Non-local geometric phase in two-photon interferometry

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Abstract – We report the experimental observation of the non-local geometric phase in Hanbury Brown-Twiss polarized intensity interferometry. The experiment involves two independent, polarized, incoherent sources, illuminating two polarized detectors. Varying the relative polarization angle between the detectors introduces a geometric phase equal to half the solid angle on the Poincaré sphere traced out by a pair of single photons. Local measurements at either detector do not reveal the effect of the geometric phase, which appears only in the coincidence counts between the two detectors, showing a genuinely non-local effect. We show experimentally that coincidence rates of photon arrival times at separated detectors can be controlled by the two-photon geometric phase. This effect can be used for manipulating and controlling photonic entanglement.

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Introduction. – Quantum interference is one of the most delicate and intriguing aspects of quantum theory. A subtle interference effect, discovered surprisingly late, is due to the geometric phase [1,2]. It was soon realized [3] that Berry’s discovery had been anticipated by Pancharatnam’s work on the interference of polarized light [4]. Pancharatnam’s work is now widely recognized as an early precursor and optical analog of the geometric phase [3,5].

Studies of the geometric phase usually deal with amplitude interferometry, which describes the interference of particles from coherent sources. Of particular interest to us here are multiparticle interference effects, discovered by Hanbury Brown and Twiss (HB-T), who performed intensity interferometry experiments using incoherent thermal sources [6]. In a previous paper [7], an experiment was proposed to detect a non-local Pancharatnam phase in polarized intensity interferometry. The idea of the experiment, cf. fig. 1, is to have two thermal sources of opposite circular polarizations, PR and PL, where R and L describe right and left hand circular polarizations, respectively. These illuminate two detectors that are covered by linear polarization analyzers, P3 and P4. A theoretical analysis of the experiment reveals that as the relative angle between the two linear polarizations, φ34, is changed, the rate of coincidence counts between the arrival time of photons at the two detectors varies sinusoidally [7].

In recent years there have been many interferometric studies involving pairs of photons [8–12]. Most of these studies use entangled photon pairs generated by spontaneous parametric down conversion and address effects like teleportation and Bell’s inequality. There has been one study of geometric phases in two-photon interference [11] involving parametrically down-converted photon pairs focussing on the interplay between geometric and dynamic phases. In contrast, in our work we use thermal sources to generate photon pairs and study the role of the geometric phase in inducing non-local cross-correlations.

Entanglement lies at the very heart of quantum physics. It is now widely recognized as a fundamental resource in the growing field of quantum information, regarding both quantum computation and communication protocols [13,14]. As with any resource, effective use requires control, and quantum information applications are especially concerned with the degree of entanglement between spatially separated quantum channels or systems. A multiparticle Aharonov-Bohm effect was discussed theoretically by Büttiker [15] who noted in the context of electronic charge transport that two-particle correlations can be sensitive to a magnetic flux even if the single-particle observables are flux insensitive. The effect of the flux is visible only in current cross-correlations and is a

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A. Martin et al.

Fig. 1: (Colour on-line) General scheme for the Pancharatnam phase-dependent two-photon intensity interferometry. Two sources illuminate two detectors by direct processes (blue, dotted lines), and by exchange processes (red, dashed-dotted lines). The sources are respectively followed by analyzers for opposite circular polarization, and analyzers for different linear polarizations are placed in front of the detectors.

Fig. 2: (Colour on-line) Path on the Poincaré sphere determining the detected geometric phase. The relative angle \( \varphi_{34} \) between the linear polarizers \( P_3 \) and \( P_4 \) sets the angular width of the enclosed surface on the sphere. \( \Omega = 4\varphi_{34} \) corresponds to the solid angle defined by the geodesic path by the geodesic path defined in \( \Omega \).

genuinely non-local and multiparticle Aharonov-Bohm effect [16]. This has been experimentally seen in intensity interferometry experiments carried out using edge currents in quantum Hall systems [17], and the theory is further developed in [18,19]. These authors propose an experiment in which Aharonov-Bohm fluxes may be used to control the orbital entanglement of electronic states. While this proposal is sound in principle, in practice, such experiments may be afflicted by decoherence [13,14]. In contrast to electronic states, photonic states are easier to generate, manipulate and distribute. Most studies related to quantum communication protocols use photons rather than massive particles. However, photons are neutral and unaffected by the magnetic fields needed in the Aharonov-Bohm effect. The idea here is to use the geometric phase as an effective magnetic field to control the degree of photonic entanglement [7].

