions her detector at position DB2. She detects clicks only at region $k$, but she cannot distinguish whether her detection was produced by a photon coming through beam $a$ or $b$. By Feynman’s dictum that quantum mechanics (QM) requires that all indistinguishable transition processes interfere $[$4$]$, the beams interfere at $k$. Note that the mirrors are so oriented as to ensure that both beams interfere constructively at all points in $k$, assuming that the laser beams diverge insignificantly across the experiment. (If they diverge considerably, then path length compensators with thickness varying suitably across the light stream are inserted to ensure that interference at $k$ is in-phase at all points, which is always possible.) On account of interference, Alice expects to find a photon with probability

$$P(k) \propto \int_w \left| \frac{\lambda'}{2} + \frac{\lambda}{2} \right|^2 dz = 1.$$  \hspace{1cm} (3)

At the same time, we still have $P(h) \propto 1/2$. We need to renormalize the probabilities because $P(k) + P(h) = 1 + (1/2) \equiv \beta = 3/2 \neq 1$. Thus, $P(k) = 1/\beta = 2/3$ and $P(h) = 1/(2\beta) = 1/3$. Consequently, in this case, $C(k) = 2C/3$ and $C(h) = C/3$.

It is somewhat surprising that we have to renormalize the probabilities. Yet, later we shall see that the renormalization is inevitable on account of the incompleteness of the DA2 measurement basis. It would not be needed if, instead of the DA2 measurement, Alice uses a Young double slit set-up, wherein beam $a$ passes through one slit, and beam $b$ through another, to interfere on a screen lying ahead. Note that this also, as with the DA2 measurement, entails that the two beams are indistinguishably registered. But, as shown in the next section, the pattern of bright and dark fringes on the screen will be such that the integrated count will be $C/2$, so that $C(h)$ also remains $C/2$. One aspect of the difference between the two measurements is that the DA2 measurement is a controlled interference, in which the relative phase of the interfering rays is pre-determined, whilst in the double slit interference, being diffraction determined, it is not.

The renormalization requires a subtle shift in the probabilistic interpretation of the wavefunction. The conventional statement of probability conservation in the region of spatial overlap of beams $a(x, t)$ and $b(x, t)$

$$\frac{\partial}{\partial t}(a + b)(a^* + b^*) = \frac{i}{\hbar} \nabla \cdot \{((a^* + b^*) \nabla (a + b)) - (a + b) \nabla (a^* + b^*)\}.$$  \hspace{1cm} (4)

contains no reference to beam $c$. This raises the question how the $a$-$b$ interference can deplete probability in arm $c$. The answer, which is elucidated by the second quantization formalism, is that amplitudes are in this case not squared probabilities, but only proportional to squared probabilities. In general, they are first order correlation functions. Thus, Eq. (4) is, strictly speaking, not the equation for probability conservation, but that for first order correlation. It gives us a measure of how correlated a photon being found at $k$ is with the input photon. In the above experiment, the enhanced correlation in $k$ is compensated for by conservation-enforcing correlation current flows into $k$ according to Eq. (6). However, correlation density at $p$ remains unchanged, leading effectively to probability flow from the lower arm to upper arms.

In a double slit experiment, the pattern of bright-dark fringes indicate phase difference between the interfering modes. These phase difference terms, when integrated over, vanish, as expected of orthogonal states in Hilbert space. Hence it is worth pointing out that the fact that only constructive interference takes place at all points in $k$ does not violate the Hilbert space orthogonality of the beams. If $X(x, t)$ is the wavefunction of beam $X$ ($X = a, b, h$), then unitarity guarantees orthogonality, i.e., that the volume integral over all configuration space

$$\int a(x, t)b^*(x, t)d^3x = 0.$$  \hspace{1cm} (5)
Thus, even though the two beams interfere constructively in a small volume surrounding region \(k\), at some other regions of their overlap away from the screen they interfere destructively. This pattern of fringes in space will be such that Eq. (3) is non-trivially satisfied.

It is interesting that by merely splitting and indistinguishably re-combining her beam, Alice boosts her probability of finding a photon from 1/2 to 2/3, while weakening Bob’s from 1/2 to 1/3 without directly acting on his beam. But this also means that Bob can instantaneously discern whether Alice measured at DB1 or DB2, depending on whether he finds count rate C/2 or C/3 on arm \(h\). Here we note that this nonlocal effect does not affect the strictly local character of unitary evolution. It is during the DA2 measurement and the consequent renormalization that a nonlocal connection between beams \(a, b\) on the one hand, and \(c\) on the other becomes manifest. It is as if the interferometer beams must somehow talk to each other, no matter how far apart they are, so that probabilities remain consistent with the correlation functions.

