

Interplay of nonclassicality and entanglement of two-mode Gaussian fields generated in optical parametric processes

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The behavior of general nonclassical two-mode Gaussian states at a beam splitter is investigated. Single-mode nonclassicality as well as two-mode entanglement of both input and output states are analyzed suggesting their suitable quantifiers. These quantifiers are derived from local and global invariants of linear unitary two-mode transformations such that the sum of input (or output) local nonclassicality measures and entanglement measure gives a global invariant. This invariant quantifies the global nonclassicality resource. Mutual transformations of local nonclassicalities and entanglement induced by the beam splitter are analyzed considering incident noisy twin beams, single-mode noisy squeezed vacuum states, and states encompassing both squeezed states and twin beams. A rich tapestry of interesting nonclassical output states is predicted.

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I. INTRODUCTION

The nonclassical properties of light have been for a long time the main topic of interest in quantum optics. The question whether a given quantum state is nonclassical (i.e., cannot be treated by the classical statistical theory) has been considered as one of the most important problems since the early days of quantum physics [1–3] (for a review see, e.g., Refs. [4–6]). For optical fields, a commonly accepted criterion for distinguishing nonclassical states from the classical ones is expressed as follows [5,7–9]: a quantum state is nonclassical if its Glauber-Sudarshan P function fails to have all the properties of a probability density. We recall that the Glauber-Sudarshan P function for an M -mode bosonic state $\hat{\rho}$ can be defined as [10,11]

$$\hat{\rho} = \int P(\boldsymbol{\alpha}, \boldsymbol{\alpha}^*) |\boldsymbol{\alpha}\rangle \langle \boldsymbol{\alpha}| d^2\boldsymbol{\alpha}, \quad (1)$$

where $|\boldsymbol{\alpha}\rangle = \prod_{m=1}^M |\alpha_m\rangle$ is given in terms of the m th-mode coherent state $|\alpha_m\rangle$, which is the eigenstate of the m th-mode annihilation operator \hat{a}_m , $\boldsymbol{\alpha}$ denotes complex multivariable $(\alpha_1, \alpha_2, \dots, \alpha_M)$, and $d^2\boldsymbol{\alpha} = \prod_m d^2\alpha_m$. It is worth noting that the negativity of the P function is necessary and sufficient for nonclassicality, while the singularity or irregularity of the P function is only a sufficient condition (i.e., it is a nonclassical witness). Thus, if the P function is more singular or more

irregular than Dirac's δ function for a given state, then it is also nonpositive (semidefinite) in the formalism of generalized functions. A standard example of such irregular functions is the P function for an n -photon Fock state (with $n = 1, 2, \dots$), which is given by the n th derivative of $\delta(\boldsymbol{\alpha})$.

Based on this definition of nonclassicality, various operational criteria (also called witnesses) have been described for testing the nonclassicality of single-mode [7,8,12] and multimode [13–15] fields. Their derivations are based either on the fields moments [13,15,16] or exploit the Bochner theorem written for the characteristic function of the Glauber-Sudarshan P function [17]. A direct reconstruction of the quasidistributions of integrated intensities is a sufficient but not necessary condition of the nonclassicality of the detected fields by the definition [18–20]. We note that nonclassicality criteria derived from the majorization theory have also been found useful [21,22].

Entanglement between two optical fields is one of the most frequently studied forms of nonclassical light. Such light emerges in various two-mode or multimode nonlinear optical processes, e.g., in spontaneous parametric down-conversion. In this process, pairs of photons composed of the signal and idler modes are created at the expense of the annihilated pump photons. This pairwise character of emitted light lies in the heart of entanglement here. The process of spontaneous parametric down-conversion has its degenerate variant called second-subharmonic generation, where both photons in a pair are emitted into the same optical mode. This gives rise to phase squeezing of the second-subharmonic field composed

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of, in general, many photon pairs. The squeezed light is also considered nonclassical as it has its phase fluctuations suppressed below the classical limit. The nonclassicality in both cases has the same origin which is pairing of photons. On the other side, the emitted photon pairs can be manipulated by linear optics. In detail, two photons from one pair present in the same mode of a squeezed state of light can be split (on a beam splitter) and contribute to the entanglement of the output fields. Also, two photons from a pair incident on different input ports of a beam splitter can “stick together” (bunch) and leave the beam splitter in the same output port (as testified in the Hong-Ou-Mandel interferometer [23]). The interconnection of these two types of fields by the means of linear optics has already been shown by Braunstein [24] and later elaborated by Adesso [25] for arbitrarily strong Gaussian states. This behavior poses a natural question whether it is possible to introduce a physical quantity that quantifies “a nonclassicality resource” present during the creation of both types of fields and later conserved during linear-optical transformations.

The answer to this question is intimately related to the quantifiers of *entanglement* and *local nonclassicality*. Several measures were proposed to quantify the entanglement in both discrete and continuous domains [26–31]. The negativity (or its logarithmic variant) is considered, probably, as the most useful at present. On the other hand, the Lee nonclassicality depth [21] is conventionally used to quantify the nonclassicality of optical fields. Alternatively, the nonclassicality of an optical field can be transcribed to entanglement using a beam splitter and quantified via an entanglement measure [32,33]. For a comparative study of these two nonclassicality measures, see, e.g., recent Refs. [34,35].

We note that, apart from the local nonclassicalities of two parts of a bipartite state, also *global nonclassicality* can naturally be defined. All these three quantities have been analyzed in Ref. [36] for intense multimode twin beams with the following result: whenever a twin beam is entangled, it is globally nonclassical. On the other hand, its signal and idler constituents are multithermal and so locally classical. A general approach for describing the relation between the entanglement and global nonclassicality of two-mode states has been proposed in Ref. [37].

Returning back to our question, we look for an invariant with respect to linear-unitary transformations (conserving the overall number of photons) that comprises both the entanglement and local nonclassicalities. This question has recently been addressed in Ref. [38] considering beam-splitter transformations and a quantity composed of the logarithmic negativity and the logarithm of nonclassicality depth. However, the introduced quantity has been found useful only under very specific conditions [39].

In this paper, we construct such an invariant for general two-mode Gaussian states arising in nonlinear processes described by the second-order susceptibility $\chi^{(2)}$. The processes of spontaneous parametric down-conversion and second-subharmonic generation represent their most important examples. As schematically shown in Fig. 1, the found invariant is decomposable into three parts characterizing in turn entanglement and two local nonclassicalities. The entanglement indicator is shown to be a monotone of the logarithmic negativity similarly to the newly defined nonclassicality measure that is

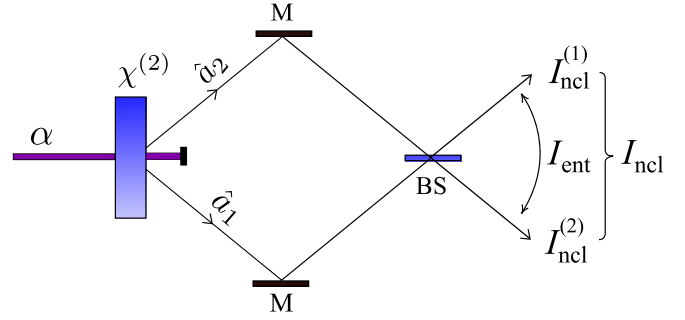


FIG. 1. Diagram showing the main goal of this paper: The local ($I_{\text{ncl}}^{(1)}$ and $I_{\text{ncl}}^{(2)}$) and global (I_{ncl}) nonclassicality invariants are analyzed in relation with the entanglement, described by the invariant I_{ent} , for the light generated by the optical parametric process (described by the second-order susceptibility $\chi^{(2)}$) and then combined at a beam splitter BS with varying transmissivity T . Here, α is the amplitude of a classical pump field, \hat{a}_1 and \hat{a}_2 are the annihilation operators of the generated light, and M denotes a mirror.

a monotone of the Lee nonclassicality depth under any linear unitary transformation.

The obtained results are potentially interesting for manipulations with nonclassicality in quantum engineering that have become substantial ingredients of a growing number of applications of quantum technologies [28,40–43].

The paper is organized as follows. In Sec. II, a model comprising parametric down-conversion and second-subharmonic generation is developed. A suitable nonclassicality invariant is suggested using local and global invariants of two-mode Gaussian fields. Its decomposition into an entanglement quantifier and local nonclassicality quantifiers is also discussed. Twin beams as they behave on a beam splitter are discussed in Sec. III. In Sec. IV, a single-mode squeezed state on a beam splitter is analyzed. Section V is devoted to the behavior of two single-mode squeezed states interfering on a beam splitter. States having both “twin-beam” and squeezed components are investigated in Sec. VI. Conclusions are drawn in Sec. VII. Quasidistributions related to the normal and symmetric ordering of operators are discussed in the Appendix.

II. GAUSSIAN STATES GENERATED IN $\chi^{(2)}$ INTERACTIONS AND THEIR INVARIANTS

We consider a nonlinear interaction Hamiltonian \hat{H}_{int} describing both parametric down-conversion and second-subharmonic generation that provide photon pairs [9] (for the scheme, see Fig. 2)

$$\hat{H}_{\text{int}} = -\hbar g_{12}^* \hat{a}_1 \hat{a}_2 - \hbar g_{11}^* \hat{a}_1^2 - \hbar g_{22}^* \hat{a}_2^2 + \text{H.c.} \quad (2)$$

In Eq. (2), the symbols \hat{a}_1 (\hat{a}_1^\dagger) and \hat{a}_2 (\hat{a}_2^\dagger) represent the annihilation (creation) operators of the fields 1 and 2, g_{12} is a nonlinear coupling constant characterizing parametric down-conversion, and g_{ii} stands for a nonlinear coupling constant of the second-subharmonic generation in the i th mode described by the second-order susceptibility $\chi^{(2)}$ of a medium. Symbol H.c. represents the Hermitian conjugated terms. Due to the presence of noise in real nonlinear processes, we also

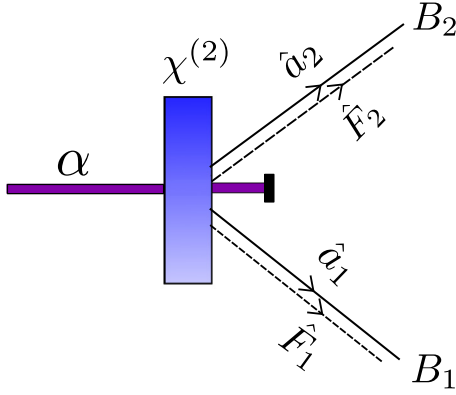


FIG. 2. Diagram of the optical parametric process described by Eq. (2): the classical pump field, with complex amplitude α , generates the signal and idler modes described by the annihilation operators \hat{a}_j and affected by the noise stochastic operators \hat{F}_j , $j = 1, 2$. For simplicity, the pump-field amplitude α is incorporated into the coupling constants g_{ij} . The mean photon number in the signal (idler) mode influenced by the noise is denoted by B_1 (B_2). In Sec. III, $B_1 = B_p + B_s$ and $B_2 = B_p + B_i$, where $B_p = \sinh^2(g_{12}t)$ is the mean number of generated photon pairs and $B_s = \langle \hat{F}_1^\dagger \hat{F}_1 \rangle$ ($B_i = \langle \hat{F}_2^\dagger \hat{F}_2 \rangle$) is the mean number of signal (idler) noise photons. In Secs. IV and V, $B_1 = \bar{B}_p^s + B_s$ and $B_2 = \bar{B}_p^i + B_i$, where \bar{B}_p^s (\bar{B}_p^i) is the mean number of squeezed photons in the signal (idler) mode.

consider the Langevin forces \hat{L}_j arising in the interaction with the reservoir chaotic oscillators characterized by means of noise photon numbers $\langle n_d \rangle$. This leads to damping processes described by the damping constants γ_j .

