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Berry Phase and Statistical Properties of the Displaced Fock and Squeezed States

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Abstract: We study the displaced Fock states and squeezed displaced states by calculating their Berry phase when their parameters are explicitly time dependent. Furthermore, we analyze their statistical properties to identify the parameter regimes associated with classical and non-classical behavior. To this end, we assume that both the displaced Fock states and the squeezed displaced states are an explicit function of time where their parameters vary in time to obtain the geometrical phase. Then, we calculate the Mandel Q-parameter and the second-order correlation function $g^{(2)}(0)$ for each state using the similarity transformations of the ladder operators that allow us to obtain the expectation values needed.

Keywords: Displaced Fock states, Squeezed displaced states, Berry phase, Mandel Q-parameter, Second-order correlation function, Classical and non-classical behavior.

1. Introduction

In quantum optics, there is a growing interest in non-classical states of light. The harmonic oscillator is a key tool for these studies. By quantizing the electromagnetic field, we can map different optical systems to a harmonic oscillator [1-5]. This allows us to explore fundamental concepts in quantum mechanics.

We can define the states of light as eigenstates of Hamiltonians in a Fock space. Some of these states show classical properties, while others show non-classical properties [6]. Studying the quantum features of these states is very important in optics [2-6].

Another useful property of these quantum states is the Berry phase. This phase appears in many physical systems [3, 6, 7]. It is a geometric phase, it is a gauge invariant, and it connects closely with differential geometry [3, 6].

Furthermore, the study of these non-classical states of light is relevant for quantum metrology. Specifically, the reduction of photon number fluctuations helps to improve the signal-to-noise ratio in optical measurements because the noise is below the classical shot-noise limit. Therefore, understanding their geometrical phases and statistical properties provides a theoretical basis for future quantum transducers. The exact implementation of these states depends on the specific sensor architecture, which is beyond the scope of this fundamental analysis.

The aim of this work is to extend the study of the displaced Fock states and squeezed displaced states. The statistical results, including the Mandel Q-parameter and the second-order correlation function $g^{(2)}(0)$, were previously reported in a preliminary form in ref. [2]. The present article extends that conference paper by calculating the Berry phase when the state parameters are explicitly time-dependent.

While we have previously used the $su(1,1)$ algebra to calculate the Berry phase for the parametric amplifier [3], this manuscript applies the method to a different physical system. Furthermore, the geometric phase for generalized squeezed coherent states was established in ref. [6]. In Section 2, we derive our equations as special cases of their general expression. The new physical contribution of our specific treatment is that we use the algebraic approach to connect the exact geometric phase directly to the parameter boundaries that define classical and non-classical statistical behavior.

This work is organized as follows. In Section 2, we compute the Berry phase of both states. Section 3 reproduces the statistical properties from our previous work to identify the regimes where non-classical behavior appears. Finally, some concluding remarks are given.

2. Berry Phase

The geometric phase appears when we study a parameterized family of quantum states. We assume these states are written in terms of a set of parameters which are time-dependent and make up a closed loop C in the parameter space. After completing the cyclic evolution, the state returns to itself up to a phase factor. The geometric phase is defined as

$$\gamma_n(t) = i \oint_C \langle \varphi_n | \nabla_R | \varphi_n \rangle \, dR \tag{1}$$

This line integral, in general, is different to zero and depends on the path instead of how quickly it is traversed [3, 7, 8].

2.1. Berry Phase of the Displaced Fock States

Since we are mapping the states of light to eigenstates of Hamiltonians which can be defined in an appropriate Fock space, we need to use its basis states called Fock states, which are states with specific number of energy quanta, defined as [1, 2, 9, 10]

$$\hat{a}^\dagger \hat{a} |n\rangle = n|n\rangle, \tag{2}$$

where the number operator is $\hat{n} = \hat{a}^\dagger \hat{a}$ and the states $|n\rangle$ are the Fock states with $n = 0,1,2, \ldots$. Furthermore, the n photon state is defined as

$$|n\rangle = \frac{(\hat{a}^\dagger)^n}{\sqrt{n!}} |0\rangle, \tag{3}$$

with a factor that leads to the correct normalization of the state [1, 2, 10].

