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Revealing the quantum nature of quasi-classical states carrying orbital angular momentum

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ABSTRACT

The observation of the quantum nature of quasi-classical states is of utmost importance since these states are realistic sources of radiation and do not have any shortage in photon numbers. Therefore, they do not face the scalability problem as much as single-photon sources do, which makes them much more robust against photon loss. The physical dynamics of these states predict that the quantum aspects of a continuous source can be observed in the weak limit. Yet, since these states are superposition of different particle states of radiation, we need post-selection to achieve this result. Practically, even after the post-selection, this high photon rate has not been observed. In this manuscript, we attempt to demonstrate that quasi-classical states supplemented with orbital angular momentum (OAM) can deliver a significant high photon rate in producing quantum signature alluding to this fact that OAM is indeed an intrinsic property of light at the quantum level as well. Through our experiment, we realize two-photon bunching with the projection of OAM from a continuous wave source. We observe that this approach can produce the result from a single-photon source with the high accuracy after the post-selection. In fact, since there is no interaction with matter, the post-selected photon rate from this approach can be considered as an upper bound for the single-photon generation based on the input power. This work is a step forward toward a more diverse and practical use of quasi-classical states in the domain of quantum optics and quantum information.

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I. INTRODUCTION

The observation of two-photon bunching and its accompanying Hong–Ou–Mandel (HOM) effect in an interferometric setting is the signature of the quantum nature of light, and upon this observation, we need a fully quantized frame to correctly describe the radiation.¹ Sources satisfying this criterion can be safely used to probe the quantum world. This non-classical behavior is at the core of many modern quantum applications such as quantum teleportation, quantum computation, and quantum metrology.^{2–4} Quantum sources, such as heralded photons from spontaneous parametric down-conversion (SPDC), suffer from low photon numbers due to the poor efficiency of non-linear processes that produce them.^{5–7} This is an especially acute limitation if further optical processing is expected. In this regard, it has been highly desirable to find alternative sources that can overcome this shortage. Radiation from a continuous wave source is described by a set of states known as quasi-classical or coherent states. Quasi-classical states are quantum states with interesting features. These states are the eigenstates of the annihilation operator, and like the vacuum state, they define the standard quantum limit (SQL), i.e., they minimize the

uncertainty relations.^{8,9} In addition, they predict correct behavior of physical quantities like the average electric field in the classical regime.^{10,11} On the other hand, it is well known that they do not pass the anticorrelation test and thus cannot directly represent a single-photon source.¹² In fact, the observation of two-photon bunching and the HOM effect have a close connection with the anticorrelation test. One can observe these phenomena with sources that can pass the anticorrelation test.

In a physical system, the existence of additional degrees of freedom (DoFs) enables a higher capability of information processing. In earlier works, polarization has been the common DoF of light to investigate an optical system, but following the work of Allen *et al.*,¹³ the orbital angular momentum (OAM) of light was added. Since this addition provides a higher-dimensional Hilbert space, it enhances the capacity of information processing.¹⁴ This allows a high-dimensional storage of quantum information on a single photon. It has been shown that such generalization provides a vast range of theoretical extensions and experimental applications.^{15–18} It empowers a higher information capacity, enhanced security, lower disturbance, and stronger support

for the entanglement against local realistic theories as well as better efficiency in quantum computing.^{19–25} Examples employing OAM as a suitable DoF alone²⁶ or alongside other DoFs, such as polarization, are frequent. For instance, hyper-entangled states involve different physical variables, such as polarization, OAM, and time energy, and are entangled in each DoF.²⁷

In this work, we present to our knowledge the first experimental observation of two-photon bunching with the radiation described by quasi-classical states carrying OAM. This is realized by controlling the total state through polarization and OAM from a continuous wave (CW) source. In previous studies involving the quantum nature of quasi-classical states, radiation did not carry OAM.^{28–33} Essentially, this experiment demonstrates that quasi-classical state supplemented by OAM can deliver a significant higher photon rate in producing the quantum signature alluding to this fact that OAM is indeed an intrinsic property of light at the quantum level. One immediate benefit of this addition is the possibility of probing multi-photon interference based on the sorting approach of spatial modes of radiation. The idea is to sort different OAM states and then look at the sorted statistics of the interference.³⁴ Non-linear sources face low photon numbers. For example, the rate of pair production for a few milliwatt of power from a pump source delivered to the non-linear crystal is about $10^5 - 10^6$.³⁵ This deficiency is a serious issue when further optical applications are needed, especially when one deals with the spatial properties of photons.³⁶ It can exceedingly challenge the outcome of the experiment. Here, we aim to show that quasi-classical states carrying OAM from a CW source can overcome this serious obstacle. For the same power from the CW source in our experiment, the photon rate is much more efficient.