Suppose, on grounds of Einstein causality, we require that Bob’s observed count rate on arm \(h\) should remain C/2. Then, because \(P(k) = 2 \times P(h)\), Alice should observe a count rate \(C\) at \(k\). The combined intensity output for Alice and Bob would then be \(3C/2\), which is greater than the input rate, in contradiction with energy conservation. On the other hand, what about somehow prohibiting the interference of the re-combining beams? This option is unlikely because the recombining beams are indistinguishably registered. It is not clear that we can prohibit interference at \(k\) without also prohibiting the familiar double slit interference, where indistinguishability between the slits guarantees interference, in agreement with Feynman’s above-mentioned dictum, and also consistent with the quantum optical formalism, as shown in the next Section.

Finally, we note that if we insert a phase shifter of phase \(\phi\) in path \(a\) at position DA1 in conjunction with detector DA2, Bob’s observed count rate is \(C(2 + \cos \phi)^{-1}\). In this way, the above experiment can be generalized for the transmission of continuum, as against discrete, signals.

### III. DERIVATION

We now derive more rigorously the result arrived at in the preceding section. The three-mode quantum optical state vector of the photon field just after leaving the beam-splitters is given by:

\[
|\Psi\rangle = |\text{vac}\rangle + \epsilon(|a\rangle + |b\rangle + \sqrt{2}|h\rangle),
\]  

where \(|\text{vac}\rangle\) is the vacuum state, \(|X\rangle (X = a, b, h)\) is the Fock state mode corresponding to some ray \(z\) in beam \(X\) of the interferometer and \(\epsilon\) depends on the source strength [5]. Let us consider configuration #1, wherein Alice positions her detector at DA1. The positive frequency part of the electric field at any of the three detectors is:

\[
E_Y^{(+)} = \lambda' e^{ikr_Y} \hat{Y} \quad (Y = a, b); \quad E_h^{(+)} = \lambda e^{ikr_h} \hat{h},
\]

where \(\hat{X}\) is the annihilation operator for the \(|X\rangle\) mode and \(r_X\) is the distance from the source to the detector (located at \(l, m\) or \(p\)) along arm \(X\) to a point \(z\) on the detector (Figure 4).

The single particle detection rate at \(z\), or intensity \(I(X) (X = a, b, h)\), obtained by Alice or Bob, is proportional to the first order correlation function, i.e, \(\langle E_Y^{-}(+)E_Y^{(+)\dagger}\rangle\), where \(\langle \cdots \rangle\) indicates the expectation value in the state \(|\Psi\rangle\). The total photon count rate \(C(X)\) is obtained by integrating the intensity over the width of beam \(X\):

\[
C(Y) \propto \int_{w} \langle E_Y^{-}(+)E_Y^{(+)\dagger}\rangle \, dz = \epsilon^2, \quad C(h) \propto \int_{w} \langle E_h^{-}(+)E_h^{(+)\dagger}\rangle \, dz = 2\epsilon^2,
\]

where we have used Eqs. (3) and (4). Thus, the chances of detection on any one of the arms \(a\) and \(b\) is half that on arm \(h\). On the other hand, let us consider configuration #2, wherein Alice positions her detector at DA2. In this case, the positive part of the electric field at \(k\) is given by:

\[
E_k^{(+)} = \lambda' \left( e^{ikr_{yk}} \hat{a} + e^{ikr_{zk}} \hat{b} \right),
\]

where \(r_{Yk} (Y = a, b)\) is the distance from the source to region \(k\) along a ray \(z\) in beam \(X\), while Bob’s field remains as in Eq. (4). The single counting rate at the region \(k\) is:

\[
C(k) \propto \int_{w} \langle E_k^{-}(+)E_k^{(+)\dagger}\rangle \, dz = 4\epsilon^2,
\]
assuming that, by the appropriate setting of path length compensator PLC in Figure 1, \( r_a = r_b + nk^{-1} \), for some integer \( n \). The expression for arm \( h \) remains as in Eq. (8). In an actual implementation, Eqs. (8) and (10) must be further modified to take into consideration the single slit diffraction effects and the profile of the laser beam.