In particular, the harmonic oscillator states which minimizes the Heisenberg uncertainty relation, $\Delta p \Delta q \geq \hbar/2$, plays a significant role in the quantum theory of radiation and in quantum optics [2, 10-12].

These states which simultaneously minimize the variance in both the position and momentum operators, are called coherent states [2, 10, 11]. In [2-5, 12], the coherent states are defined as the action of the displacement operator on the vacuum state

$$|\alpha\rangle = \hat{D}(\alpha)|0\rangle, \tag{4}$$

where $\hat{D}(\alpha) = e^{\alpha \hat{a}^\dagger - \alpha^* \hat{a}}$ and $\alpha = |\alpha| e^{i\phi}$.

To compute the Berry phase, we assume that the parameter of the displaced Fock states is time dependent

$$|\alpha(R(t))\rangle = \hat{D}(\alpha(R(t)))|n\rangle, \tag{5}$$

where $\hat{D}(\alpha(R(t))) = e^{\alpha(R(t))\, \hat{a}^\dagger - \alpha^*(R(t))\, \hat{a}}$ and $\alpha\big(R(t)\big) = \alpha\big(\phi(t)\big) = |\alpha| e^{i\phi(t)}$. So, the geometrical phase is given by

$$\gamma = i \oint_C \langle \alpha\big(R(t)\big) | \nabla | \alpha\big(R(t)\big) \rangle \, dR, \tag{6}$$

Using Eq. (5), we obtain the following expression

$$\gamma = i \oint_C \langle n | \hat{D}^\dagger\big(\alpha(R)\big) \frac{d}{dR} \hat{D}\big(\alpha(R)\big) | n \rangle \, dR \tag{7}$$

Using the Hadamard identity which is defined as

$$\begin{aligned} e^A B e^{-A} = B + [A,B] + \tfrac{1}{2!}\big[A,[A,B]\big] + \\ + \tfrac{1}{3!}\Big[A,\big[A,[A,B]\big]\Big] + \cdots, \end{aligned} \tag{8}$$

we can calculate the similarity transformation of the operator $\frac{d}{dR}$ as follows

$$\begin{aligned} e^{\alpha^* \hat{a} - \alpha \hat{a}^\dagger} \frac{d}{dR} e^{\alpha \hat{a}^\dagger - \alpha^* \hat{a}} = \\ = \frac{d}{dR} - \dot{\alpha}^* \hat{a} + \dot{\alpha} \hat{a}^\dagger + \tfrac{1}{2}(\alpha^* \dot{\alpha} - \alpha \dot{\alpha}^*) \end{aligned}$$

Therefore, by substituting we obtain that the Berry phase is given as

$$\gamma = \tfrac{i}{2} \oint_C (\alpha^* \dot{\alpha} - \alpha \dot{\alpha}^*) \, dR, \tag{9}$$

where $\alpha^* \dot{\alpha} - \alpha \dot{\alpha}^* = 2i|\alpha|^2 \dot{\phi}$ and the displacement parameter satisfies ϕ: $[0,2\pi]$

$$\gamma = \tfrac{i}{2} \oint_0^{2\pi} 2i \, |\alpha|^2 d\phi = -2\pi |\alpha|^2 \tag{10}$$

Therefore, the Berry phase has been calculated. This result is a special case of the general expression presented in ref. [6], obtained here specifically for the displaced Fock states using the translation properties of the operators.

2.2. Berry Phase of the Squeezed Displaced States

Another harmonic oscillator state which minimizes the Heisenberg uncertainty relation are the squeezed states where the standard deviation of one observable

is squeezed, and the other conjugate observable is stretched exhibiting properties that closely resemble the classical harmonic oscillator [2, 10-12]. These states are defined by the combined action of the displacement and squeezing operators on the vacuum state

$$|\alpha,\xi\rangle = \hat{D}(\alpha)\hat{S}(\xi)|0\rangle, \tag{11}$$

where $\hat{D}(\alpha) = e^{\alpha\hat{a}^\dagger - \alpha^*\hat{a}}$ and $\hat{S}(\xi) = e^{\frac{1}{2}\left(\xi^*\hat{a}^2 - \xi\hat{a}^{\dagger 2}\right)}$ with $\xi = re^{i\theta}$ which satisfies $0 \le r < \infty$ and $0 \le \theta \le 2\pi$ [2-5,10-14].