The Heisenberg-Langevin operator equations corresponding to the Hamiltonian \hat{H}_{int} are derived in the following matrix form:

$$\frac{d\hat{\mathbf{a}}}{dt} = \mathbf{M}\hat{\mathbf{a}} + \hat{\mathbf{L}} \quad (3)$$

using the vectors $\hat{\mathbf{a}} = (\hat{a}_1, \hat{a}_1^\dagger, \hat{a}_2, \hat{a}_2^\dagger)^T$ and $\hat{\mathbf{L}} = (\hat{L}_1, \hat{L}_1^\dagger, \hat{L}_2, \hat{L}_2^\dagger)^T$, and the matrix

$$\mathbf{M} = \begin{bmatrix} -\gamma_1/2 & 2ig_{11} & 0 & ig_{12} \\ -2ig_{11} & -\gamma_1/2 & -ig_{12} & 0 \\ 0 & ig_{12} & -\gamma_2/2 & 2ig_{22} \\ -ig_{12} & 0 & -2ig_{22} & -\gamma_2/2 \end{bmatrix}. \quad (4)$$

The Langevin operators \hat{L}_1 and \hat{L}_2 introduced in Eq. (3) obey the following relations:

$$\begin{aligned} \langle \hat{L}_i(t) \rangle &= \langle \hat{L}_i^\dagger(t) \rangle = 0, \\ \langle \hat{L}_i^\dagger(t) \hat{L}_j(t') \rangle &= \delta_{ij} \langle n_d \rangle \delta(t - t'), \\ \langle \hat{L}_i(t) \hat{L}_j^\dagger(t') \rangle &= \delta_{ij} (\langle n_d \rangle + 1) \delta(t - t'), \end{aligned} \quad (5)$$

where δ_{ij} stands for the Kronecker symbol and δ denotes the Dirac delta function.

The solution of Eq. (3) for the operators \hat{a}_1 and \hat{a}_2 is conveniently written in the following matrix form using suitable evolution matrices \mathbf{U} and \mathbf{V} and a stochastic operator

vector $\hat{\mathbf{F}}$ (for details see, e.g., [44]):

$$\begin{bmatrix} \hat{a}_1(t) \\ \hat{a}_2(t) \end{bmatrix} = \mathbf{U}(t) \begin{bmatrix} \hat{a}_1(0) \\ \hat{a}_2(0) \end{bmatrix} + \mathbf{V}(t) \begin{bmatrix} \hat{a}_1^\dagger(0) \\ \hat{a}_2^\dagger(0) \end{bmatrix} + \hat{\mathbf{F}}(t). \quad (6)$$

Specific forms of the general evolution matrices \mathbf{U} and \mathbf{V} are discussed in the sections below. The elements of the stochastic operator vector $\hat{\mathbf{F}} \equiv (\hat{F}_1, \hat{F}_2)$ are derived as linear combinations of the Langevin forces \hat{L}_j and \hat{L}_j^\dagger that reflect the “deterministic” solution described by the matrices \mathbf{U} and \mathbf{V} [44].

Statistical properties of the emitted fields, in a given state $\hat{\rho}$, are described by the Glauber-Sudarshan P function, given by Eq. (1) or, equivalently, by the normal quantum characteristic function $C_{\mathcal{N}}$ defined as

$$C_{\mathcal{N}}(\beta_1, \beta_2) = \langle \exp(\beta_1 \hat{a}_1^\dagger + \beta_2 \hat{a}_2^\dagger) \exp(-\beta_1^* \hat{a}_1 - \beta_2^* \hat{a}_2) \rangle, \quad (7)$$

where the symbol $\langle \dots \rangle$ denotes quantum averaging including both system and reservoir. Using the solution given in Eq. (6) and the initial vacuum states in both fields, the normal characteristic function $C_{\mathcal{N}}$ attains the following form:

$$\begin{aligned} C_{\mathcal{N}}(\beta_1, \beta_2) &= \exp \left[-B_1 |\beta_1|^2 - B_2 |\beta_2|^2 + \left(\frac{C_1}{2} \beta_1^{*2} \right. \right. \\ &\quad \left. \left. + \frac{C_2}{2} \beta_2^{*2} + D_{12} \beta_1^* \beta_2^* + \bar{D}_{12} \beta_1 \beta_2 + \text{c.c.} \right) \right], \quad (8) \end{aligned}$$

where the auxiliary functions are defined as follows:

$$\begin{aligned} B_j &= \langle \Delta \hat{a}_j^\dagger \Delta \hat{a}_j \rangle = \sum_{k=1,2} |V_{jk}|^2 + \langle \hat{F}_j^\dagger \hat{F}_j \rangle, \\ C_j &= \langle (\Delta \hat{a}_j)^2 \rangle = \sum_{k=1,2} U_{jk} V_{jk} + \langle \hat{F}_j^2 \rangle, \\ D_{12} &= \langle \Delta \hat{a}_1 \Delta \hat{a}_2 \rangle = \sum_{k=1,2} U_{1k} V_{2k} + \langle \hat{F}_1 \hat{F}_2 \rangle, \\ \bar{D}_{12} &= -\langle \Delta \hat{a}_1^\dagger \Delta \hat{a}_2 \rangle = -\sum_{k=1,2} V_{1k}^* V_{2k} - \langle \hat{F}_1^\dagger \hat{F}_2 \rangle. \quad (9) \end{aligned}$$

The normal characteristic function given in Eq. (8) can conveniently be rewritten into its matrix form $C_{\mathcal{N}}(\boldsymbol{\beta}) = \exp(\boldsymbol{\beta}^\dagger \mathbf{A}_{\mathcal{N}} \boldsymbol{\beta} / 2)$ using the covariance matrix $\mathbf{A}_{\mathcal{N}}$ related to the normal ordering [45] (for different possibilities in describing the generated fields, see Table I):

$$\mathbf{A}_{\mathcal{N}} = \begin{bmatrix} -B_1 & C_1 & \bar{D}_{12}^* & D_{12} \\ C_1^* & -B_1 & D_{12}^* & \bar{D}_{12} \\ \bar{D}_{12} & D_{12} & -B_2 & C_2 \\ D_{12}^* & \bar{D}_{12}^* & C_2^* & -B_2 \end{bmatrix}, \quad (10)$$

and the column vector $\boldsymbol{\beta} = (\beta_1, \beta_1^*, \beta_2, \beta_2^*)^T$.

The covariance matrix $\mathbf{A}_{\mathcal{N}}$ related to the normal ordering determines the *global nonclassicality* of a two-mode Gaussian state via the Lee nonclassicality depth τ . The nonclassicality depth τ is defined with the help of the maximal positive eigenvalue $\lambda_+(\mathbf{A}_{\mathcal{N}})$ of the matrix $\mathbf{A}_{\mathcal{N}}$ as follows:

$$\tau = \max[0, \lambda_+(\mathbf{A}_{\mathcal{N}})]. \quad (11)$$

TABLE I. Schematic diagram for the relations between (a) two quasiprobability distributions (quasidistributions), i.e., the Glauber-Sudarshan P and Wigner W functions for a given two-mode state $\hat{\rho}$, (b) characteristic functions $C_{\mathcal{N}}$ and $C_{\mathcal{S}}$, and (c) covariance matrices $A_{\mathcal{N}}$ and $A_{\mathcal{S}}$ assuming here that $\hat{\rho}$ is a Gaussian state for normal and symmetric orderings, respectively. Their interrelations (as marked by left-right arrows) are given in the Appendix. The single arrow indicates that the calculation of the P function from the Wigner function is more complicated (it can be done via the relation between $C_{\mathcal{S}}$ and $C_{\mathcal{N}}$) than the trivial calculation of the Wigner function from the P function (as marked by double arrow).

Ordering	Quasidistribution		Characteristic function		Covariance matrix of a Gaussian state
Normal	$P(\alpha_1, \alpha_2) \equiv W^{(s=1)}(\alpha_1, \alpha_2)$	\iff	$C_{\mathcal{N}}(\beta_1, \beta_2)$	\longleftrightarrow	$A_{\mathcal{N}}$
	$\Downarrow \Uparrow$		$\Downarrow \Uparrow$		$\Downarrow \Uparrow$
Symmetric	$W(\alpha_1, \alpha_2) \equiv W^{(s=0)}(\alpha_1, \alpha_2)$	\iff	$C_{\mathcal{S}}(\beta_1, \beta_2)$	\longleftrightarrow	$A_{\mathcal{S}}$

We note that the nonclassicality depth τ , according to its definition [21], gives the amount of noise photons present equally in both modes and needed to conceal the nonclassical character of the state.

The covariance matrix $A_{\mathcal{N}}$ of the two-mode field can be written in the following block form:

$$\begin{aligned} \mathbf{A}_{\mathcal{N}} &= \begin{bmatrix} \mathbf{B}_1 & \mathbf{D}_{12} \\ \mathbf{D}_{21} & \mathbf{B}_2 \end{bmatrix}, \\ \mathbf{B}_j &= \begin{bmatrix} -B_j & C_j \\ C_j^* & -B_j \end{bmatrix}, \quad j = 1, 2, \\ \mathbf{D}_{12} &= \begin{bmatrix} \bar{D}_{12}^* & D_{12} \\ D_{12}^* & \bar{D}_{12} \end{bmatrix}, \\ \mathbf{D}_{21} &= \begin{bmatrix} \bar{D}_{12} & D_{12} \\ D_{12}^* & \bar{D}_{12}^* \end{bmatrix}. \end{aligned} \quad (12)$$

This form points out at the existence of three local invariants I_j , $j = 1, 2, 3$, that do not change under any local linear unitary transformation applied in mode 1 or 2. The local invariants I_j are expressed as

$$I_1 = \det(\mathbf{B}_1), \quad I_2 = \det(\mathbf{B}_2), \quad I_3 = \det(\mathbf{D}_{12}). \quad (13)$$

Moreover, there exist two global invariants I and Δ preserved under arbitrary linear unitary transformations and applied to both modes:

$$I = \det(\mathbf{A}_{\mathcal{N}}), \quad \Delta = I_1 + I_2 + 2I_3. \quad (14)$$

Whereas the global invariant I encompasses the whole complex structure of the matrix $\mathbf{A}_{\mathcal{N}}$ and, as such, is not useful in our considerations, the global invariant Δ reflects the block structure of the matrix $\mathbf{A}_{\mathcal{N}}$ and lies in the center of our attention.

Moreover, the global invariant Δ includes the additive local invariants I_1 and I_2 that indicate the nonclassical behavior of the reduced states of modes 1 and 2, respectively. Indeed, the determinants defining these invariants occur in the Fourier transform of the normal characteristic functions of the reduced states directly related to their local Glauber-Sudarshan P functions. If a determinant fails to be positive, then the corresponding Glauber-Sudarshan P function does not exist as a non-negative function. Thus, the value of determinant I_j can be used to quantify the *local nonclassicality* of the reduced state in mode j as it is a monotone of the local

Lee nonclassicality depth τ_j . The local Lee nonclassicality depth τ_j is defined along the formula (11) that provides the relation

$$\tau_j = \max(0, |C_j| - B_j), \quad j = 1, 2. \quad (15)$$

Using Eq. (15), we arrive at the monotonic relation between the local nonclassicality depth τ_j and local nonclassicality invariant (NI) I_j if we assume τ_j to be continuous:

$$I_j = -\tau_j(\tau_j + 2B_j). \quad (16)$$

We can redefine the local symplectic invariant in Eq. (16) as $I_{\text{nc}}^{(j)} = -I_j$ in order to deal with positive values when quantifying the local nonclassicality. We note that not only the positive values of this local NI $I_{\text{nc}}^{(1)}$ are useful for quantifying the local nonclassicality, also the negative values of this invariant are important as they quantify the ‘‘robustness’’ of the classicality of a local state.