II. THEORY

From the dynamics of quasi-classical states, we can observe and find interesting facts. By definition, they are eigenstates of annihilation operator,

$$\hat{a}|\alpha\rangle = \alpha|\alpha\rangle, \quad (1)$$

in terms of number state, the coherent state $|\alpha\rangle$ is

$$|\alpha\rangle = \exp\left(-\frac{1}{2}|\alpha|^2\right) \sum_{n=0}^{\infty} \frac{\alpha^n}{\sqrt{n!}} |n\rangle. \quad (2)$$

It is clear from Eq. (2) that quasi-classical states do not represent single-photon sources since these states are the superposition of different particle states. The probability distribution of these states follows Poisson distribution,

$$P(m) = e^{-|\alpha|^2} \frac{|\alpha|^{2m}}{m!} = e^{-\mu} \frac{\mu^m}{m!}, \quad (3)$$

where

$$\mu = \langle \alpha | \hat{a}^\dagger \hat{a} | \alpha \rangle$$

is the average photon number over coherent states.

A specific line that we want to emphasize is the validity of particle nature of radiation using these states, i.e., the limit of quantum observations predicted by such states. The quantity of interest is the average electric field over these states for a monochromatic light that propagates along the z direction with $\mathbf{E}$ and $\mathbf{B}$ fields oscillating along x and y ,

respectively. The electric field in the quantized frame with the frequency ω and the wave vector $\mathbf{k}$ ($\omega/k = c$) is

$$\hat{E}_x(\mathbf{r}, t) = i \left(\frac{\hbar\omega}{2\epsilon_0 V} \right)^{\frac{1}{2}} [\hat{a} e^{i(\mathbf{k}\cdot\mathbf{r} - \omega t)} - \hat{a}^\dagger e^{-i(\mathbf{k}\cdot\mathbf{r} - \omega t)}], \quad (4)$$

where V is the volume of quantization. While clearly the average field is zero over the number state $|n\rangle$, the coherent state $|\alpha\rangle$ will produce what one observes in reality,

$$\begin{aligned} \langle \alpha | \hat{E}_x(\mathbf{r}, t) | \alpha \rangle &= i \left(\frac{\hbar\omega}{2\epsilon_0 V} \right)^{\frac{1}{2}} [\alpha e^{i(\mathbf{k}\cdot\mathbf{r} - \omega t)} - \alpha^* e^{-i(\mathbf{k}\cdot\mathbf{r} - \omega t)}] \\ &= 2|\alpha| \left(\frac{\hbar\omega}{2\epsilon_0 V} \right)^{\frac{1}{2}} \sin(\omega t - \mathbf{k}\cdot\mathbf{r} - \theta), \end{aligned} \quad (5)$$

in the last line we represent the coherent state in the polar form where $|\alpha\rangle = |\alpha|e^{i\theta}$. The uncertainty of the electric field over the coherent state is

$$\Delta \hat{E}_x = \left(\frac{\hbar\omega}{2\epsilon_0 V} \right)^{\frac{1}{2}} \equiv \mathcal{E}^{(1)}, \quad (6)$$

which is equal to the uncertainty calculated from the vacuum state $|0\rangle$, and $\mathcal{E}^{(1)}$ is the single-photon amplitude of the electric field. The behavior of this mean field is plotted in Fig. 1. Now, the main point that we want to look upon is the domain of validity of the quantized picture. For this purpose, we consider the ratio of the uncertainty bounds to the amplitude of the average field,

$$\frac{2\Delta \hat{E}_x}{2|\alpha|\mathcal{E}^{(1)}} = \frac{1}{|\alpha|} = \frac{1}{\sqrt{\mu}}. \quad (7)$$

For large values of μ , as expected, this ratio tends to zero, indicating a classical limit. However, a peculiar behavior arises when the average photon number μ becomes less than one. For example, with $\mu = 10^{-2}$, the ratio is 10, and for $\mu = 10^{-4}$, this ratio increases to 100. In the limit of $\mu \ll 1$, the average field tends to zero as well. Here, we emphasize that the significance of quantum effects becomes prominent and evident at low values of μ . It is at this level that quantum phenomena emerge, even though the quantum nature of radiation is inherently present at all times. Previously, we mentioned that the average field over the number state, i.e., $\langle \hat{E} \rangle_n$ is zero, and now, we observe from the