Comparing Eqs. (5) and (10), we find that by measuring at \( k \), Alice has a two-fold greater chance of finding a photon than does Bob at \( h \). The total correlation is \( \beta \equiv 6 \hat{c}^2 \). Since the total power has to be conserved, Alice’s controlled interference boosts her observed rate to \( 4\hat{c}^2C/\beta = 2C/3 \), thereby producing a corresponding 2/3 factor depletion in Bob’s photon counts. We thus confirm the result of the preceding section. Similarly, we can also verify the result for the phase shifter being inserted into path \( a \) at DA1.

Suppose we replace DA2 with a Young double slit system, wherein \( a \) passes through one slit, and \( b \) through another. The electric field, \( E_y^{(+)} \), at some point \( y \) on the screen of size \( S \), would be

\[
E_y^{(+)} = \frac{1}{\sqrt{S}} \left( e^{i k(\rho_{aa} + r_{ay})} a + e^{i k(\rho_{bb} + r_{by})} b \right),
\]

where \( \rho_{aa} (\rho_{bb}) \) is the distance from the source to the upper (lower) slit along beam \( a (b) \), and \( r_{ay} (r_{by}) \) is the distance from the upper (lower) slit to point \( y \). The count rate over all \( y \) on the screen is:

\[
C(S) \propto \int_S \left( E_y^{(-)} E_y^{(+)} \right) dt = \int_S \frac{2\hat{c}^2}{S} \left( 1 + \cos[r_{ay}(t) - r_{by}(t)] \right) dt \sim 2\hat{c}^2,
\]

assuming that, by the appropriate setting of path length compensator PLC in Figure 1, \( r_{aa} = r_{bb} + nk^{-1} \), for some integer \( n \). Comparing Eqs. (8) and (12), we find \( C(S) = C(h) \). Hence, no nonlocal effect arises in this case.

The crucial point is to note that the field \( E_k^{(+)} \) in Eq. (9) automatically incorporates Feynman’s dictum that the modes \( a \) and \( b \) should superpose. Thus, we are not free to suppress interference at \( k \) in the DA2 set-up in order to forbid the nonlocality. The electric fields in Eqs. (9) and (11) have been written down in a similar way, by including all modes that are incident at a given point. There is no way, then, to modify \( E_k^{(+)} \) in Eq. (9) without also modifying \( E_y^{(+)} \) in Eq. (11). Since the latter follows straightforwardly from the second quantization formalism \( \Psi \), and is experimentally confirmed, so too are Eq. (9), and the correlation due to it, inevitable. What remains is to better understand the need for renormalization, which is taken up in the next section.

### IV. QUANTUM MECHANICAL PICTURE

Although interference experiments rightly belong to the domain of quantum optics (QO), many of them can usually be translated into quantum mechanical language (eg., cf. Ref. [5] as regards a number of multiphoton interferometric experiments, and Ref. [9] as regards the delayed choice experiment). Sometimes the latter version, even though not entirely valid, can be easier to physically interpret. To this end, let us revert back to first quantization formalism and view the optical modes in the experiment as Hilbert space eigenstates. Furthermore, for simplicity, we ignore the finite width of the beam, so that we replace Eq. (8) by the simplified version:

\[
|\psi\rangle = \frac{1}{2} \left( |a\rangle + |b\rangle + \sqrt{2} |h\rangle \right),
\]

where now \(|X\rangle (X = a, b, h)\) represent the state corresponding to the photon being found on beam \( X \). They constitute the eigenstates of the 3-dim Hilbert space of the beams.

We then interpret the measurements at DA1 and DA2 as corresponding to two different “which beam?” observables. In this notation, observable DA1 is given by the spectral decomposition

\[
\hat{O}_1 = l|a\rangle \langle a| + m|b\rangle \langle b| + p|h\rangle \langle h|,
\]

where \( p \) is the region on the beam \( h \) where Bob’s detector is placed. According to the von Neumann projection postulate \( \Psi \), measurement non-unitarily reduces the state vector to an eigenstate by the action of one of the projection operators

\[
\hat{P}_a \equiv |a\rangle \langle a|, \quad \hat{P}_b \equiv |b\rangle \langle b|, \quad \hat{P}_c \equiv |h\rangle \langle h|.
\]
The probability to find the photon at \( l, m \) or \( p \), given by \( P(l) = P(m) = P(p)/2 = 1/4 \), is obtained by the expectation value of the corresponding projector in the state \( |\psi\rangle \) of Eq. (13).