Returning back to the last term I_3 in the global invariant Δ , this term describes solely the mutual quantum correlations between the fields 1 and 2. As such, it has to play a crucial role in the description of the entanglement between two fields. To reveal and quantify this entanglement and the role of local invariant I_3 here, we apply for a while the phase space (x, p) approach for describing the fields in the symmetric ordering of field operators corresponding to the Wigner formalism (see Table I and then the Appendix). The reason is technical and is given by the fact that we know how to derive the covariance matrix of a Gaussian state obtained by the partial transposition of the original state. According to Simon [46], the partial transposition means to replace p by $-p$. The covariance matrix of the partially transposed state then provides us the logarithmic negativity E_N that is a commonly used measure for the entanglement. Moreover, it provides as an entanglement measure useful in our considerations.

In detail, the covariance matrix $A_{\mathcal{S}}$ expressed in the symmetric ordering is obtained in its block structure as follows:

$$\begin{aligned} \mathbf{A}_{\mathcal{S}} &= \begin{bmatrix} \mathbf{B}_{S1} & \mathbf{D}_{\mathcal{S}} \\ \mathbf{D}_{\mathcal{S}}^T & \mathbf{B}_{S2} \end{bmatrix}, \\ \mathbf{B}_{Sj} &= \begin{bmatrix} B_j + \text{Re}(C_j) + 1/2 & \text{Im}(C_j) \\ \text{Im}(C_j) & B_j - \text{Re}(C_j) + 1/2 \end{bmatrix}, \\ & j = 1, 2, \\ \mathbf{D}_{\mathcal{S}} &= \begin{bmatrix} \text{Re}(D_{12} - \bar{D}_{12}) & \text{Im}(D_{12} + \bar{D}_{12}) \\ \text{Im}(D_{12} - \bar{D}_{12}) & -\text{Re}(D_{12} + \bar{D}_{12}) \end{bmatrix}. \end{aligned} \quad (17)$$

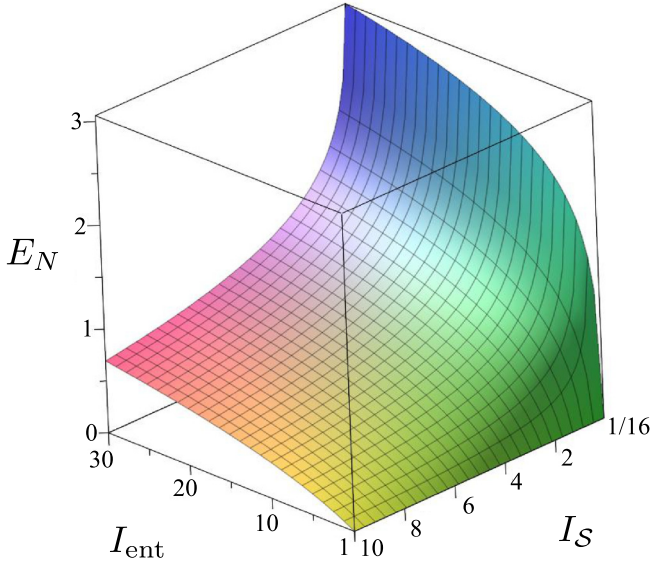


FIG. 3. Logarithmic negativity E_N as a function of entanglement indicator I_{ent} , given by Eq. (19), and global nonclassicality invariant I_S , given by Eq. (18).

The covariance matrix \mathbf{A}_S , similarly as its normally ordered counterpart, has three local invariants $I_{S_j}, j = 1, 2, 3$, and two global ones denoted as I_S and Δ_S :

$$I_{S1} = \det(\mathbf{B}_{S1}), \quad I_{S2} = \det(\mathbf{B}_{S2}), \quad I_{S3} = \det(\mathbf{D}_S),$$

$$I_S = \det(\mathbf{A}_S), \quad \Delta_S = I_{S1} + I_{S2} + 2I_{S3}. \quad (18)$$

Moreover, the comparison of the formulas for the invariants I_3 and I_{S3} shows that $I_3 = I_{S3}$.

Following Refs. [31,46,47], the entanglement criterion can be expressed through the positivity of the entanglement indicator (EI) I_{ent} defined in terms of the invariants in Eq. (18) as follows:

$$I_{\text{ent}} = \frac{1}{4}(I_{S1} + I_{S2} - 2I_{S3}) - I_S - \frac{1}{16}. \quad (19)$$

As we show below the EI I_{ent} is a monotonous function of logarithmic negativity E_N , which can be derived from the symplectic eigenvalue \tilde{d}_- of the partially transposed (PT) matrix \mathbf{A}_S^{PT} along the formula [48] (see Fig. 3)

$$E_N = \max[0, -\ln(2\tilde{d}_-)]. \quad (20)$$

According to Eq. (20), a state is entangled if $\tilde{d}_- < \frac{1}{2}$. In turn, the symplectic eigenvalue \tilde{d}_- is found as

$$\tilde{d}_- = \frac{1}{\sqrt{2}} \sqrt{\tilde{\Delta}_S - \sqrt{\tilde{\Delta}_S^2 - 4I_S}}, \quad (21)$$

where $\tilde{\Delta}_S = I_{S1} + I_{S2} - 2I_{S3}$. Combining Eqs. (19) and (21), we arrive at the relation between the symplectic eigenvalue \tilde{d}_- and entanglement indicator I_{ent} :

$$\tilde{d}_- = \frac{1}{\sqrt{2}} \sqrt{I' - \sqrt{I'^2 - 4I_S}}, \quad (22)$$

where $I' = 4I_S + 4I_{\text{ent}} + \frac{1}{4}$.

Assuming the global invariant I_S is fixed, the relation (22) shows that the larger is the entanglement indicator I_{ent} , the

smaller is the symplectic eigenvalue \tilde{d}_- and, according to formula (20), also the larger is the logarithmic negativity E_N . As a consequence, the entanglement indicator I_{ent} represents an alternative to the logarithmic negativity E_N in quantifying entanglement. We illustrate the monotonous dependence of the logarithmic negativity E_N on the entanglement indicator I_{ent} in Fig. 2. We note that a simple analytical formula between the logarithmic negativity E_N and entanglement indicator I_{ent} is derived for pure states ($I_S = \frac{1}{16}$) assuming $I_{\text{ent}} > 0$:

$$E_N = \ln(2\sqrt{I_{\text{ent}}} + \sqrt{1 + 4I_{\text{ent}}}). \quad (23)$$

As we look for a relation among the local invariants $I_{\text{nc1}}^{(1)}$ and $I_{\text{nc1}}^{(2)}$ and the entanglement indicator I_{ent} (see Fig. 1), we eliminate the invariants $I_3 = I_{S3}$ from Eqs. (13) and (19) by their comparing. This leaves us with the relation

$$I_{\text{nc1}}^{(1)} + I_{\text{nc1}}^{(2)} + 2I_{\text{ent}} = \frac{1}{2}\Delta_S - \Delta - 2I_S - \frac{1}{8}. \quad (24)$$

As only the global invariants occur at the right-hand side of Eq. (24), the relation $I_{\text{nc1}}^{(1)} + I_{\text{nc1}}^{(2)} + 2I_{\text{ent}}$ is invariant under any global linear unitary transformation.

Equation (24) can be transformed into the central result of our paper, if we define a new quantity I_{nc1} , which is a global nonclassicality invariant:

$$I_{\text{nc1}} = I_{\text{nc1}}^{(1)} + I_{\text{nc1}}^{(2)} + 2I_{\text{ent}}. \quad (25)$$

In the derivation of this equation, it is useful to recall the property that the local determinants for the normally ordered CF, I_3 , and the symmetrically ordered CF, I_{S3} , are equal $I_3 = I_{S3}$, and given by Eqs. (14) and (18). Thus, we have

$$\begin{aligned} I_{\text{nc1}} &= I_{\text{nc1}}^{(1)} + I_{\text{nc1}}^{(2)} + 2I_{\text{ent}} \\ &= -I_1 - I_2 + \frac{1}{2}(I_{S1} + I_{S2} - 2I_{S3}) - 2I_S - \frac{1}{8} \\ &= -I_1 - I_2 - 2I_{S3} + \frac{1}{2}(I_{S1} + I_{S2} + 2I_{S3}) - 2I_S - \frac{1}{8} \\ &= -\Delta + \frac{1}{2}\Delta_S - 2I_S - \frac{1}{8}. \end{aligned} \quad (26)$$

Equation (25) means that the local nonclassicality invariants $I_{\text{nc1}}^{(1)}$ and $I_{\text{nc1}}^{(2)}$ together with the entanglement indicator I_{ent} form the global NI I_{nc1} . Any linear unitary transformation in general modifies both the local NIs $I_{\text{nc1}}^{(1)}$ and $I_{\text{nc1}}^{(2)}$ and the entanglement invariant I_{ent} only in such a way that it preserves the value of the global NI I_{nc1} . Whenever I_{nc1} is positive, the analyzed state is nonclassical due to the local nonclassicality of the reduced states or its entanglement. The negative values of the global NI I_{nc1} do not necessarily mean that a given state is classical, as we will see below.

In the next sections, we analyze the nonclassicality and entanglement of several kinds of important quantum states from the point of view of their transformation by a beam splitter. The division of the global NI into the EI and the local NIs is in the center of our attention. In general, six regions differing in the occurrence of entanglement and local nonclassicalities can be defined (see Table II). All these regions are found in the examples analyzed in the next sections, as indicated in Table II.

We note that an invariant based on the second-order intensity moments and, as such, describing intensity autocorrelations and cross-correlations has been suggested in Ref. [49]

TABLE II. Regions of different entanglement and local nonclassicalities observed in the figures of Secs. III–VI.

Case/region	Entanglement	Nonclassicality of one mode	Nonclassicality of another mode	Figures
I	Yes	Yes	Yes	6, 10
II	Yes	Yes	No	6(b)
III	Yes	No	No	6, 10
IV	No	Yes	Yes	6, 10
V	No	Yes	No	6(b)
VI	No	No	No	6, 10

for two-mode fields with specific mode correlations and unitary transformations. Later, this invariant was experimentally analyzed in Ref. [50]. Here, we describe the propagation of fields through the beam splitter (see Fig. 1) described by the real transmissivity T and the phase ϕ through the unitary transformation characterized by the matrix \mathbf{U} :

$$\mathbf{U} = \begin{pmatrix} \sqrt{T} & 0 & -\sqrt{R}e^{i\phi} & 0 \\ 0 & \sqrt{T} & 0 & -\sqrt{R}e^{-i\phi} \\ \sqrt{R}e^{-i\phi} & 0 & \sqrt{T} & 0 \\ 0 & \sqrt{R}e^{i\phi} & 0 & \sqrt{T} \end{pmatrix}; \quad (27)$$

$R = 1 - T$ is the reflectivity of the beam splitter. The covariance matrix \mathbf{A}^{out} at the output of the beam splitter is obtained as $\mathbf{A}^{\text{out}} = \mathbf{U}^\dagger \mathbf{A} \mathbf{U}$.