Following a similar procedure to calculate the Berry phase of the displaced squeezed states, we need to assume that their parameters are time dependent

$$\begin{gathered}|\alpha(R(t)),\xi(R(t))\rangle = \\ = \hat{D}(\alpha(R(t)))\hat{S}(\xi\big(R(t)\big))|0\rangle,\end{gathered} \tag{12}$$

where $\hat{D}(\alpha(R(t))) = e^{\alpha(R(t))\,\hat{a}^\dagger - \alpha^*(R(t))\,\hat{a}}$ and $\hat{S}(\xi(R(t))) = e^{\frac{1}{2}(\xi^*(R(t))\hat{a}^2 - \xi\left(R(t)\right)\hat{a}^{\dagger 2})}$. So, the geometrical phase is given by

$$\gamma = i\oint_C \langle\alpha(R),\xi(R)|\nabla|\alpha(R),\xi(R)\rangle\, dR \tag{13}$$

Using Eq. (12), we obtain the following expression

$$\gamma = i\oint_C \langle 0|\,\hat{S}^\dagger(\xi)\hat{D}^\dagger(\alpha)\frac{d}{dR}\hat{D}(\alpha)\hat{S}(\xi)|0\rangle\, dR \tag{14}$$

Considering the generators of the $su(1,1)$ lie algebra defined as [3-5, 12]

$$\begin{gathered}K_0 = \frac{1}{2}\left(\hat{a}^\dagger\hat{a} + \frac{1}{2}\right), K_+ = \frac{1}{2}\hat{a}^{\dagger 2}, \\ K_- = \frac{1}{2}\hat{a}^2,\end{gathered} \tag{15}$$

which satisfy $[K_0, K_\pm] = \pm K_\pm$ and $[K_-, K_+] = 2K_0$, we simplify the calculation of the similarity transformations

$$\begin{gathered}e^{\xi K_+ - \xi^* K_-}\frac{d}{dR}e^{\xi^* K_- - \xi K_+} = \\ = \frac{d}{dR} - \left(\dot{\xi}K_+ - \dot{\xi}^* K_-\right) + \left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right)K_0 - \\ -\frac{1}{3}\left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right)(\xi K_+ + \xi^* K_-) + \\ + \frac{1}{3}\left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right)|\xi|^2 K_0 + \\ -\frac{1}{15}\left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right)|\xi|^2(\xi K_+ + \xi^* K_-) + \\ + \frac{2}{45}\left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right)|\xi|^4 K_0 - \cdots,\end{gathered}$$

$$\begin{gathered}e^{\frac{1}{2}(\xi\hat{a}^{\dagger 2} - \xi^*\hat{a}^2)}(-\dot{\alpha}^*\hat{a} + \dot{\alpha}\hat{a}^\dagger)\, e^{\frac{1}{2}(\xi^*\hat{a}^2 - \xi\hat{a}^{\dagger 2})} = \\ = -\xi\dot{\alpha}^*\hat{a}^\dagger + \xi^*\dot{\alpha}\hat{a} + |\xi|^2\dot{\alpha}^*\hat{a} - |\xi|^2\dot{\alpha}\hat{a}^\dagger - \\ - |\xi|^2\xi\dot{\alpha}^*\hat{a}^\dagger + |\xi|^2\xi^*\dot{\alpha}\hat{a} + \cdots,\end{gathered}$$

$$\begin{gathered}e^{\frac{1}{2}\left(\xi\hat{a}^{\dagger 2} - \xi^*\hat{a}^2\right)}\left(\frac{1}{2}(\alpha^*\dot{\alpha} - \right. \\ \left. -\alpha\dot{\alpha}^*)\right)\, e^{\frac{1}{2}(\xi^*\hat{a}^2 - \xi\hat{a}^{\dagger 2})} = \frac{1}{2}(\alpha^*\dot{\alpha} - \alpha\dot{\alpha}^*)\end{gathered}$$