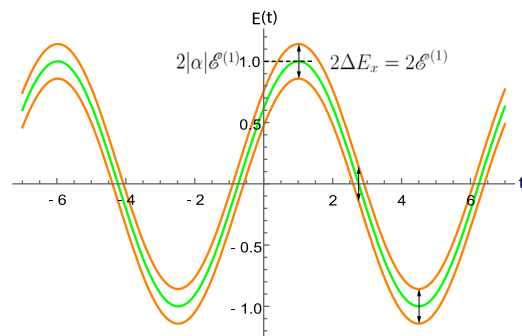


Fig. 1. Average electric field in the case of the coherent state with the uncertainty that is equal to the uncertainty of the vacuum state.

dynamics of the coherent states that this fact manifests itself at the limit of $\mu \ll 1$. This observation suggests that the quantum aspect of a continuous source represented by the dynamics of quasi-classical states occurs in the weak limit.

This implies that an observer can perceive the quantum aspect of radiation by approaching this weak limit without any interaction with matter. In this limit, the observer discovers that most of the radiation is devoid of photons, indicating that the dominant contribution arises from the vacuum state. For instance, with $\mu = 10^{-2}$, the probability of the single-photon state $|1\rangle$ is 10^{-2} relative to the vacuum $|0\rangle$, and the probability of observing two-photon state $|2\rangle$ is approximately 10^{-4} relative to the vacuum state. Now, depending on the precision that we want to reach, we can reduce the intensity of the source delivered to the optical system and achieve the desired μ . The only limitation to achieve a higher photon rate in this approach is the coherence time and coherence length of the source. A shorter coherence time corresponds to a higher photon rate. It is worth noting that since there is no interaction with matter through nonlinear processes like SPDC, in this approach the efficiency will be higher compared to other sources that rely on non-linearity for generating single-photon states. This is the case because we reach the quantum limit only by decoupling the source from the optical system. Now, there is an issue that needs to be addressed and a suitable solution needs to be found. Indeed, the quantum regime attainable through the weak limit is still not a perfect single-photon state. We have a superposition of different particle states with Poisson probability distribution. We need to isolate and separate the single-photon state. One method of separation that allows this is building an interferometric system. However, we will encounter a significant challenge in that case. The particle nature of radiation is unknown in the interferometric setting, and we cannot proceed to the measurement on the photon events unless we can resolve and address this problem. Mathematically, this challenge arises from this reality that the quasi-classical state $|\alpha\rangle$ is not an eigenstate of the number state operator $\hat{n}$. In fact, the solution is simple but profound, and it comes from the foundation of quantum mechanics itself. Introducing a temporal randomization in the interferometric setting and having a monochromatic radiation will pave the way, and it yields a consistent reading of the photon events. The details of this

process will be discussed in the following lines. Next, we develop the details of our approach, then describe our experimental setup, and lastly present the results.

As mentioned above, coherent states cannot directly produce the quantum mechanical effects observed by single-photon sources such as two-photon bunching or entanglement. The anticorrelation test does not allow quasi-classical states to represent single-photon sources because there is a nonzero probability of joint detection. Nevertheless, we deduced from our previous work the viability of observing purely quantum mechanical effects using these states in an interferometry setup and after post-selection. In this approach, we introduce a uniform time delay to one of the interfering beams. In an interferometric setting like Mach-Zehnder interferometry with arms a and b (see Fig. 3), the phase randomization process of the coherent state $|\alpha\rangle = |\sqrt{\mu}e^{i\theta}\rangle$ with the mean photon number $\mu = |\alpha|^2$ makes it possible to integrate over many temporal modes to ensure that the state of this signal is a mixed state with the Poisson distribution P_n in the number states representation,

$$\rho = \int \frac{d\theta}{2\pi} |\sqrt{\mu}e^{i\theta}\rangle \langle \sqrt{\mu}e^{i\theta}| = \sum_{n=0}^{\infty} P_n |n\rangle \langle n|. \quad (8)$$

However, this is an emergent behavior that comes from the foundation of quantum mechanics itself, i.e., uncertainty principle. By randomizing the temporal phase, we gain a stable measure on the energy content of the propagating beams through the interferometer. Since our radiation is monochromatic, the uncertainty of energy time evolves to a number state-phase uncertainty,

$$\Delta E \Delta t = \hbar \Rightarrow \Delta n \Delta \theta_t = 1, \quad (9)$$

which means that upon the act of randomization in the temporal phase θ_t , we achieve a reliable reading of the photon counts from the interferometer. Figure 2 demonstrates this fact that is related to the records from our time-controller interface. Here, the coincidence window is 10 ns, the integration time is 100 ms, and the acquisition time is 5 s. Without temporal randomization (left plot), the number state is perfectly random and occupies all possibilities, and then, from the middle plot to the right one, randomization process was applied in a way that there is higher randomization rate in the right plot compared