As a step towards achieving this goal, we report in this letter an experiment based on two independent polarized thermal sources for a demonstration of the non-local Pancharatnam phase effect. Such a non-local effect can only be observed in the intensity cross-correlations \( G^{(2)} \).

A previously reported experiment involving single photons prepared in mixed polarization states led to lower-order correlations \( G^{(1)} \) [20]. Here we take advantage of the setup suggested in [7] and show that the geometric phase effect only appears in the coincidence counts between spatially separated detectors when the relative angle between the polarization analysers is tuned. The two-photon Pancharatnam phase is non-local in the precise sense that it cannot be observed when local measurements at either detector are performed. The coincidence counts between the two detectors are modulated by a phase which has a geometric component as well as a dynamical (or propagation) phase.

Ideally, the experiment should use a priori incoherent thermal sources, as in the HB-T experiment. However, natural thermal light sources would necessitate very short integration times (shorter than the coherence time), which is hard to implement experimentally. Standard laser sources, which are usually used in amplitude interferometry, are unsuitable since they are coherent and do not show the HB-T effect. In this experiment, we use a laser source and introduce incoherence by passing the laser beam through a ground glass plate [21]. This gives us a source of light in which the amplitude and phase are randomly varying, mimicking a thermal light source intensity distribution. Using such a (pseudo-thermal) source, we demonstrate the non-local Pancharatnam phase [7].

**Theory.** – The central quantity of interest is the polarized intensity-intensity correlation function measured by the coincidence counts for the detectors \( D_3 \) and \( D_4 \), given by

\[
C = G^{(2)}_{34} = \frac{\langle N_3N_4 \rangle}{\langle N_3 \rangle \langle N_4 \rangle},
\]

\( N_3 \) and \( N_4 \) being the photon numbers detected at \( D_3 \) and \( D_4 \) per time per bandwidth. The detailed theory is outlined in [7]. The theoretical prediction for \( C \) is

\[
C = \frac{3}{2} + \frac{1}{2} \cos \left[ k(r_{13} - r_{14}) - (r_{21} - r_{24}) + \frac{\Omega}{2} \right],
\]

where \( r_{ij} \), with \( i = 1, 2 \) and \( j = 3, 4 \), denote the optical path lengths between the sources \( S_1 \) and \( S_2 \) and detectors \( D_3 \) and \( D_4 \). \( \Omega \) is the solid angle enclosed by the geodesic path shown in fig. 2.

This theoretical calculation can be summarized as follows. There is a process in which a photon from \( S_1 \) reaches detector \( D_3 \) and a photon from \( S_2 \) reaches \( D_4 \) (direct process, shown in blue). There is also a corresponding exchange process in which a photon from \( S_1 \) reaches \( D_4 \) and one from \( S_2 \) reaches \( D_3 \) (exchange process, shown in pink). These two processes are indistinguishable, and in the absence of any “welcher Weg” information, we have to combine the amplitudes for these processes with a relative phase \( \phi \), which depends on the optical
pseudo-thermal light sources (S) (Fig. 1) is schematically shown in Fig. 3. Two independent polarizing beamsplitter; APD: avalanche photodiode. and λ/2: quarter- and half-waveplates, respectively; BS: non-polarizing beamsplitter; PBS: polarizing beamsplitter; L: lens; PBS: polarizing beamsplitter; APD: avalanche photodiode.

Our pseudo-thermal sources emit monochromatic, thermal light, with an adjustable coherence time. This is accomplished by focusing a 852 nm external cavity diode laser on a rotating, ground glass disk, producing path difference between the two processes. The quantum interference between these processes causes the HB-T effect. Only the interference between direct and exchange processes gives rise to oscillations, the remaining terms effect. Only the interference between direct and exchange processes causes the HB-T effect of time sampling frequency of 1 MHz.