Since $\xi = |r|e^{i\theta}$, therefore $\left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right) = 2ir^2 d\theta$ and the similarity transformation can be written as

$$\begin{gathered}e^{\xi K_+ - \xi^* K_-}\frac{d}{dR}e^{\xi^* K_- - \xi K_+} = \\ \frac{d}{dR} + 2i\, sinh^2(r)\, d\theta K_0 - \left(\dot{\xi}K_+ - \dot{\xi}^* K_-\right) - \\ -\frac{1}{3}\left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right)(\xi K_+ + \xi^* K_-) - \\ -\frac{1}{15}\left(\xi^*\dot{\xi} - \xi\dot{\xi}^*\right)|\xi|^2(\xi K_+ + \xi^* K_-) - \cdots\end{gathered}$$

Now, using the action of these operators

$$\begin{gathered}K_+|k,n\rangle = \sqrt{(n+1)(2k+n)}|k,n+1\rangle, \\ K_-|k,n\rangle = \sqrt{n(2k+n-1)}|k,n-1\rangle, \\ K_0|k,n\rangle = (k+n)|k,n\rangle,\end{gathered}$$

the Berry phase is given as

$$\gamma = \mathrm{i}\oint_C \left(\frac{1}{2}(\alpha^*\dot{\alpha} - \alpha\dot{\alpha}^*) + \frac{1}{2}i\, sinh^2(r)\right)d\theta, \tag{16}$$

where the first term is the same as the Berry phase of the displaced Fock states. For the closed path C, we assume that the phase parameters ϕ and θ vary simultaneously from 0 to 2π. Thus, integrating over this cycle gives

$$\begin{gathered}\gamma = \frac{i}{2}\oint_0^{2\pi} 2i\,|\alpha|^2 d\theta + \frac{i}{2}\oint_0^{2\pi} i\, sinh^2(r)d\theta = \\ = -2\pi|\alpha|^2 - \pi sinh^2(r)\end{gathered} \tag{17}$$

The Berry phase has been calculated for this system. This equation is a special case of the general geometric phase established in [6], derived here through the explicit time-dependent evolution of the $su(1,1)$ generators.

3. Statistical Properties

Depending on the presence of some quantum features such as the relative fluctuation in photon number, the quantum states of light can be classified as classical and non-classical states [1, 2, 4].

Specifically, the Mandel Q-parameter allows us to measure the deviation from the Poisson distribution to distinguish quantum processes from classical ones because the photon number distribution for a coherent field is Poissonian, so, any distribution which is narrower than Poissonian must correspond to a non-classical field [1, 2, 4, 13]. This parameter is defined as

$$Q = \frac{\langle\hat{n}^2\rangle - \langle\hat{n}\rangle^2}{\langle\hat{n}\rangle} - 1 = \frac{\langle(\Delta n)^2\rangle - \langle\hat{n}\rangle}{\langle\hat{n}\rangle}, \tag{18}$$

where it describes a super-Poissonian distribution for $Q > 0$, a Poissonian distribution for $Q = 0$, and a sub-Poissonian distribution for $Q < 0$ [1, 2, 4, 13].

The second-order correlation function $g^2(0)$ quantifies the intensity fluctuations and provides

information on the bunching or the anti-bunching effects, is defined as

$$g^2(0) = \frac{\langle \hat{n}^2 \rangle - \langle \hat{n} \rangle}{\langle \hat{n} \rangle^2} = \frac{Q}{\langle \hat{n} \rangle} + 1, \quad (19)$$

where the bunching effect appears when $g^2(0) > 1$ and the anti-bunching effect when $g^2(0) < 1$ [1, 2, 4, 13-15].