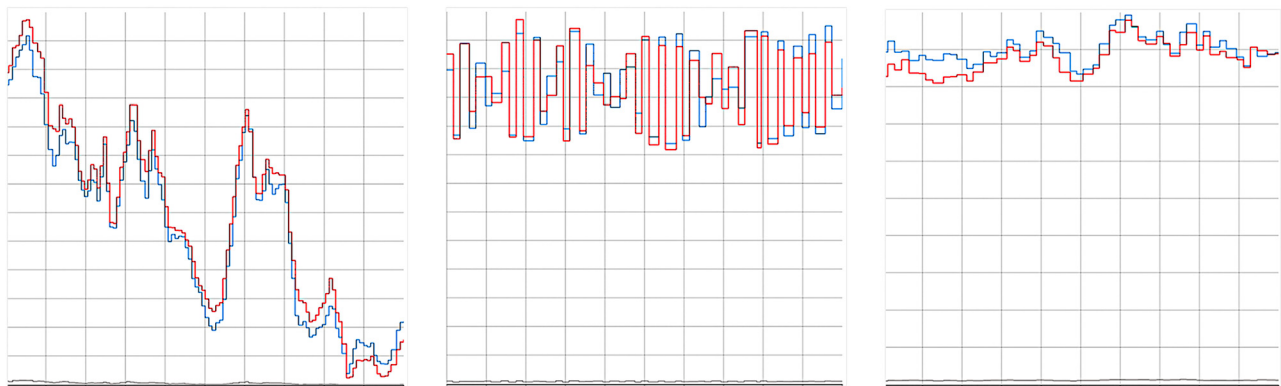


Fig. 2. Uncertainty measurements between temporal phase θ_t and number state n (vertical axis represents n , and the horizontal shows θ_t). Left: without any randomization, the number state is random. Middle: with randomization, there is a tight bound on n . Right: with a higher randomization rate relative to the middle case, there is a tighter bound on the number state n .

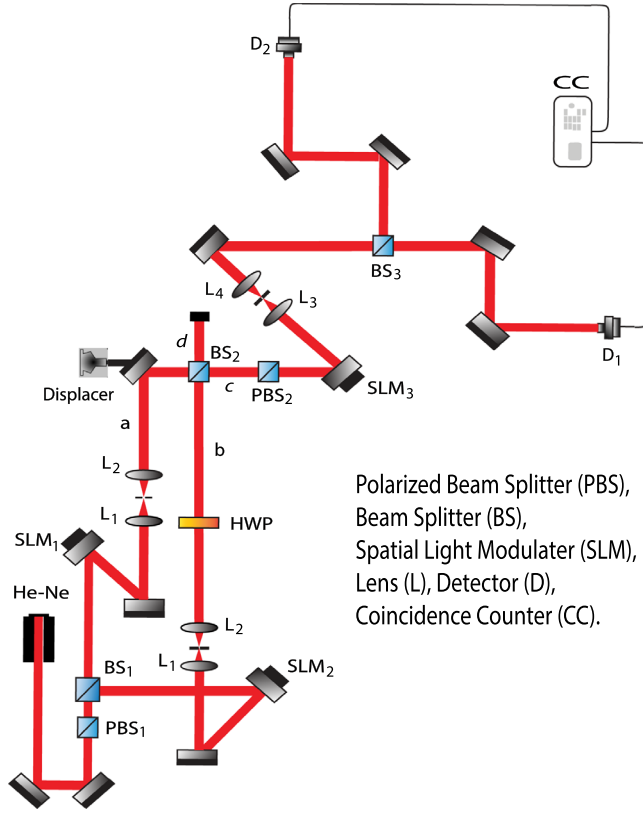


FIG. 3. Experimental setup for the observation of two-photon bunching using quasi-classical states carrying OAM. SLM₁ in arm *a* and SLM₂ in arm *b* of interferometer projects OAM states. After BS₂, one interfering beam is projected to the SLM₃ to do the measurement on the OAM space. HWP in arm *b* rotates the polarization of the projection.