III. TWIN BEAM

These beams are generated by parametric down-conversion from the vacuum into which photon pairs are ideally emitted. For this reason, only the terms B_1, B_2 , and D_{12} in the normal characteristic function $C_{\mathcal{N}}$ are nonzero. The evolution matrices \mathbf{U} and \mathbf{V} in Eq. (6) have the following nonzero elements:

$$\begin{aligned} U_{11}(t) &= U_{22}(t) = \cosh(gt), \\ V_{12}(t) &= V_{21}(t) = i \exp(i\theta) \sinh(gt). \end{aligned} \quad (28)$$

The coefficients B_1 and B_2 can be expressed as $B_1 = B_p + B_s$ and $B_2 = B_p + B_i$, where $B_p = \sinh^2(gt)$ gives the mean number of generated photon pairs and $B_s = \langle \hat{F}_1^\dagger \hat{F}_1 \rangle$ ($B_i = \langle \hat{F}_2^\dagger \hat{F}_2 \rangle$) denotes the mean number of signal (idler) noise photons coming from the reservoir (see Fig. 2). On the other hand, the parameter D_{12} characterizing mutual correlations depends only on the mean number B_p of photon pairs as $D_{12} = i\sqrt{B_p(B_p + 1)}$ ($\theta = 0$ is assumed without the loss of generality).

The general formulas for the local NIs $I_{\text{nc1}}^{(j)}$, entanglement invariant I_{ent} , and the global NI I_{nc1} attain the following forms for twin beams:

$$\begin{aligned} I_{\text{nc1}}^{(1)} &= 4TR(B_p^2 + B_p) - [B_p + TB_s + RB_i]^2, \\ I_{\text{nc1}}^{(2)} &= 4TR(B_p^2 + B_p) - [B_p + TB_i + RB_s]^2, \\ I_{\text{ent}} &= -[(B_s + B_i)^2 - (T - R)^2](B_p^2 + B_p) \\ &\quad - 2B_p B_s B_i (B_s + B_i) - (B_s^2 + B_s)(B_i^2 + B_i) \\ &\quad - TR(B_s + B_i)^2, \end{aligned} \quad (29)$$

$$\begin{aligned} I_{\text{nc1}} &= 2B_p - (B_s + B_i)^2 [2(B_p^2 + B_p) + 1] \\ &\quad - 2B_p(1 + 2B_s B_i)(B_s + B_i) \\ &\quad - 2B_s B_i (B_s + B_i + B_s B_i). \end{aligned} \quad (30)$$

We first discuss the behavior of noiseless twin beams for which $B_s = B_i = 0$. In this case, the global NI I_{nc1} equals $2B_p$ and

$$\begin{aligned} I_{\text{nc1}}^{(j)} &= 4TR(B_p^2 + B_p) - B_p^2, \quad j = 1, 2 \\ I_{\text{ent}} &= (T - R)^2(B_p^2 + B_p). \end{aligned} \quad (31)$$

As suggested by the formula in Eq. (31), the local NIs $I_{\text{nc1}}^{(j)}$ can be decomposed into two terms. The negative term reflects classical thermal statistics of photon pairs in a twin beam with its photon bunching effect and as such suppresses the nonclassical behavior of the twin beam. On the other hand, the positive term refers to squeezing appearing at the individual output ports of the beam splitter. The squeezing effect originates in pairing of photons in individual output ports caused by “sticking of two photons from a pair together” (photon bunching) at the beam splitter [5]. Photon pairs with both photons in one output port contribute to the local nonclassicality of the field in this port. On the other hand, when two photons from one photon pair occur in different output ports, they contribute to the entanglement. “A given individual photon pair” is, thus, responsible either for the local nonclassicality in one of the output ports or for their entanglement. Never for both. Propagation through the beam splitter can, thus, be viewed as the process of breaking photon pairs (antibunching) arriving at the same input port and gluing (bunching) of photons from a given pair coming from different input ports. Whereas the first process disturbs local squeezing and supports entanglement, the second process strengthens squeezing at the expense of entanglement. The global NI I_{nc1} is equal twice the number B_p of photon pairs and, as such, indicates an appropriate choice of this nonclassicality resource quantifier.

Detailed analysis of the formulas in Eq. (31) shows that the local marginal states are nonclassical only if the transmissivity T lies in certain interval around $\frac{1}{2}$:

$$T \in \left(\frac{1}{2} - \frac{1}{2\sqrt{B_p + 1}}, \frac{1}{2} + \frac{1}{2\sqrt{B_p + 1}} \right). \quad (32)$$

It holds that the larger is the mean photon-pair number B_p , the narrower is the interval. The optimal transmissivity T maximizing the local NIs $I_{\text{nc1}}^{(j)}$ equals $\frac{1}{2}$. In this case, the entanglement of the incident twin beam is completely and

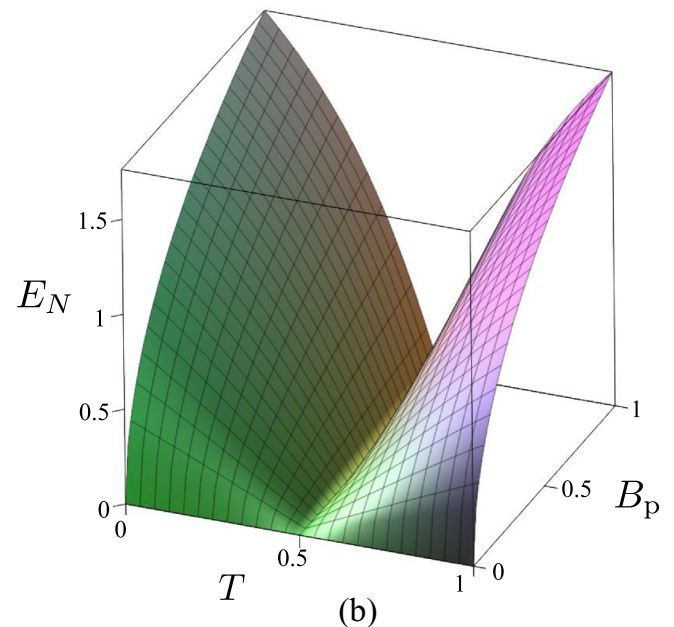
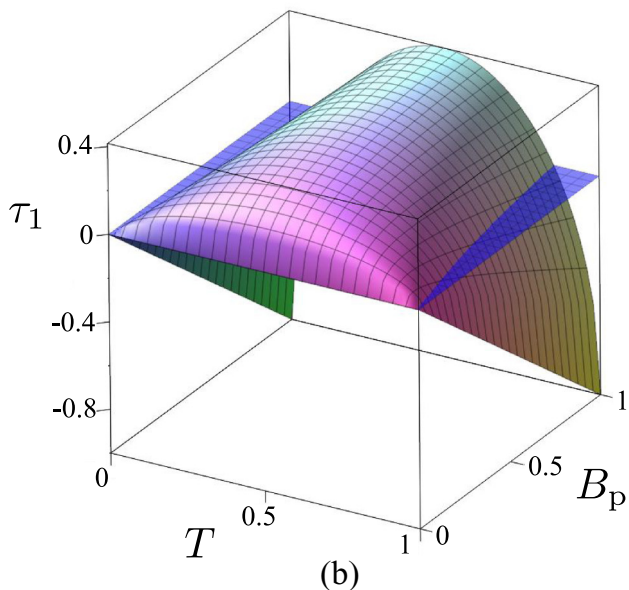
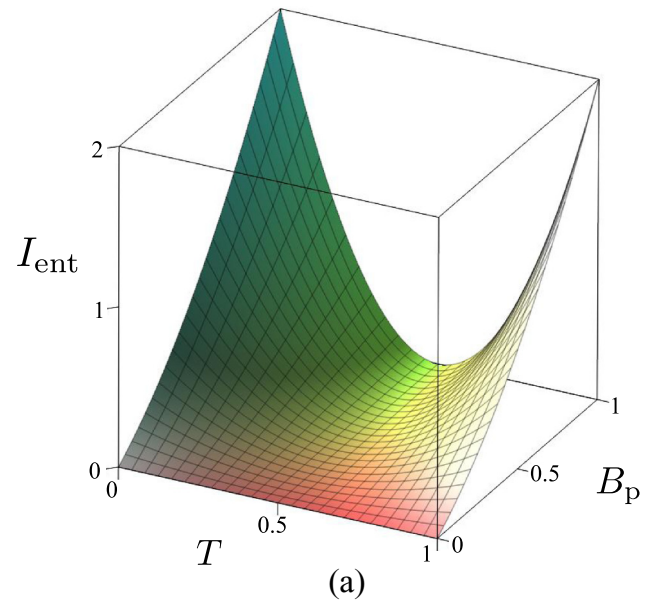
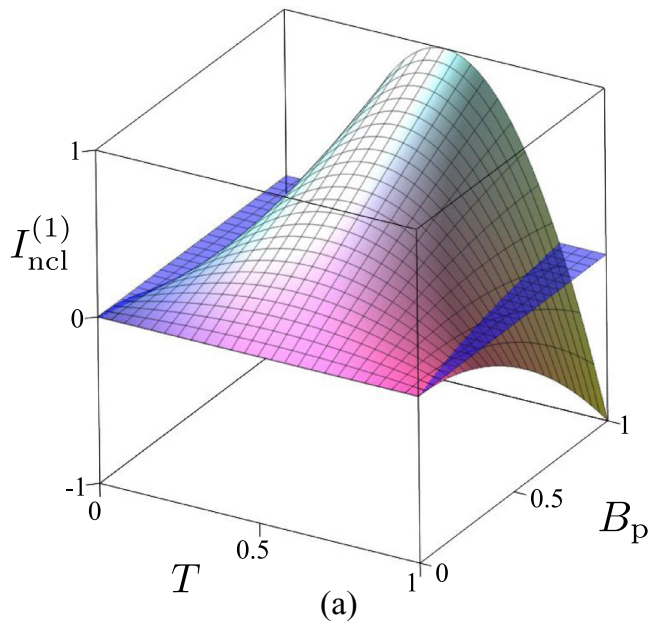


FIG. 4. (a) Local nonclassicality invariant $I_{\text{ncl}}^{(1)}$ and (b) continuous Lee nonclassicality depth τ_1 (including negative values) at the output port 1 of the beam splitter as a function of the mean photon-pair number B_p and the beam-splitter transmissivity T for pure twin beam states. In panels (a) and (b), the blue dark gray plain surface at $I_{\text{ncl}}^{(1)} = 0$ and $\tau_1 = 0$ shows the boundary between the classical and nonclassical domains.

FIG. 5. (a) Entanglement invariant E_I and (b) logarithmic negativity E_N after the beam splitter transformation considered as functions of the mean photon-pair number B_p and the beam-splitter transmissivity T for pure twin beams states.

equally transferred into the local nonclassicalities of the two output modes. On the other hand, the twin beam loses its entanglement only when $T = \frac{1}{2}$. In this case, all the incident photon pairs stick together (bunch) at the beam splitter suppressing completely their entanglement. Hand in hand, the local NIs $I_{\text{ncl}}^{(1)} = I_{\text{ncl}}^{(2)}$ attain their maximal values. This can be interpreted such that the initial entanglement is transferred into the squeezing of the marginal output fields [51]. These effects are shown in Figs. 4(a) and 5(a) for the dependencies of the local NI $I_{\text{ncl}}^{(1)}$ and EI I_{ent} on the transmissivity T

and mean photon-pair number B_p . The commonly used Lee nonclassicality depth τ_1 and the logarithmic negativity E_N are shown for comparison in Figs. 4(b) and 5(b). We note that whereas the values of the Lee nonclassicality depth τ_1 cannot exceed $\frac{1}{2}$, the values of the local NI $I_{\text{ncl}}^{(1)}$ can be arbitrarily large depending on the intensity of the twin beam.