3.1. Statistical Properties of the Displaced Fock State

By applying the Hadamard identity from Eq. (8), we find the translation properties for the annihilation and creation operators [2-5, 12-14]

$$\hat{D}^\dagger(\alpha)\hat{a}\hat{D}(\alpha) = \hat{a} + \alpha, \quad (20)$$

$$\hat{D}^\dagger(\alpha)\hat{a}^\dagger\hat{D}(\alpha) = \hat{a}^\dagger + \alpha^* \quad (21)$$

Next, we calculate the required expectation values to find the Mandel Q-parameter. We use the standard relations [2, 12]

$$\hat{a}|n\rangle = \sqrt{n}|n-1\rangle, \quad (22)$$

$$\hat{a}^\dagger|n\rangle = \sqrt{n+1}|n+1\rangle \quad (23)$$

Consequently, the expectation values are

$$\langle \hat{a}^\dagger\hat{a} \rangle = n + |\alpha|^2, \quad (24)$$

$$\langle (\hat{a}^\dagger\hat{a})^2 \rangle = n^2 + 4|\alpha|^2 n + |\alpha|^4 + |\alpha|^2 \quad (25)$$

This gives the following Mandel Q-parameter

$$Q = \frac{n(2|\alpha|^2-1)}{n+|\alpha|^2} \quad (26)$$

As illustrated in Fig. 1, a sub-Poissonian distribution occurs when $|\alpha|^2 < 1/2$. For larger values, we observe a super-Poissonian distribution. This transition applies strictly for $n > 0$. For the special case of $n = 0$, we obtain $Q = 0$ for all values of α. This is expected because the displaced vacuum state is exactly a coherent state. Fig. 2 displays the Mandel Q-parameter depending on the quantum number n and the displacement α.

Following this, we determine the second-order correlation function $g^2(0)$

$$g^2(0) = \frac{n(2|\alpha|^2-1)}{(n+|\alpha|^2)^2} + 1 \quad (27)$$

Antibunching effects are present when $|\alpha|^2 < 1/2$. This indicates that small displacements preserve the non-classical nature of the Fock state. However, for large displacements, classical behavior takes over. Fig. 3 shows that the exact transition happens at $|\alpha|^2 = 1/2$, where the system acts like a coherent state.

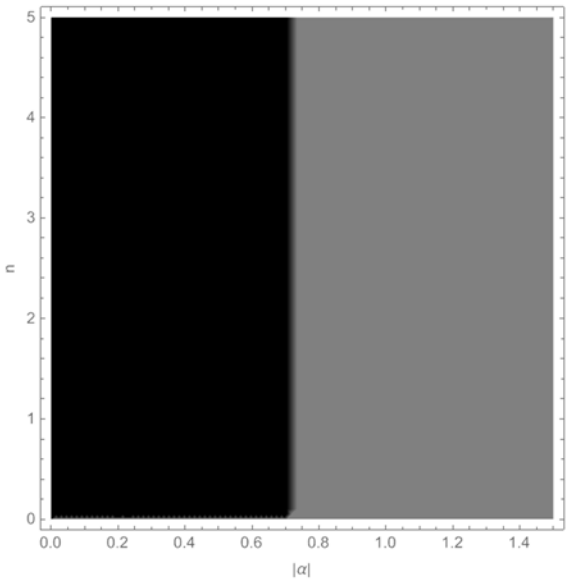


Fig. 1. Phase diagram of the Mandel Q-parameter for the displaced Fock state. Black region: $Q < 0$. Gray region: $Q > 0$.

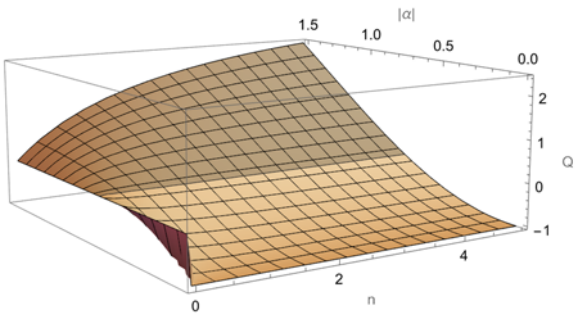


Fig. 2. Mandel Q-parameter for the displaced Fock state.