to the middle one; hence, there is a tighter bound on the number state n . In the absence of this action, it is not possible to gather consistent data. Once this randomization is employed, we can filter the shares of higher-order photon numbers coming from each arm. This means that one measures the two-photon events, while both arms *a* and *b* are open, and since the state is mixed, one also needs to consider the possibility of photon pairs originating from each specific arm. We deal with this situation by finding separate contributions, i.e., photon pairs only coming from arm *a*, by closing arm *b*, and then measuring two-photon counts only from arm *b*, while arm *a* is closed. In this manner, we do three measurements for each data point and then subtract the second and the third counts from the first one (both arms being open) to retain the single-photon contributions to the interference. At this stage, it is necessary to point out some facts regarding the role of phase randomization. It is clear that we have the particle nature of radiation once there is no interference, i.e., radiation comes only from arm *a* or *b*. In these two cases, we do not need any randomization, and the act of measurement will integrate over the possible modes and Eq. (8) is satisfied. However, we do not have a meaningful measure on the number states in the case of both arms being open and having interference. Therefore, phase randomization in this case is essential, and we gain a

stable measure of the number state due to this process. In this regard, the uncertainty principle as in Eq. (9) is the underlying basis for the realization of Eq. (8).

The density operator representing the mixed state before the second balanced beam splitter (BS₂) is

$$\rho = \sum_{i,j} P_a(i)P_b(j)|i_{al_1}, j_{bl_2}\rangle\langle i_{al_1}, j_{bl_2}|, \quad (10)$$

where the probabilities $P_a(i)$ and $P_b(j)$ are Poisson distributions,

$$P_s(n) = \frac{\mu_s^n}{n!} e^{-\mu_s}, \quad s = a, b. \quad (11)$$

Here, μ_a and μ_b are the mean photon numbers at each input port of BS₂. State $|i_{al_1}, j_{bl_2}\rangle$ describes an eigenstate with $i(j)$ photon number of OAM $l_1(l_2)$ at the input port *a*(*b*). We suppressed the polarization DoF here. The goal is to separate two-photon events, among which we are merely interested in two-photon counts either at *c* or *d* ports. The two-photon events at port *c* are

$$\begin{aligned} \mathcal{N}_2(\mu_a, \mu_b) &= \mu_a \mu_b P(0_d, 2_c | 1_a, 1_b) + \frac{\mu_a^2}{2!} P(0_d, 2_c | 2_a, 0_b) \\ &+ \frac{\mu_b^2}{2!} P(0_d, 2_c | 0_a, 2_b) + \mathcal{O}(\mu^3). \end{aligned} \quad (12)$$

The first term on the right-hand side is related to the desired state, and the other two terms are unwanted. The contributions to the latter are coming from

$$\begin{aligned} \mathcal{N}_2(\mu_a, 0) &= \frac{\mu_a^2}{2} P(0_d, 2_c | 2_a, 0_b) + \mathcal{O}(\mu^3), \\ \mathcal{N}_2(0, \mu_b) &= \frac{\mu_b^2}{2} P(0_d, 2_c | 0_a, 2_b) + \mathcal{O}(\mu^3). \end{aligned} \quad (13)$$

The first line in Eq. (13) is the number of counts only coming from arm *a*, and second one is just measuring the counts from arm *b*. As argued earlier, one needs to remove these two terms from the signal, which leaves the pure contribution from the desired state,

$$\mu_a \mu_b P(0_d, 2_c | 1_a, 1_b) \simeq \mathcal{N}_2(\mu_a, \mu_b) - \mathcal{N}_2(\mu_a, 0) - \mathcal{N}_2(0, \mu_b). \quad (14)$$

To justify this approximation, it is important to mention that since the beam's intensity is low, the higher-order terms, $\mathcal{O}(\mu^3)$, can be ignored. Let us address an obvious limitation of this method. The approximation in Eq. (14) is valid for a weak source. However, as we argued earlier, this is a fundamental bound caused by the quantum nature of such radiation. The coherence time of the source is 1 ns. Therefore, with $\mu = 10^{-2}$, phase randomized weak coherent states (PRWCSs) can be safely practiced resulting in a photon rate of 10^7 , which is readily attainable with this approach as we shall see later. However, suppose we have to build more complex optical systems, then definitely this approach shines in photon production and establishes its potential in high photon loss scenarios.

The desired (post-selected) state before BS₂ is

$$|\Psi\rangle = |\psi_S\rangle \otimes |\psi_L\rangle, \quad (15)$$

where $|\psi_S\rangle$ and $|\psi_L\rangle$ are the state functions of the polarization and OAM subspaces, respectively,³⁷

$$|\psi_S\rangle = \frac{1}{\sqrt{2}} (|1_{a1s}, 1_{b2s}\rangle + e^{2i\phi} |1_{a2s}, 1_{b1s}\rangle) \quad (16)$$

and

$$|\psi_L\rangle = \frac{1}{\sqrt{2}} (|1_{a1l}, 1_{b2l}\rangle + e^{2i\phi} |1_{a2l}, 1_{b1l}\rangle), \quad (17)$$