Now, we consider general noisy twin beams. It has been shown in Ref. [36] that whenever the overall noise $B_s + B_i$ exceeds one, the twin beam is unentangled and, thus, it cannot generate any nonclassical feature. Even if $B_s + B_i < 1$, the mean photon-pair number B_p has to be sufficiently large, as

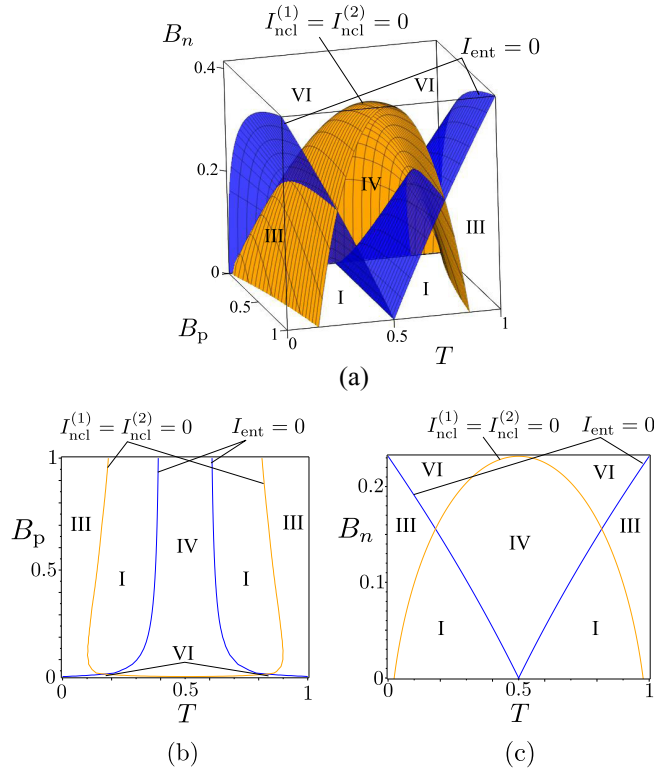


FIG. 6. Diagram (a) shows the nonclassicality and entanglement invariants for the twin beams states occurring at the output ports of a beam splitter depending on the mean noise photon number B_n , mean photon-pair number B_p , and transmissivity T according to Eq. (29) for $B_n \equiv B_s = B_i$. The surfaces are plotted at $I_{\text{ncl}}^{(1)}(B_n, B_p, T) = 0$ (orange light gray surface), $I_{\text{ncl}}^{(2)}(B_n, B_p, T) = 0$ (orange light gray), and $E_I(B_n, B_p, T) = 0$ (blue dark surface) indicating six different regions specified in the text and Table II. Diagrams (b) and (c) show the perpendicular cross sections of diagram (a) taken at chosen values of $B_n = 0.1$ and $B_p = 0.1$, respectively.

given by

$$B_p > \frac{B_s B_i}{1 - (B_s + B_i)}. \quad (33)$$

Then, the incident noisy twin beam is entangled and is capable to provide its entanglement and local nonclassicality after the beam splitter. However, the general analysis of Eqs. (29) and (30) leads to the conclusion that the noise only degrades the nonclassical behavior independently whether it is manifested by local nonclassicality or entanglement. The stronger the noise, the weaker the nonclassical features.

To provide a deeper insight into the role of noise, we analyze two special cases: in the first one, the noise is equally divided into both modes of the incident twin beam, while noise occurs only in one mode of the incident twin beam in the second case.

When noise occurs in both modes of the incident twin beam ($B_n \equiv B_s = B_i$), the globally nonclassical output states can be divided into three groups. They are displayed in the “phase diagram” in Fig. 6. In this diagram, the surfaces $I_{\text{ncl}}^{(1)}(B_n, B_p, T) = 0$ and $I_{\text{ent}}(B_n, B_p, T) = 0$ are shown. They identify four different regions belonging to different groups of states (see Table II for details). The states exhibiting both

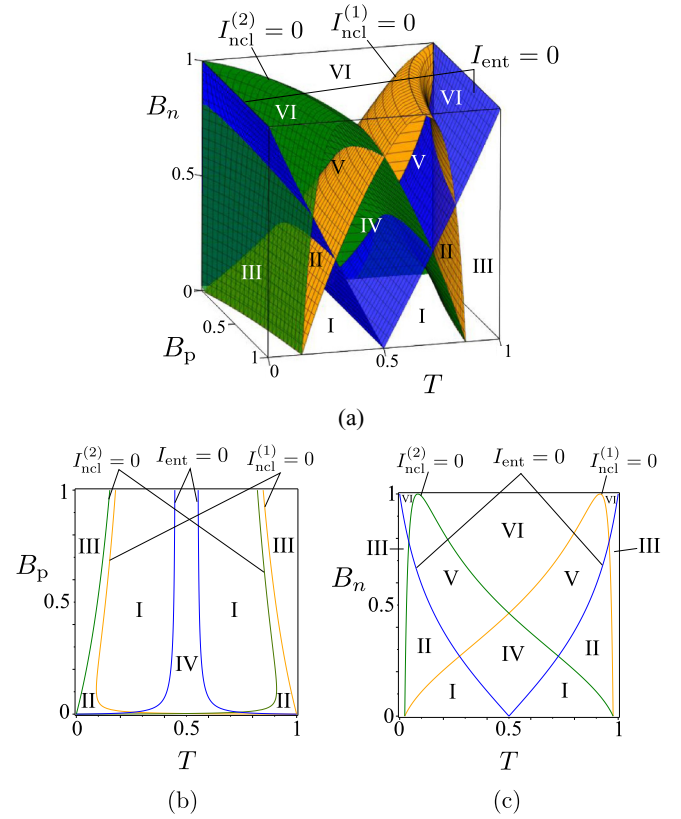


FIG. 7. Diagram (a) shows the nonclassicality and entanglement invariants for the twin beams states occurring at the output ports of a beam splitter depending on the mean noise photon number B_n , mean photon-pair number B_p , and transmissivity T according to Eq. (29) for $B_s = 0$ and $B_n = B_i$. The surfaces are plotted at $I_{\text{ncl}}^{(1)}(B_n, B_p, T) = 0$ (orange light gray surface), $I_{\text{ncl}}^{(2)}(B_n, B_p, T) = 0$ (green dark gray surfaces), and $E_I(B_n, B_p, T) = 0$ (blue dark surface) indicating six different regions specified in the text and Table II. Diagrams (b) and (c) show the perpendicular cross sections of diagram (a) taken at fixed values of $B_n = 0.1$ and $B_p = 0.1$, respectively. These cross sections are analogous to those in Figs. 6(b) and 6(c).

entanglement and local nonclassicality occur in region I. In region III, the states are entangled but locally classical. The locally nonclassical and unentangled states are found in region IV. In region VI, the unentangled and locally classical states exist.

The presence of noise in only one mode of the incident twin beam ($B_s = 0, B_i \equiv B_n \neq 0$) leads to asymmetry between the output modes. This is shown in Fig. 7, where the surfaces $I_{\text{ncl}}^{(1)}(B_n, B_p, T) = 0$ and $I_{\text{ncl}}^{(2)}(B_n, B_p, T) = 0$ behave differently. The symmetry, with respect to the plane for $T = \frac{1}{2}$, which is clearly visible in Fig. 6, does not exist in Fig. 7. As a consequence, two additional groups of states are found in the diagram. In region V, there are unentangled states with only one marginal field exhibiting local nonclassicality. The entangled states with only one locally nonclassical field are found in region II. In detail, mode I (2) is locally nonclassical for $T < \frac{1}{2}$ ($T > \frac{1}{2}$). We note that the EI I_{ent} is not sensitive to the noise asymmetry, as shown by the surface $I_{\text{ent}}(B_n, B_p, T) = 0$ in Fig. 7. It is worth noting that

positive values of the GNI I_{ncl} are exhibited when either entanglement or local nonclassicality or even both are found. The negative values of the global NI I_{ncl} do not necessarily mean classicality. The state with the negative GNI I_{ncl} can still be globally nonclassical due to either its entanglement or local nonclassicality, but not both. The diagram in Fig. 6(a) can serve as an example. The surface $I_{\text{ncl}}(B_n, B_p, T) = 0$ lies naturally in-between the surfaces $I_{\text{ncl}}^{(1)}(B_n, B_p, T) = 0$, and $I_{\text{ent}}(B_n, B_p, T) = 0$ and its position identifies the globally nonclassical states with $I_{\text{ncl}} < 0$.

IV. SQUEEZED VACUUM STATE WITH NOISE

Here, we consider a squeezed vacuum state [5] mixed with the noise incident on one input port of the beam splitter, whereas the second input port is left in the vacuum state. In this case, the nonzero elements of evolution matrices \mathbf{U} and \mathbf{V} in Eq. (6) are given as ($\theta = 0$ is assumed)

$$\begin{aligned} U_{11}(t) &= \cosh(gt), & U_{22}(t) &= 1, \\ V_{11}(t) &= i \exp(i\theta) \sinh(gt). \end{aligned} \quad (34)$$

The nonzero parameters of the normal characteristic function $C_{\mathcal{N}}$ in Eq. (10) are B_1 and C_1 as given by $B_1 = \tilde{B}_p^{\text{sq}} + B_s$ and $C_1 = i\sqrt{\tilde{B}_p^{\text{sq}}(\tilde{B}_p^{\text{sq}} + 1)}$. The symbol \tilde{B}_p^{sq} denotes the mean number of squeezed photons and the symbol B_s stands for the mean number of the signal noise photons (see also Fig. 2). The local NIs $I_{\text{ncl}}^{(j)}$ and EI I_{ent} are easily expressed in terms of the global NI I_{ncl} as follows:

$$\begin{aligned} I_{\text{ncl}}^{(1)} &= T^2 I_{\text{ncl}}, & I_{\text{ncl}}^{(2)} &= R^2 I_{\text{ncl}}, & I_{\text{ent}} &= T R I_{\text{ncl}}, \\ I_{\text{ncl}} &= \tilde{B}_p^{\text{sq}}(1 - 2B_s) - B_s^2. \end{aligned} \quad (35)$$

As the local NIs $I_{\text{ncl}}^{(1)}$ and $I_{\text{ncl}}^{(2)}$, as well as the EI I_{ent} are linearly proportional to the global NI I_{ncl} , the global nonclassicality of the output states immediately guarantees both local nonclassicalities and entanglement. This occurs only for the positive values of the global NI I_{ncl} . According to Eq. (35), $I_{\text{ncl}} > 0$ provided that the mean noise photon number B_s in the signal mode is sufficiently small:

$$B_s < \sqrt{\tilde{B}_p^{\text{sq}}(\tilde{B}_p^{\text{sq}} + 1)} - \tilde{B}_p^{\text{sq}}. \quad (36)$$

Following Eq. (35), the mean noise photon number B_s in the signal mode has to be smaller than 1. Also, the more intense is the squeezed state, the smaller is the number B_s of accepted noise photons. We note that the condition, given in Eq. (35), can immediately be revealed when the global Lee nonclassicality depth τ is analyzed. As an illustration, the dependencies of the local NIs $I_{\text{ncl}}^{(1)}$ and $I_{\text{ncl}}^{(2)}$ and the EI I_{ent} on the beam-splitter transmissivity T are plotted in Fig. 8 for the incident noiseless squeezed states. The greatest values of EI I_{ent} are reached for the balanced beam splitter ($T = \frac{1}{2}$). However, some incident photon pairs are not broken (i.e., split) by the beam splitter and give rise to nonzero local nonclassicalities $I_{\text{ncl}}^{(1)}$ and $I_{\text{ncl}}^{(2)}$ even in this case.