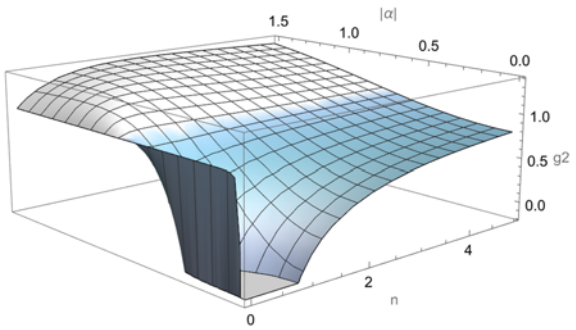


Fig. 3. Second-order correlation function $g^2(0)$ for the displaced Fock state.

3.2. Statistical Properties of the Squeezed Displaced State

With the help of Eqs. (8), (20) and (21), we derive the properties of the operators for this state [2, 12-14]

$$\hat{S}^{\dagger}(\xi)\hat{D}^{\dagger}(\alpha)\hat{a}\hat{D}(\alpha)\hat{S}(\xi) = \mu\hat{a} - v\hat{a}^{\dagger} + \alpha, \quad (28)$$

$$\hat{S}^{\dagger}(\xi)\hat{D}^{\dagger}(\alpha)\hat{a}^{\dagger}\hat{D}(\alpha)\hat{S}(\xi) = \mu\hat{a}^{\dagger} - v^{*}\hat{a} + \alpha^{*}, \quad (29)$$

where $\mu = \cosh(r)$ and $v = e^{i\theta}\sinh(r)$. Using these, the expectation values needed for the Mandel Q-parameter are

$$\langle \hat{a}^{\dagger}\hat{a} \rangle = |\alpha|^2 + sinh^2(r), \quad (30)$$

$$\langle (\Delta n)^2 \rangle = |\alpha|^2 \begin{pmatrix} e^{-2r}cos^2\left(\phi - \frac{\theta}{2}\right) + \\ +e^{2r}sin^2\left(\phi - \frac{\theta}{2}\right) \end{pmatrix} + \\ +2sinh^2(r)(sinh^2(r) + 1) \quad (31)$$

We define $\alpha = |\alpha|e^{i\phi}$. Then, the Mandel Q-parameter is

$$Q = \frac{|\alpha|^2(e^{-2r}-1)+2sinh^4(r)+sinh^2(r)}{|\alpha|^2+sinh^2(r)} \quad (32)$$

Here, we set $\phi = \theta/2$. This aligns the displacement with the squeezed quadrature to minimize the photon number fluctuations [2, 14].

In this case, it is possible to obtain both sub- and super-Poissonian distribution depending on the values of r and α, as shown in Fig. 4 and Fig. 5 [14, 15]. The exact boundary between these two regimes is determined analytically when the numerator of Eq. (32) is equal to zero: $|\alpha|^2(e^{-2r} - 1) + 2sinh^4(r) + sinh^2(r) = 0$. Supported by this condition, the super-Poissonian regime dominates for larger values of r. In contrast, the sub-Poissonian distribution appears for small values of r and $|\alpha|^2 \neq 0$.

Finally, we compute the second-order correlation function $g^2(0)$

$$g^2(0) = \frac{|\alpha|^2(e^{-2r}-1)+2sinh^4(r)+sinh^2(r)}{(|\alpha|^2+sinh^2(r))^2} + 1 \quad (33)$$

Depending on r and α, this state exhibits either bunching or antibunching effects, as seen in Fig. 6 [2, 13]. Bunching occurs exclusively when r is large. Conversely, antibunching is present for small r values if $|\alpha|^2 \neq 0$. This shows that adding a small amount of squeezing generates non-classical behavior.

4. Conclusions

In this article, we studied the geometric phase and the statistical properties of both displaced Fock states and squeezed displaced states. By assuming that the state parameters change over time and complete a closed path, we successfully calculated the exact Berry phase for both systems.

Through the Mandel Q-parameter and the second-order correlation function, we mapped the regions of classical and non-classical behavior. For the displaced Fock states, we found that non-classical features like sub-Poissonian statistics and antibunching dominate when $|\alpha|^2 < 1/2$. The transition to classical coherent behavior happens exactly at $|\alpha|^2 = 1/2$ for $n > 0$.