Experiment. – Our experimental realization of Fig. 1 is schematically shown in Fig. 3. Two independent pseudo-thermal light sources (S1 and S2) are spatially filtered using single-mode fibers, and their polarizations are adjusted to right-hand and left-hand circular (P3 and P2), ϕ34, and are modulated by a geometric phase of half the solid angle Ω covered on the Poincaré sphere by the four polarizations involved (see Fig. 2). The closed path on the Poincaré sphere is a property of a pair of photons. The effect of the geometric phase on the coincidence rates of photons at the detectors D3 and D4 is what we measure in the experiment below.

Non-local geometric phase in two-photon interferometry

Fig. 3: (Colour on-line) Experimental setup for demonstrating the Pancharatnam phase at the two-photon level. S1 and S2: independent pseudo-thermal light sources; SMF: single-mode optical fibre; L: lens; PBS: polarizing beamsplitter; λ/4 and λ/2: quarter- and half-waveplates, respectively; BS: non-polarizing beamsplitter; APD: avalanche photodiode.

Fig. 4: (Colour on-line) Autocorrelation functions (G_{11}^{(2)}(τ) and G_{22}^{(2)}(τ)) of the pseudo-thermal light sources S1 and S2 (red and green), and their cross-correlation function (G_{12}^{(2)}(τ), blue), as functions of time using sampling frequency of 1 MHz.

The light received at P becomes uncorrelated with itself as the beam traverses the plate. For a Gaussian beam profile incident on the glass plate, we expect the intensity autocorrelation function \( G^{(2)}(\tau) = \langle I(0)I(\tau) \rangle / \langle I(0) \rangle^2 \) to show a Gaussian decay in time [24]. More precisely, \( G^{(2)}(\tau) = 1 + \exp(-\pi(\tau/\tau_c)^2) \), where the coherence time \( \tau_c \) is on the order of the time taken by a point on the glass plate to traverse the beam width. Such a Gaussian decay of correlations is indeed observed in our experiment.

The average intensities have been chosen such that both the probability of having more than one photon during the dead time of the APDs (45 ns), and the probability of having a dark count, are negligible. The effective coherence times of the sources, and the absence of correlations between them, have been tested by measuring their autocorrelations (\( G^{(2)}_{ii}(\tau) \)) and the cross-correlation function (\( G^{(2)}_{ij}(\tau) \)). The corresponding experimental results are shown in Fig. 4. As predicted, both sources feature a Gaussian correlation decay as a function of time, associated with coherence times of about 7–800 μs. In addition, they exhibit no cross-correlation, proving they are actually independent. Therefore, we have estimated that sampling frequencies down to 100 kHz will ensure good
measurements of $G^{(2)}(0)$, while $10^5$ samples give us a negligible error. These parameters have been chosen in order to minimize the measurement time, and thus avoiding any unwanted path length fluctuations in the interferometer, as well as frequency drifts of the laser.

To observe the interference pattern in the coincidence counts $\mathcal{C}$ (see eq. (2)), we rotate the $P_3$ analyzer’s half-waveplate by an angle $\theta$ from 0 to $2\pi$, thereby continuously scanning the polarization (and thus the angle $\varphi_{34}$, from 0 to $4\pi$). The recorded signal is shown in fig. 5, as well as the expected fringe pattern with an ideal visibility.

Discussion. – Apart from its intrinsic interest as a demonstration of the non-local geometric phase in optics, one of our prime motivations is to realize a controllable and reconfigurable circuit enabling the generation and the manipulation of all the four photonic entangled Bell states [14]. This can be done by extending our experiment in the following way. Let us first replace the two thermal sources of fig. 1 by two simultaneously heralded single-photon sources. We suppose that when two photons are emitted at the same time, i.e., one per source, they have opposite circular polarization. Second, we interface these polarized single-photon emitters and the current setup by a time-bin preparation circuit acting on the two photons [13], where time-bin here is the means for creating coherent superpositions of states in the time domain. This way, using beamsplitters and identical delay lines for the two photons, we can make these photons travel over short, $|s\rangle$, or long, $|l\rangle$, paths, with certain probability amplitudes, and accordingly prepare time-bin input qubits for the geometric phase setup. We prepare incident photons initially with $P_H$ and $P_L$ polarization in the time-bin superpositions $|\varphi\rangle = \alpha|s\rangle + \beta|l\rangle$ and $|\psi\rangle = \alpha'|s\rangle + \beta'|l\rangle$, respectively. Here, $\alpha$, $\beta$, $\alpha'$, and $\beta'$ are complex numbers related to the beamsplitting ratios. In this configuration [23], the output state is given by