The strength of squeezing in a given mode is commonly characterized by a principal squeeze variance λ [52], which is

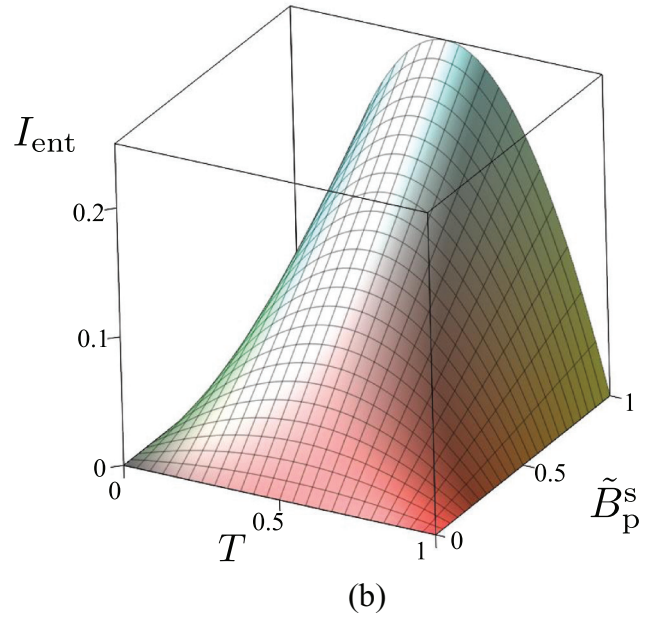
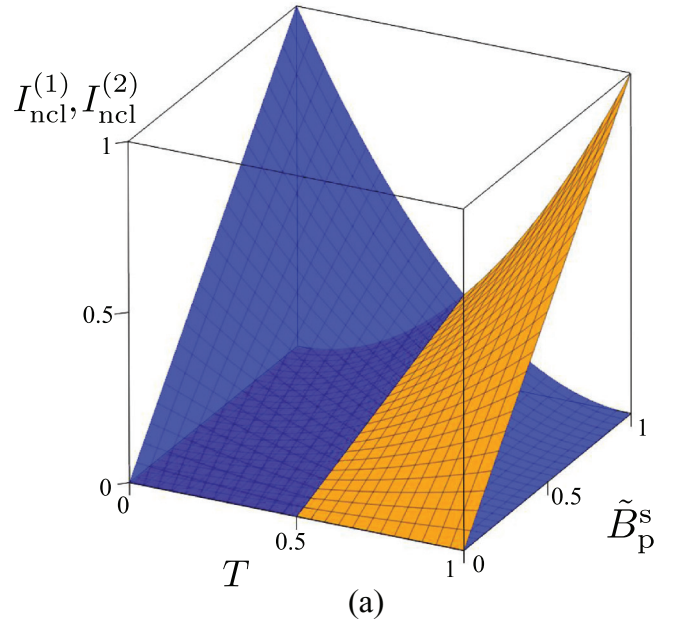


FIG. 8. Invariant nonclassicality parameters: (a) the local nonclassicality invariants $I_{\text{ncl}}^{(1)}$ (orange light gray surface) and $I_{\text{ncl}}^{(2)}$ (blue dark gray surface), and (b) the entanglement invariant I_{ent} versus the mean number \tilde{B}_p^{sq} of squeezed photons and the beam-splitter transmissivity T according to Eq. (35) assuming $B_s = 0$.

here given by

$$\lambda_j = 1/2 + B_j - |C_j|. \quad (37)$$

When a given output mode $j = 1, 2$ is locally nonclassical, it is also squeezed, which corresponds to $\lambda_j < \frac{1}{2}$. According to the relation between the local NI $I_{\text{ncl}}^{(j)}$ and the principal squeeze variance λ_j derived by combining Eqs. (16) and (37),

$$I_{\text{ncl}}^{(j)} = (1/2 - \lambda_j)(2B_j + 1/2 - \lambda_j), \quad (38)$$

the smaller is the value of the principal squeeze variance λ_j below $\frac{1}{2}$, the greater is the value of the local NI $I_{\text{ncl}}^{(j)}$.

V. TWO SQUEEZED VACUA

Two independent squeezed states are generated by the Hamiltonian given in Eq. (2) provided that the process of parametric down-conversion does not occur in the nonlinear medium ($g_{12} = 0$). The solution of the evolution governed by the Hamiltonian (2) gives us the following nonzero elements of the evolution matrices \mathbf{U} and \mathbf{V} :

$$\begin{aligned} U_{11} &= \cosh(2g_{11}t), & V_{11} &= i \exp(i\kappa_1) \sinh(2g_{11}t), \\ U_{22} &= \cosh(2g_{22}t), & V_{22} &= i \exp(i\kappa_2) \sinh(2g_{11}t), \end{aligned} \quad (39)$$

where κ_1 and κ_2 are arbitrary phases. The nonzero coefficients of the incident covariance matrix $\mathbf{A}_{\mathcal{N}}$ are given as $B_{1,2} = \tilde{B}_p^{s,i} + B_{s,i}$ and $C_{1,2} = \exp(i\theta_{1,2})\sqrt{\tilde{B}_p^{s,i}(\tilde{B}_p^{s,i} + 1)}$, $\theta_j = \kappa_j + \pi/2$ for $j = 1, 2$, where \tilde{B}_p^s (\tilde{B}_p^i) stands for the mean number of squeezed photons in the signal (idler) mode, whereas the corresponding mean signal (idler) noise photon number is denoted as B_s (B_i).

After the beam splitter, the local NIs $I_{\text{ncl}}^{(j)}$, EI I_{ent} , and global NI I_{ncl} acquire the form

$$\begin{aligned} I_{\text{ncl}}^{(1)} &= T^2 \tilde{B}_p^s (\tilde{B}_p^s + 1) + R^2 \tilde{B}_p^i (\tilde{B}_p^i + 1) + TR \bar{D}'_{12} \\ &\quad \times \cos(\theta_1 - \theta_2) - [T \tilde{B}_p^s + R \tilde{B}_p^i + T B_s + R B_i]^2, \\ I_{\text{ncl}} &= B_1 + B_2 - 2B_s B_i [2B_1(1 + \tilde{B}_p^i) + 2\tilde{B}_p^i(1 + B_1) \\ &\quad + B_i(1 + 2B_1) + B_s(1 + 2B_2)] \\ &\quad - 2(B_s B_1 + B_i B_2) - (B_s + B_i)^2, \\ I_{\text{ent}} &= TR [-\bar{D}'_{12} \cos(\theta_1 - \theta_2) + (\tilde{B}_p^s + \tilde{B}_p^i + 2\tilde{B}_p^s \tilde{B}_p^i) \\ &\quad - (B_s + B_i)^2 - 2(\tilde{B}_p^s - \tilde{B}_p^i)(B_s - B_i)] \\ &\quad + B_s B_i [2\tilde{B}_p^s(1 + B_i) + 2\tilde{B}_p^i(1 + B_s) \\ &\quad + 4\tilde{B}_p^s \tilde{B}_p^i + (1 + B_s)(1 + B_i)], \end{aligned} \quad (40)$$

where $\bar{D}'_{12} = 2\sqrt{\tilde{B}_p^s(\tilde{B}_p^s + 1)\tilde{B}_p^i(\tilde{B}_p^i + 1)}$, $B_1 = \tilde{B}_p^s + B_s$, $B_2 = \tilde{B}_p^i + B_i$, and, for simplicity, we assumed $\phi = 0$ in Eq. (27).

The formula for $I_{\text{ncl}}^{(2)}$ is obtained from that for $I_{\text{ncl}}^{(1)}$ in Eq. (40) with the substitution $s \leftrightarrow i$.

The global NI I_{ncl} does not depend on the relative phase $\Delta\theta = \theta_1 - \theta_2$ of two incident squeezed states, while the local NIs $I_{\text{ncl}}^{(j)}$ and EI I_{ent} change with the relative phase $\Delta\theta$. The case of two equally intense incident noiseless squeezed states, as graphically analyzed in Fig. 9, shows that the phase difference $\Delta\theta$ plays a crucial role in distributing the nonclassicality between the output entanglement and local nonclassicalities. If the phases θ_1 and θ_2 are equal, the incident photon pairs stick (bunch) ideally together due to the interference at the beam splitter and the incident locally nonclassical squeezed states are moved into the output ports. No photon pair is broken and so no entanglement is observed. On the other hand, if $\Delta\theta = \pi$, then some incident photon pairs are broken and, thus, the output squeezing (as well as local nonclassicalities) is weaker. The broken photon pairs give rise to the entanglement. The value of EI I_{ent} is maximal for the transmissivity $T = \frac{1}{2}$. In this case, all the photon pairs are broken, their signal and idler photons occur in different output ports, and, as a consequence,

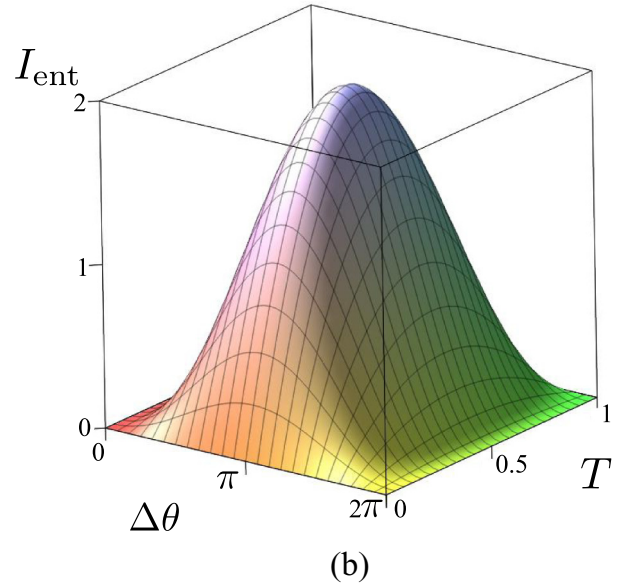
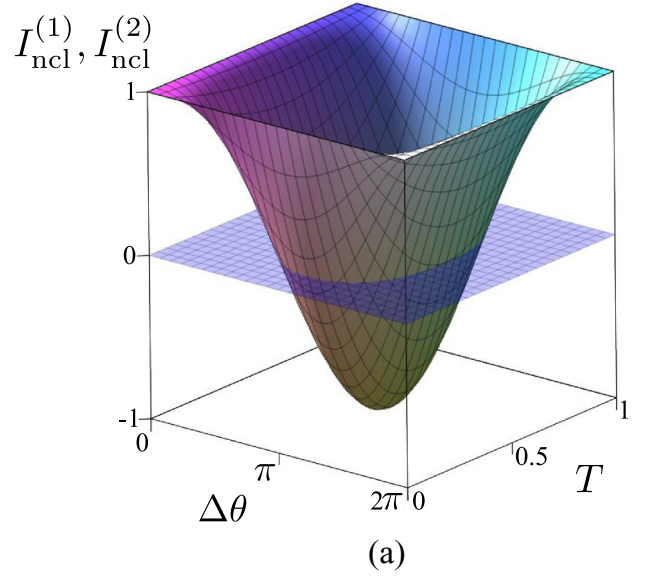


FIG. 9. (a) Local nonclassicality invariants $I_{\text{ncl}}^{(1)} = I_{\text{ncl}}^{(2)}$ and (b) entanglement invariant I_{ent} versus the phase difference $\Delta\theta$ and beam-splitter transmissivity T for two noiseless squeezed states according to Eq. (40); $\tilde{B}_p^s = \tilde{B}_p^i = 1$. In panel (a), the blue surface at $I_{\text{ncl}}^{(1)} = I_{\text{ncl}}^{(2)} = 0$ shows the boundary between classical and nonclassical states.

the ideal conditions for entanglement generation are met. Hand in hand, the vanishing local NIs $I_{\text{ncl}}^{(j)}$ are found (see Fig. 9).