where $|1_{a1s}, 1_{b2s}\rangle$ denotes one-particle states of light with the polarization s coming from paths a and b , respectively. The mismatched angle between arms a and b is ϕ . After the second BS, the states of interfering beams in each subspace become

$$|\psi_S\rangle = \frac{ie^{i\phi}}{\sqrt{2}} [(|1_{d1s}, 1_{c2s}\rangle - |1_{c1s}, 1_{d2s}\rangle) \sin \phi + (|1_{c1s}, 1_{c2s}\rangle + |1_{d1s}, 1_{d2s}\rangle) \cos \phi] \quad (18)$$

and

$$|\psi_L\rangle = \frac{ie^{i\phi}}{\sqrt{2}} [(|1_{d1l}, 1_{c2l}\rangle - |1_{c1l}, 1_{d2l}\rangle) \sin(l\phi) + (|1_{c1l}, 1_{c2l}\rangle + |1_{d1l}, 1_{d2l}\rangle) \cos(l\phi)]. \quad (19)$$

We utilize Laguerre–Gaussian (LG) modes having OAM dependence of the form $\exp(il\phi)$ and each photon carrying $l\hbar$ quanta of OAM. The probability of two photons being detected in the same output ports c and d is given by the second terms in each subspace corresponding to Eqs. (18) and (19), respectively. Therefore, the probability of having two-photon bunching at each port $c(d)$ is

$$P_{c_1c_2} = P_{d_1d_2} = \frac{1}{2} \cos^2 \phi = \frac{1}{4} [1 + \cos(2\phi)]; \quad l = 1. \quad (20)$$

In this work, the OAM quanta from SLM₁ and SLM₂ were fixed at $l_1 = l_2 = l = 1$, and the state containing this OAM number was projected to the SLM₃ generating OAM “−1,” i.e., $l_3 = -1$. Thereby, we can use the projection of these states to get the resultant signal with zero OAM and within the first diffraction order after SLM₃. The main reason for this choice of OAM quanta relates to this fact that $|l| = 1$ is the first non-trivial (non-zero) OAM state. Also, for this case, the observation probability of the desired phenomenon is the same in both subspaces. However, it is not necessary to have these values for the OAM. Indeed, as can be seen from the state $|\psi_L\rangle$, if one knows what the correct order of diffraction is, then with any arbitrary OAM state, the probability of having two-photon coalescence comes from the second term with the scaled phase in the cos function as $\cos(l\phi)$. In the experiment, since we aim to filter the OAM states through projection, it is clear that the right choice is to have $l_1 = l_2 = -l_3$ as depicted in Fig. 4. In this figure, we have recorded OAM projection up to $|l| = 10$. Moreover, it is shown that OAM generation and detection can be carried out consistently using SLM up to $l = 600$.³⁸ Thus, the generation of higher OAMs is not an issue here. As noted, having the value of OAM as $|l| = 1$ will produce the same probability in both polarization and OAM subspaces.

III. EXPERIMENT

The experimental setup is illustrated in Fig. 3. A Mach–Zehnder type interferometer is built involving two identical spatial light modulators (SLM₁ and SLM₂) as the projectors of photon states after reflecting the light from the source. The experiment implements a He–Ne

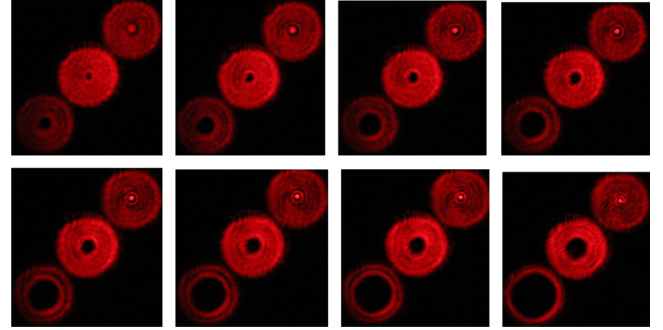


Fig. 4. Projection of beams carrying OAM after SLM₃. The zero OAM state is observed within the first diffraction order as a result of the projection of interfering beams carrying $l_1 = l_2 = 1, \dots, 7, 10$ with $l_3 = -1, \dots, -7, -10$ (numerating from the top left to the bottom right).