Now we can easily check the operational equivalence between the present QM picture and the quantum optical derivation in Section II: the projectors \( \hat{P}_Z \) \( (Z = l, m, p) \) are equivalent to the electric field operators \( E_X^{-}E_X^{(+)} \) \((X = a, b, h)\), in the sense that the \( \hat{P}_Z \)'s have, respectively, the same expectation value in the state \( |\psi\rangle \) of Eq. (13), as the \( E_X^{-}E_X^{(+)} \)'s have in the quantum optical state \( |\Psi\rangle \) in Eq. (1), in view of Eq. (8) (apart from a \( 4c^2 \) factor).

Measurement in the DA2 basis represents an incomplete measurement, because it cannot distinguish between beams \( a \) and \( b \). In a non-interferometric setting, the probability of an incomplete measurement on \( a \) and \( b \) is given by the measurement of the projector

\[
|a\rangle\langle a| + |b\rangle\langle b| = \hat{P}_l + \hat{P}_m.
\]

If valid, it would indeed suffice to prohibit the nonlocality discussed above. However, lacking cross-beam terms, it would also imply that the beams \( a \) and \( b \) don’t interfere at \( k \), which, as noted earlier, is not in keeping with Feynman’s aforementioned dictum (that indistinguishable transition processes should interfere), nor consistent with the quantum optical formalism.

Taking a clue from the form of the electric field operator \( E_k^{(+) \,\ast} \) given by Eq. (1), it is straightforward to find that, to be equivalent to the operator \( E_k^{(-)}E_k^{(+)} \) in Eq. (1), the incomplete projector corresponding to Alice’s DA2 measurement should have the form:

\[
\hat{P}_k \equiv \hat{P}_{l+m} = (\alpha|a\rangle + \beta|b\rangle)(\alpha^\ast\langle a| + \beta^\ast\langle b|) \neq \hat{P}_l + \hat{P}_m,
\]

in view of Eq. (13). In Eq. (17), we set the phase factors \( \alpha = e^{ik(r_+k)} \) and \( \beta = e^{ik(r_-k)} \) in order that \( \langle \hat{P}_k \rangle \) with respect to state \( |\psi\rangle \) should yield an (unnormalized) probability \((= 1)\) in consonance with Eq. (3) and Eq. (14), apart from the \( 4c^2 \) factor. \( \hat{P}_k \) is the required incomplete projector that takes into account the interference between beams \( a \) and \( b \). It yields the correct expression even when the two beams are out of phase by \( \phi \).

We are now in a position to appreciate why the renormalization is inevitable. The Hilbert space available to the experiment is spanned by \( |X\rangle \) \((X = a, b, h)\). The projectors \( \{\hat{P}_l, \hat{P}_m, \hat{P}_p\} \) for the DA1 measurement basis are complete in the sense that

\[
\hat{P}_l + \hat{P}_m + \hat{P}_p = I,
\]

where \( I \) is the identity operator. On the other hand, the projectors for the DA2 measurement, namely \( \{\hat{P}_k, \hat{P}_p\} \), clearly do not constitute a complete basis set since

\[
\hat{P}_k + \hat{P}_p \neq I.
\]

They complete a subspace of lower dimensionality \((= 2)\), spanned by vectors \( \alpha|a\rangle + \beta|b\rangle \) and \( |h\rangle \), to which \( |\psi\rangle \) is projected when \( \hat{O}_2 \) is measured. That the completeness of basis is necessary to prohibit nonlocal transmission of classical information was first explicitly noted in Ref. [4].

One can press ahead with the quantum mechanical picture and, in analogy with Eq. (14), write down a formal spectral decomposition for the DA2 ‘observable’

\[
\hat{O}_2 = k\hat{P}_k + p\hat{P}_p.
\]