It is remarkable that the global NI I_{ncl} for the equally intense noiseless squeezed states is given formally by the same formula as that valid for the noiseless twin beams considering the mean photon-pair number B_p instead of $\tilde{B}_p^s = \tilde{B}_p^i \equiv \tilde{B}_p$. However, the incident twin beam serves as a source of locally nonclassical (squeezed) states, whereas the incident squeezed states provide entangled states at the output of the beam splitter. The comparison of graphs in Figs. 4(a) and 5(a) with those in Figs. 10(a) and 10(b) reveals that the incident noiseless squeezed states generate entangled states for an arbitrary value of the transmissivity T , but the incident noiseless twin beams

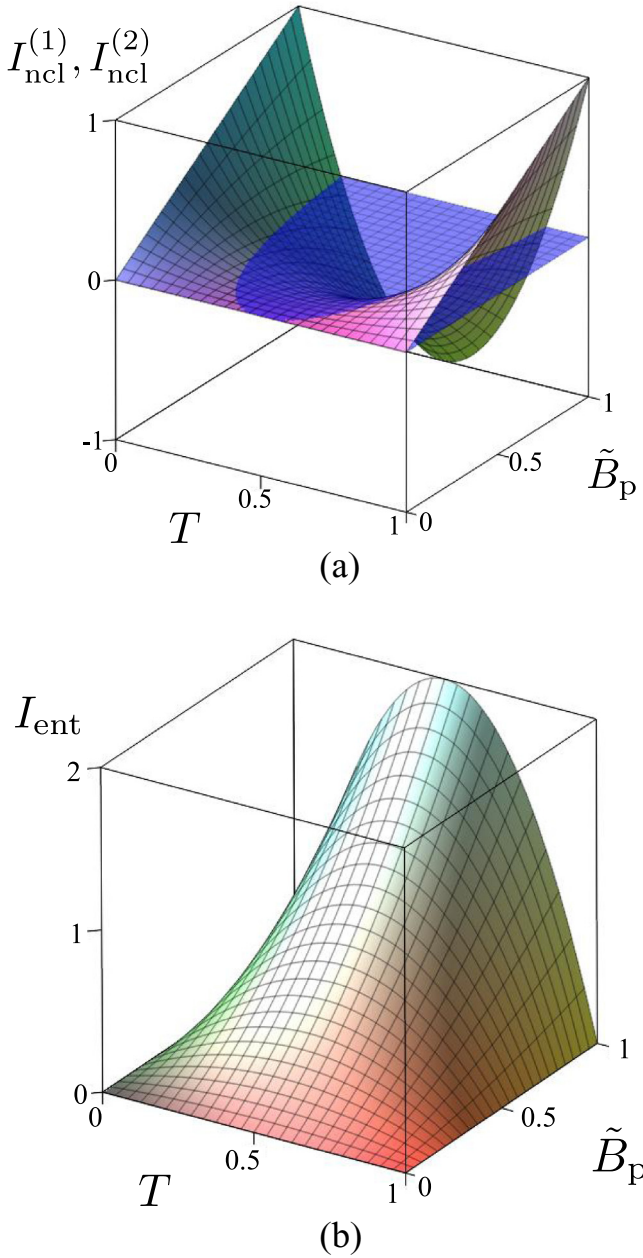


FIG. 10. (a) Local nonclassicality invariant $I_{\text{ncl}}^{(1)}$ and (b) entanglement invariant I_{ent} versus the beam-splitter transmissivity T and mean number \tilde{B}_p of squeezed photons for two noiseless squeezed states according to Eq. (40); $\tilde{B}_p \equiv \tilde{B}_p^s = \tilde{B}_p^i$; $\Delta\theta = \pi$. In panel (a) the blue surface at $I_{\text{ncl}}^{(1)} = I_{\text{ncl}}^{(2)} = 0$ shows the boundary between classical and nonclassical states.

are capable of the generation of the output squeezed states only in a certain interval of the transmissivity T depending on the intensity.

Similarly as for the twin beams, the noise diminishes the global NI I_{ncl} [see the formula for I_{ncl} in Eq. (40)]. Considering the incident states with $\tilde{B}_p^s = \tilde{B}_p^i$ and $B_s = B_i$, the presence of noise leads to the occurrence of the three different types of globally nonclassical states already discussed in the connection with the noisy twin beams with symmetric noise. Regions corresponding to different types of the output states

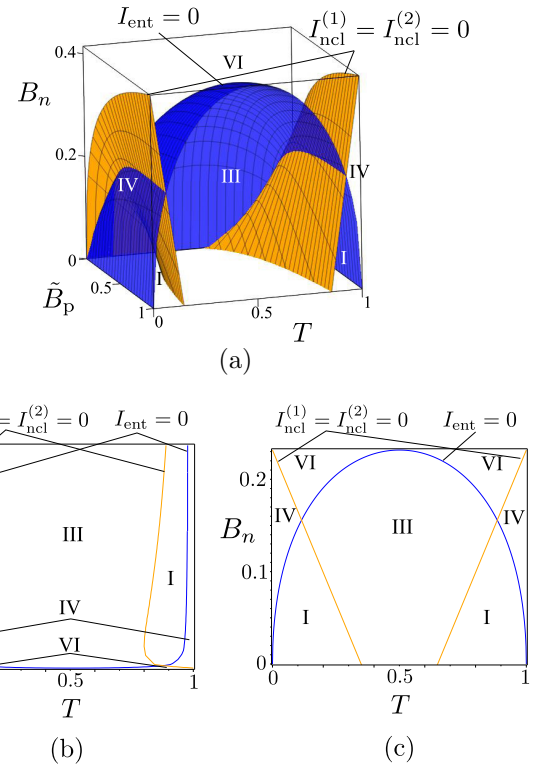


FIG. 11. Diagram (a) shows the nonclassicality and entanglement invariants for the two squeezed vacua occurring at the output ports of the beam splitter versus the mean noise photon number B_n , mean number \tilde{B}_p of squeezed photons, and transmissivity T assuming $B_n \equiv B_s = B_i$ and $\tilde{B}_p \equiv \tilde{B}_p^s = \tilde{B}_p^i$ and $\Delta\theta = \pi$. Surfaces at $I_{\text{ncl}}^{(j)}(B_n, \tilde{B}_p, T) = 0$ ($j = 1, 2$) (orange light gray) and $I_{\text{ent}}(B_n, \tilde{B}_p, T) = 0$ (blue dark gray) are shown surrounding different regions specified in Table II. Diagrams (b) and (c) show the perpendicular cross sections of diagram (a) taken at given values of $B_n = 0.1$ and $\tilde{B}_p = 0.1$, respectively. These cross sections can be compared with those in panels (b) and (c) in Figs. 6 and 7.

are shown in the diagram in Fig. 11(a) that can be compared with that of Fig. 6(a).

VI. TWIN BEAM MIXED WITH SQUEEZED STATES

Finally, we analyze an interplay of noiseless twin beams and equally populated noiseless squeezed states ($\Delta\theta = 0$) in forming the output state at the beam splitter with phase ϕ . Such state is generated by the Hamiltonian (2) assuming $g_{11} = g_{22} = g$ and described by the following elements of the evolution matrices \mathbf{U} and \mathbf{V} :

$$\begin{aligned} U_{11} &= U_{22} = \cosh(g_{12}t) \cosh(2gt), \\ V_{11} &= V_{22} = i \cosh(g_{12}t) \sinh(2gt), \\ U_{12} &= U_{21} = \sinh(g_{12}t) \sinh(2gt), \\ V_{12} &= V_{21} = i \sinh(g_{12}t) \cosh(2gt). \end{aligned} \quad (41)$$

Introducing the mean photon-pair number B_p as $B_p = \sinh^2(g_{12}t)$ and mean number \tilde{B}_p of squeezed photons per mode, $\tilde{B}_p = \sinh^2(2gt)$, the coefficients of the covariance

matrix $\mathbf{A}_{\mathcal{N}}$ are found in the form

$$\begin{aligned} B_1 &= B_2 = B_p + \tilde{B}_p + 2B_p\tilde{B}_p, \\ C_1 &= C_2 = i\sqrt{\tilde{B}_p(\tilde{B}_p + 1)(2B_p + 1)}, \\ D_{12} &= i\sqrt{B_p(B_p + 1)(2\tilde{B}_p + 1)}, \\ \tilde{D}_{12} &= -2\sqrt{B_p(B_p + 1)\tilde{B}_p(\tilde{B}_p + 1)}. \end{aligned} \quad (42)$$

The local NIs $I_{\text{ncl}}^{(j)}$, EI I_{ent} , and global NI I_{ncl} are then derived as follows:

$$\begin{aligned} I_{\text{ncl}}^{(1,2)} &= [1 - 4TR \sin^2(\phi)]\tilde{B}_p(\tilde{B}_p + 1) + 4TRB_p(B_p + 1) \\ &\quad - (\tilde{B}_p - B_p)^2 \pm K, \\ K &= 4\sqrt{TR} \cos(\phi) \sqrt{B_p(B_p + 1)\tilde{B}_p(\tilde{B}_p + 1)}, \\ I_{\text{ent}} &= (T - R)^2 B_p(B_p + 1) + 4TR \sin^2(\phi)\tilde{B}_p(\tilde{B}_p + 1), \\ I_{\text{ncl}} &= 2(B_p + \tilde{B}_p + 2B_p\tilde{B}_p). \end{aligned} \quad (43)$$

The formula for the global NI I_{ncl} , given in Eq. (43), shows that both parametric down-conversion and second-subharmonic generation contribute to the global NI making I_{ncl} always positive. Moreover, both processes enhance each other in producing larger values of the global NI. The greater is the mean photon-pair number B_p and also the greater is the mean number \tilde{B}_p of squeezed photons, the greater is the global NI I_{ncl} (see Fig. 12). Additionally, both LNI $I_{\text{ncl}}^{(j)}$ and EI I_{ent} become dependent on the phase ϕ of the beam splitter.

Provided that the phases of the incident squeezed states equal ($\phi = n\pi$, $n \in \mathbb{Z}$), photons in pairs stick together (bunch) completely when propagating through the beam splitter and so they cannot contribute to the entanglement in the output state. In this case, the entanglement originates only in photon pairs of the incident twin beam. When $T = \frac{1}{2}$, all photons

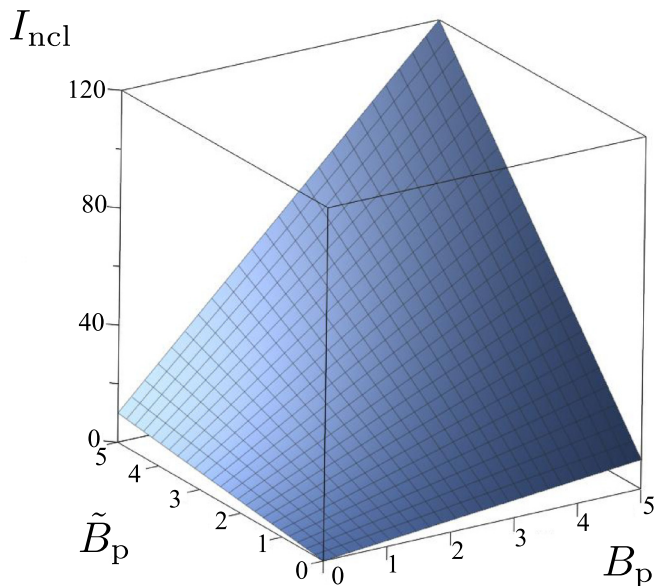


FIG. 12. Global nonclassicality invariant I_{ncl} as a function of the mean photon-pair number B_p and mean number \tilde{B}_p of squeezed photons considering the noiseless twin beams and squeezed states.