laser coupled with the single-mode fiber as a source of quasi-classical modes. After collimation of the beam, to achieve a uniform polarization, we use a polarized beam splitter (PBS₁). Then, the first balanced beam splitter (BS₁) is placed to split radiation into two beams for projection on the first and the second SLMs. These SLMs flatten the wavefront and project the same OAM quanta, for instance, “ l .” Through identical lenses and arm lengths, we bring the two beams together at the position of the BS₂. We thus have the interference pattern with two beams of the same polarization and OAM. Although in the figure the arm lengths may appear different, we have made sure that they have the same length (the aspect ratio of the arms a and b is not preserved). After BS₂, we project one of the interfering beams to the third SLM (SLM₃) loaded with the opposite OAM, i.e., “ $-l$.” We look for the result of this projection that has zero OAM, to be observed at the first diffraction order as in Fig. 4. We can then transfer this signal after splitting it by the third balanced beam splitter (BS₃) to the single-photon avalanche detectors (D₁ and D₂) using either a pinhole/aperture and multi-mode fibers or the single-mode fibers that act like filters to only pass the basic Gaussian mode. The photon rate coming out of the single-mode fiber as a result of OAM projection from a few milliwatt (mW) of power delivered to the interferometer is about $10^{11} - 10^{12}$. It is also possible to enhance this rate further with a more efficient alignment and optical manipulation. In practice, we can configure SLMs involved so that most of diffracted radiation will appear in the first order, and then, we can decrease definitely the power delivery to the system and enhance the efficiency. In fact, we have reached the interference with a visibility over 99% through the single-mode fiber after adding OAM. Before OAM addition, this number was less than 90%. This fact supports the previous statement that with additional DoFs in a physical system, we gain a better ground for information processing.

One critical component in our setup is a displacer in arm a that introduces a phase shift about 200π to the beam in this path. It creates a uniform temporal phase randomization between interfering beams, which plays a central role in the experiment as we shall explain shortly. This delay in one arm is much smaller than the coherence length of the source to ensure the indistinguishability of the interfering photons. The resultant signals known as phase randomized coherent states (PRCSs) are the main source of our investigation. We have a considerable photon rate, but the quantum effects will be observed in the weak

limit. Hence, we decouple the source to reduce the intensity of radiation delivered to the interferometer, and thus have PRWCSs. As stated above, we have the photon rate of 10^{12} after each SMF, and this rate is the result of a few milliwatt power delivered to the interferometer from the source, but this rate is not satisfying the weak limit. Therefore, we reduce the power to reach the weak limit as well as to have the rate compatible with the detectors in the experiment. An order of 10^6 (MHz) can be chosen with $\mu = 10^{-3}$. This corresponds to a few nanowatt power delivery to the optical system.

IV. RESULTS AND DISCUSSION

After explaining the concept and the math behind the PRWCSs, we are at the stage to present the results confirming the two-photon bunching using the coherent states carrying OAM. The data have been taken using polarization and OAM control of the signal. A half-wave plate (HWP) is inserted in arm *b* of the interferometer to vary the polarization phase mismatch ϕ between two interfering beams, and OAM control occurs through projection and the filtering process. Through these steps, we can check for the desired phenomenon. To clear the information about the polarization of the interfering beams, we include another polarized beam splitter (PBS₂) before SLM₃. In practice, the HWP rotates the polarization of the projection. The data showing two-photon bunching are presented in Fig. 5. Here, the red dashed curve is theoretical prediction, and the solid blue curve is the best fit to the experimental data. Each error bar shows one standard deviation from the mean at each experimental data point. Furthermore, we aim to present information regarding the photon rates associated with the experiment. In both previous work,^{28,29} the photon statistics were at a very low rate, which means we explored up to 4×10^5 Hz photon rate in the first project and up to 2×10^5 Hz for the second one. In the current experiment for the single-detection rate of 4×10^5 Hz, we have approximately 2200 overall coincidence counts (joint detections) per second when both arms are open, prior to subtracting unwanted states. After the subtraction of the unwanted states, we have approximately 1300 coincidence counts per second as the absolute (post-selected coincidence) counts. The data

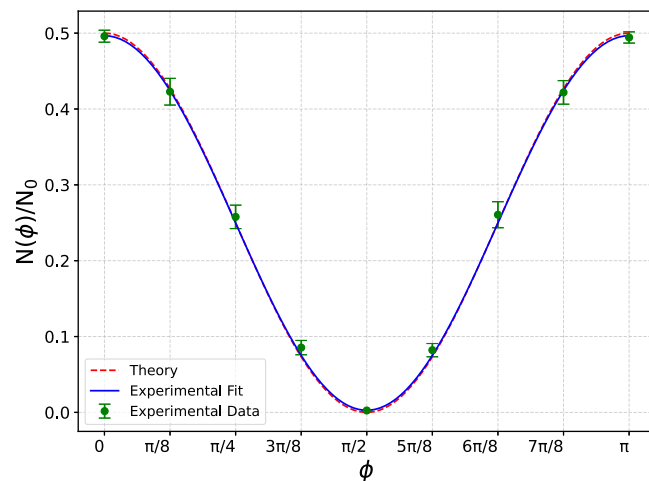


FIG. 5. Experimental realization of two-photon bunching that shows normalized counts versus polarization mismatched angle ϕ in $\pi/8$ steps. Here, the red dashed curve represents the theory [Eq. (20)], and the solid blue curve shows the fit to the experimental data points.