The normalization of a state vector is tailored to yield the true probabilities only for, and necessarily for, measurements in a complete basis. This is true irrespective of whether the measurements are complete or incomplete (i.e., represented by the projectors of the type Eq. (16)), projective or effect valued. Hence, measurement in the incomplete basis of \( \hat{O}_2 \) requires the renormalization. For the same reason, \( \langle \hat{O}_2 \rangle \) yields an expectation value that must be adjusted for normalization (here by a factor of \( 2/3 \)). Reverting back to second quantization, we attribute the renormalization to the fact that the electric field operators, \( E_k^{(+) \,\ast} \) and \( E_p^{(+) \,\ast} \), in the DA2 setup do not determine a complete basis in the Fock space available to the experiment. We note that as an intrinsic limitation of the first quantization picture, the projectors \( \hat{P}_X \) \((X = l, m, p, k)\) are to be used only for calculating probabilities, and not for predicting the state in which the photon is left if \( \hat{P}_X \) is measured, a situation that is independent of the dynamical and statistical properties of QM.
No renormalization is needed for the Young double slit measurement on beams $a$ and $b$, because it does not correspond to an incomplete basis. It merely scrambles path information. The projector for a given point $y$ on the double slit screen in the ‘beam’ basis is $\hat{P}_k(y) \equiv S^{-1}(\alpha(y)|a\rangle + \beta(y)|b\rangle)(\alpha^*(y)|a\rangle + \beta^*(y)|b\rangle)$. The completeness of the basis set follows from seeing that

$$\int_S S^{-1}(\alpha(y)|a\rangle + \beta(y)|b\rangle)(\alpha^*(y)|a\rangle + \beta^*(y)|b\rangle)dy = \hat{P}_l + \hat{P}_m,$$

where $\alpha(y)$ and $\beta(y)$ are the phase of the modes $a$ and $b$ at point $y$. Let us disambiguate the twin usages of the word completeness in the present context: $\hat{O}_2$ corresponds to an incomplete measurement in an incomplete basis. It is of course the latter feature that is responsible for the nonlocality. On the other hand, $\hat{O}_1$ corresponds to an incomplete measurement in a complete basis.

By the action of $\hat{P}_k$, the state vector is said to undergo ‘coherent projection’ or ‘coherent reduction’ into a Hilbert space of lower dimensionality. In contrast, no such reduction occurs for ‘projective reduction’, obtained by the action of $\hat{P}_l + \hat{P}_m$. Coherent reduction is characteristic of controlled incomplete interferometric measurements. It brings further richness to quantum measurement and the ways that it can be used to exploit the essentially quantum phenomenon of superposition. What is encouraging is that quantum optical tests for it are well feasible, and can indeed be already gleaned from existing interferometric experiments.

That Alice may transmit a nonlocal classical signal by choosing to measure in a complete or incomplete basis (in general, bases incomplete in different eigenstates) is not incompatible with the no-signaling arguments of Refs. [10–18], because, as pointed out by Ghirardi et al. [10], the assumption of measurement in a complete basis is implicit in them. Moreover, they deal with entangled multipartite rather than single particle systems. In view of the transparent arguments leading to, but surprising conclusions following from, the general measurement reported here, its experimental test is quite valuable. As an indirect evidence for its feasibility, an entanglement-based version of this effect was reported in Ref. [19], wherein it is argued that the results reported by Zeilinger [20] correspond to measurement in incomplete bases. However, we note that from the viewpoint of implementation and practical application the single particle version is much simpler.

V. CONCLUSION

Feynman noted that the central mystery of QM—namely, superposition— is encapsulated by the double-slit interference [1]. Coherent reduction, as discussed above, adds a further perspective to this ‘mystery’. If confirmed, it makes for an interesting addition to the basic quantum communication acts: the nonlocal transmission of a bit without concomitant communication, classical or quantum, along a causal channel.

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FIG. 1. A beam of light of width $w$, entering from the left, is split into two at beam splitter BS1. The outgoing upper beam is in turn split into two beams, $a$ and $b$ at beam splitter BS2. Alice (Bob) observes $a$ and $b$ ($h$). The mirrors M0, MA45 (MB45) and MA60 (MB60) are used to fold ray $a$ ($b$), and get them to converge to region $k$. Alice can choose to detect her two beams either indistinguishably at $k$, by positioning her detector at DB2, or at regions $l$ and $m$, by positioning it at DB1. Path length compensator PLC on path $b$ and careful orientation of the mirrors ensure that beams $a$ and $b$ illuminate region $k$ in-phase. Bob detects beam $h$ using a detector positioned at region $p$. Mirrors M0, MA45 and MB45 are tilted at 45$^\circ$ to the horizontal, MA60 and MB60, at 60$^\circ$. Hence region $k$, assuming that the light beams spread very little during their transit, has a vertical extent $w' = w \sec 30^\circ$. 