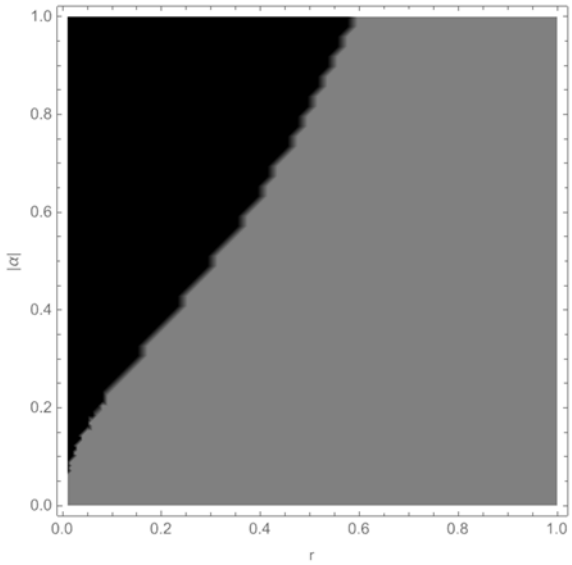


Fig. 4. Phase diagram of the Mandel Q-parameter for the squeezed displaced state. Blue region: $Q < 0$. Red region: $Q > 0$.

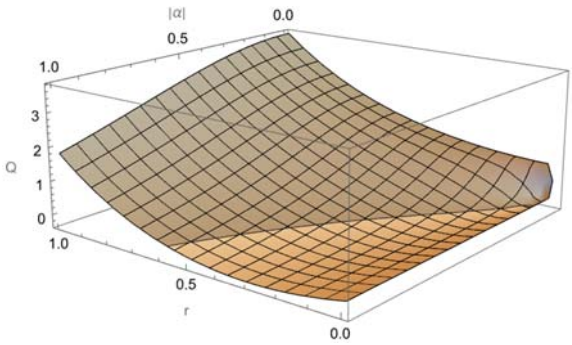


Fig. 5. Mandel Q-parameter for the squeezed displaced state.

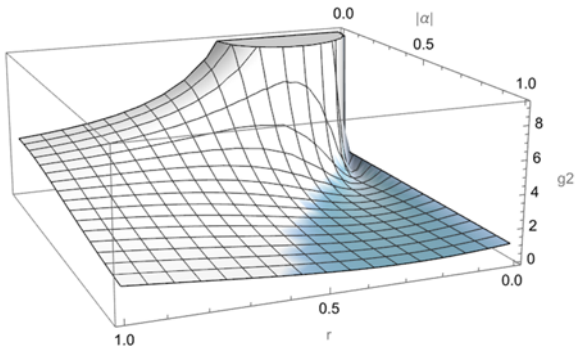


Fig. 6. Second-order correlation function $g^2(0)$ for the squeezed displaced state.

For the squeezed displaced states, we minimized the photon number fluctuations by setting $\phi = \theta/2$. Our analysis shows that this state can present both

sub-Poissonian and super-Poissonian distributions, as well as bunching and antibunching effects. A key finding is that a small squeezing parameter r is enough to produce non-classical effects, provided that the displacement is not zero.

The non-classical effects identified in this work, such as sub-Poissonian distributions and antibunching, are resources for quantum metrology. By finding the parameter regimes where the quantum noise is below the standard limit, our results show a theoretical advantage for improving the signal-to-noise ratio. The specific experimental implementation of these states depends on the transducer architecture, which is left for future applied research.

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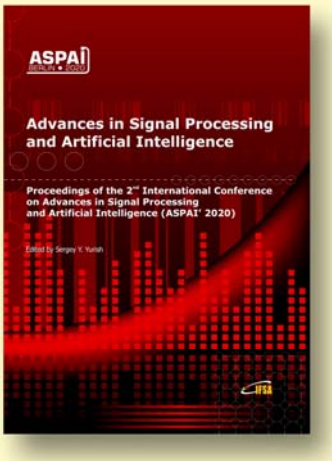
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