$$|T\rangle = \frac{1}{\sqrt{2}} \left( |\varphi\rangle |\psi\rangle + e^{i\phi} |\psi\rangle |\varphi\rangle \right),$$

where $\phi$ can be tuned using the geometric phase.

On one hand, simply setting the input states to $|\varphi\rangle = |s\rangle$ and $|\psi\rangle = |l\rangle$ permits preparing time-bin entangled states of the form $(|s\rangle + e^{i\phi}|l\rangle)/\sqrt{2}$. Such states are otherwise non-trivial to produce by usual experimental means based on pulsed parametric down-conversion associated with a preparation interferometer [11]. Here, the geometric phase allows switching between the $|\Psi^+\rangle$ and the $|\Psi^-\rangle$ Bell states.

On the other hand, substituting in eq. (4) superposition input states of the form $|\varphi\rangle = (|s\rangle + e^{i\phi_1}|l\rangle)/\sqrt{2}$ and $|\psi\rangle = (|s\rangle + e^{i\phi_2}|l\rangle)/\sqrt{2}$, where $\phi_1$ and $\phi_2$ are the relative phases between short and long-time-bin paths for either photon, leads to

$$|T\rangle = \frac{1}{\sqrt{2}} \left[ \Phi(\pi + 2\phi_1) \cos \frac{\phi}{2} 
+ e^{i\phi}(\pi + 2\phi_1) |\Psi^-\rangle \sin \frac{\phi}{2} \right],$$

in which we have set $\phi_2 = \pi + \phi_1$. By setting $\phi = 0$, and changing the two other phases, we have access to any superposition of $|\Phi^-\rangle$ and $|\Phi^+\rangle$, and to these two basis states as well. This simple scheme therefore leads to the possibility of acting on and/or controlling the degree of entanglement. Regarding experimental realizations, note that the heralded photon sources and the entanglement preparation and control circuits could be realized in the near future using integrated optics, as discussed in [25]. This would make our scheme particularly suitable for practical realizations.

In the example above entanglement is generated by particle exchange effects rather than by interactions. The latter degree of entanglement can be quantified either using Bell’s inequality or by the von Neumann entropy of the reduced density matrix, after tracing over one of the subsystems (3 or 4). A straightforward calculation of the von Neumann entropy shows that it does depend on the geometric phase. Since the geometric phase is achromatic, we can apply the same phase over all the frequencies in the band of interest by tuning $\varphi_{34}$ and generate entangled photon pairs with the degree of accuracy and the control needed for a source of photonic entanglement.

There have been measurements [26] of the geometric phase in intensity interferometry reported previously, where the geometric phase is also manifest in the lower-order correlations, and local in effects at each detector. This may be viewed as two copies of the single-photon geometric phase, which either add or cancel (depending on their relative sign) rather than a genuine two-particle phase. As a result the phase observed in this experiment is not half the solid angle on the Poincaré sphere, but twice this value or zero. There is also a difference in the
nature of the source. The experiment described in [26] uses a pair of entangled photons as a source, while our sources are incoherent and thermal, just like in the original HB-T experiment.

Conclusion. – To conclude, we have experimentally demonstrated a simple generalization of the HB-T effect, making use of the vector nature of light in order to produce a non-local geometric phase. The only conceptual difference between the proposed experiment and the classic HB-T experiment is the presence of polarizers at the sources and detectors. These polarizers cause a geometric phase to appear in the coincidence counts between the two detectors that receive linearly polarized light. Neither the single count rates, nor the self-correlations of individual detectors show any geometric phase effects. These appear solely in the cross-correlation between the count rates of the detectors. This experimental demonstration will open up possibilities for entanglement tuning via the geometric phase.

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