acquisition period is 5 s and over the coincidence window of 10 ns. In the current work, we have pushed the single rate to the limit of our detector at 1.5×10^6 photons per second for a long usage, resulting in the double detection about 26 000/s, and the absolute counts regarding the pure contributions of the single-photons around 11 000/s. Also, it is possible to use our detector for a short time and go up to 4 MHz. At this limit, we have approximately 86 000/s absolute counts. This happens at the maxima of our result corresponding to 0° or 180° . However, with the aforementioned source, we can enhance the single rate up to 30 MHz with just 2% error in Eq. (14). In fact, this is a significant increment of the photon rate after adding OAM. The explanation of this observation is really simple. The new degree of freedom, i.e., OAM, removes much classical noise from the signal as seen in Fig. 4. In the case of having a source with a shorter coherence time, it is possible to further increase this rate, potentially surpassing the current photon rate of commercially available single-photon detectors. Therefore, we think that this approach has the potential to advance our current capabilities and opens up new opportunities for exploring quantum phenomena. In Fig. 6, we have plotted the absolute coincidence counts versus the average single-detection rate in the unit of kHz. Our intuition is that this increment in the absolute coincidence counts will be saturated as we increase the photon rate, and this occurs when we have the contribution from higher-order terms. At that limit, a CW source no longer can be used to extract the single-photon contribution. This is the only fundamental limitation of this approach, and as stated, by having a source with the shorter coherence time, one can enhance the single and double photodetections rate. At this point, we would like to address this idea that in this approach the post-selected photon rate in producing a quantum signature can set an upper bound on the single-photon generation based on the input power. In the experiment, we have more than 85 kHz absolute coincidence counts only from a few nanowatt of power delivered to the interferometer with three diffractive SLMs. As mentioned earlier, it is also possible to reduce this initial power even further with more efficient configuration of these SLMs (at least by one order). To the best of our knowledge, currently no single-photon pair based on non-linearity can be practically produced from 1 nW of pump power.

V. CONCLUSION

In conclusion, what we have observed with quasi-classical states carrying OAM is the experimental realization of the quantum behavior of such states, reproducing one of the basic phenomena in the domain

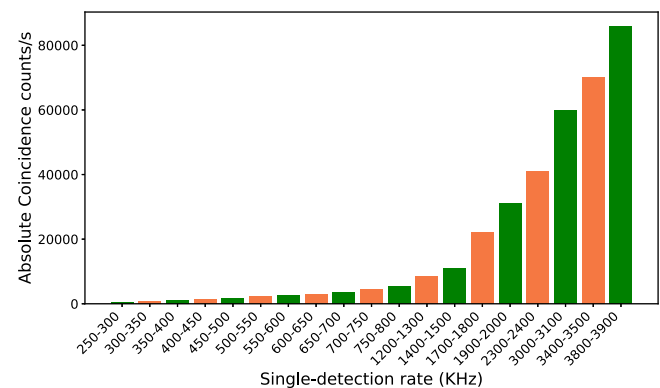


FIG. 6. Absolute coincidence counts as a function of single-detection rate.

of quantum optics, namely, two-photon bunching with a considerable coincidence rate and the possibility to enhance it further. Based on the input power, we think that this source sets an upper bound for any single-photon source simply because there is no interaction with matter. We have also experimentally observed the existence of the number state-phase uncertainty, even though there is currently no unitary operator representing the phase. The standard extension of this toolbox is that the PRWCSs can serve to investigate manifold quantum aspects of light pertinent to modern quantum technologies as well as to search for a deeper understanding of fundamental physics in the settings compatible with the process of post-selection, which is needed to isolate the single-photon contribution of coherent states in the weak limit.

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AUTHOR DECLARATIONS

Conflict of Interest

The author has no conflicts to disclose.

Author Contributions

Moslem MahdaviFar: Conceptualization (equal); Formal analysis (equal); Investigation (equal); Methodology (equal); Writing – original draft (equal); Writing – review & editing (equal).

DATA AVAILABILITY

Data underlying the results presented in this work are not publicly available at this time but may be obtained from the author upon reasonable request.

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