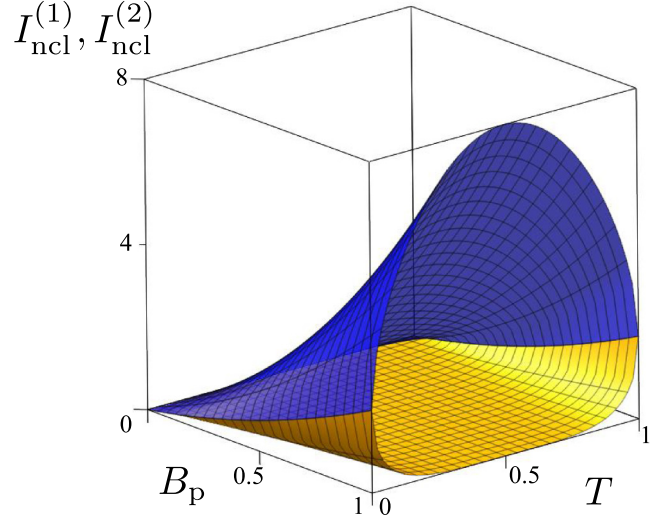


FIG. 13. Local nonclassicality invariants $I_{\text{ncl}}^{(1)}$ (blue dark upper surface) and $I_{\text{ncl}}^{(2)}$ (orange light colored lower surface) versus the beam-splitter transmissivity T and mean photon-pair number B_p assuming $B_p = \tilde{B}_p$ appropriate for the noiseless twin beams and squeezed states according to Eq. (43) assuming $\phi = 0$.

in pairs from the twin beam are glued and so the output state is separable. Contrary to this, the local NIs $I_{\text{ncl}}^{(j)}$ depend on both mean photon-pair number B_p and mean number \tilde{B}_p of squeezed photons. The fields characterizing photon pairs in individual output ports and originating in the incident squeezed states and the incident twin beam interfere, causing the asymmetry between the output ports. Depending on the parity of n one obtains the maximal local NI $I_{\text{ncl}}^{(1)}$ ($I_{\text{ncl}}^{(2)}$) if $n = 2k$ ($n = 2k + 1$), $k \in \mathbb{Z}$. This asymmetry is the largest for $T = \frac{1}{2}$. In this case, the bunched photon pairs are completely missing in one output port due to completely destructive interference. On the other hand, constructive interference provides the greatest number of the bunched photon pairs in the other output port, guaranteeing the largest attainable value of its local NI $I_{\text{ncl}}^{(j)}$. This behavior is quantified in the graph in Fig. 13.

If $\phi = \frac{\pi}{2} + n\pi$, the local NIs are equal ($I_{\text{ncl}}^{(1)} = I_{\text{ncl}}^{(2)}$) and the state at the beam-splitter output ports acquires a symmetry. Under these phase relations, also the incident squeezed photon pairs contribute, together with the twin-beam photon pairs, to the entanglement. It is worth noting that for $B_p = \tilde{B}_p$ all the state quantifiers are the same: $I_{\text{ncl}}^{(1)} = I_{\text{ncl}}^{(2)} = I_{\text{ent}} = B_p(B_p + 1)$.

VII. CONCLUSIONS

Local and global invariants of the general two-mode Gaussian states have been used to construct a specific local nonclassicality quantifier and entanglement quantifier. These quantifiers applied, respectively, to the single-mode marginal states and the whole two-mode state add together to give a quantity that is invariant under global linear unitary transformations. This invariant then quantifies the nonclassicality resources of Gaussian states. Remarkably, this invariant is linearly proportional to the number of photon pairs in the

noiseless Gaussian states. The general results have been used to study the beam-splitter transformations of fields composed of photon pairs and additional noisy photons. Twin beams, squeezed states, as well as their combinations have been considered as important examples. The behavior of photon pairs at the beam splitter causing their breaking or gluing (i.e., antibunching or bunching) has been used to explain the flow of nonclassical resources between local nonclassicalities (implying squeezing) and entanglement. A complete transfer of the entanglement of incident twin beams into the squeezing of the output modes has been observed. Also, the complete transfer of the incident squeezing into the entanglement of the output fields can be reached. The role of noise in the transfer of the nonclassicality invariant via the beam splitter has been elucidated on several examples.

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APPENDIX: QUASIPROBABILITY DISTRIBUTIONS AND CHARACTERISTIC FUNCTIONS

For the completeness and clarity of our presentation, here we give a few well-known formulas relating the quantities given in Table II, as derived by Cahill and Glauber [53]. This approach is a generalization of the standard Wigner and Glauber formalisms.

The Cahill-Glauber s -parametrized (or s -ordered) quasiprobability distribution (QPD, quasidistribution), $\mathcal{W}^{(s)}(\boldsymbol{\alpha})$ for an N -mode bosonic state $\hat{\rho}$ can be defined for a real parameter $s \in [-1, 1]$ as

$$\mathcal{W}^{(s)}(\boldsymbol{\alpha}) = \text{Tr}[\hat{\rho} \hat{T}^{(s)}(\boldsymbol{\alpha})], \quad (\text{A1})$$

which is the mean value of the operator $\hat{T}(\boldsymbol{\alpha})$ defined as the Fourier transform

$$\hat{T}^{(s)}(\boldsymbol{\alpha}) = \int \hat{D}^{(s)}(\boldsymbol{\beta}) \exp\left(\sum_n \alpha_n \beta_n^* - \text{c.c.}\right) d^2 \boldsymbol{\beta}' \quad (\text{A2})$$

of the s -parametrized multimode displacement operator given by

$$\begin{aligned} \hat{D}^{(s)}(\boldsymbol{\beta}) &= \prod_n \hat{D}^{(s)}(\beta_n) \\ &= \prod_n \exp\left(\beta_n \hat{a}_n^\dagger - \beta_n^* \hat{a}_n + \frac{s}{2} |\beta_n|^2\right). \end{aligned} \quad (\text{A3})$$

Here, \hat{a}_n (\hat{a}_n^\dagger) is the bosonic annihilation (creation) operator for the n th mode ($n = 1, 2, \dots, N$). The complex multivariable $\boldsymbol{\alpha} \equiv \{\alpha_n\} = (\alpha_1, \alpha_2, \dots, \alpha_N)$ is applied here as in Eq. (1), and, analogously $\boldsymbol{\beta} \equiv \{\beta_n\} = (\beta_1, \beta_2, \dots, \beta_N)$. The symbol c.c. denotes the complex conjugate term, and the integration is performed over $d^2 \boldsymbol{\beta}' \equiv d^2 \{\beta_n/\pi\} = \pi^{-N} \prod_n d^2 \beta_n = \pi^{-N} \prod_n d(\text{Re} \beta_n) d(\text{Im} \beta_n)$.

In the special cases for $s = 1, 0, -1$, the QPD $\mathcal{W}^{(s)}(\boldsymbol{\alpha})$ reduces to the popular Glauber-Sudarshan P , Wigner W , and Husimi Q functions corresponding to the normal, symmetric, and antinormal orderings, respectively. Our analysis in the paper is focused on the normally and symmetrically ordered functions. We recall that the standard definition of nonclassicality is based on the nonpositivity of the P function.

The statistical operator $\hat{\rho}$ corresponding to a given QPD can be calculated as follows:

$$\hat{\rho} = \int \mathcal{W}^{(s)}(\boldsymbol{\alpha}) \hat{T}^{(-s)}(\boldsymbol{\alpha}) d^2 \boldsymbol{\alpha}', \quad (\text{A4})$$

which, in the special case for $s = 1$, reduces to Eq. (1) describing the P representation of a given state $\hat{\rho}$.

The N -mode s -parameterized characteristic function $\mathcal{C}^{(s)}(\boldsymbol{\beta})$ for a given state $\hat{\rho}$ can be defined as

$$\mathcal{C}^{(s)}(\boldsymbol{\beta}) = \text{Tr}[\hat{\rho} \hat{D}^{(s)}(\boldsymbol{\beta})], \quad (\text{A5})$$

which is the mean value of the multimode displacement operator $\hat{D}^{(s)}(\boldsymbol{\beta})$. We recall that our description of nonclassicality and entanglement is based on two special cases of these characteristic functions. Specifically, the normal characteristic function C_N , defined in Eq. (7), is the special case of Eq. (A5) for $s = 1$ assuming two-mode ($N = 2$) field. While the symmetric characteristic function C_S is given by Eq. (A5) for $s = 0$ and $N = 2$.

For Gaussian states, which are solely analyzed in this paper, the characteristic function $\mathcal{C}^{(s)}(\boldsymbol{\beta})$ can uniquely be defined via the covariance matrices, which are given by Eq. (10) for normal ordering ($s = 1$) and by Eq. (17) for symmetric ordering ($s = 0$) for a two-mode case.

By comparing the definitions in Eqs. (A1) and (A5), it is easy to conclude that the s -parameterized QPD and characteristic function for any $s \in [-1, 1]$ are related via the Fourier transform, i.e.,

$$\mathcal{W}^{(s)}(\boldsymbol{\alpha}) = \int \mathcal{C}^{(s)}(\boldsymbol{\beta}) \prod_n \exp(\alpha_n \beta_n^* - \alpha_n^* \beta_n) d^2 \boldsymbol{\beta}', \quad (\text{A6})$$

$$\mathcal{C}^{(s)}(\boldsymbol{\beta}) = \int \mathcal{W}^{(s)}(\boldsymbol{\alpha}) \prod_n \exp(\alpha_n^* \beta_n - \alpha_n \beta_n^*) d^2 \boldsymbol{\alpha}', \quad (\text{A7})$$

where the integration over $d^2 \boldsymbol{\alpha}'$ is defined analogously to $d^2 \boldsymbol{\beta}'$, as in Eq. (A2). The normalization conditions are as follows:

$$\int \mathcal{W}^{(s)}(\boldsymbol{\alpha}) d^2 \boldsymbol{\alpha}' = \mathcal{C}^{(s)}(\boldsymbol{\beta} = \mathbf{0}) = 1. \quad (\text{A8})$$

The relation between the QPDs $\mathcal{W}^{(s_1)}(\boldsymbol{\alpha})$ and $\mathcal{W}^{(s_2)}(\boldsymbol{\alpha})$, assuming $s_2 < s_1$, is simply given by

$$\begin{aligned} \mathcal{W}^{(s_2)}(\boldsymbol{\alpha}) &= \left(\frac{2}{s_1 - s_2}\right)^M \int \mathcal{W}^{(s_1)}(\boldsymbol{\beta}) \\ &\times \exp\left(-\frac{2}{s_1 - s_2} \sum_n |\alpha_n - \beta_n|^2\right) d^2 \boldsymbol{\beta}'. \end{aligned} \quad (\text{A9})$$

This means that the QPD $\mathcal{W}^{(s_2)}(\boldsymbol{\alpha})$ with any parameter $s_2 \in [-1, 1]$ can easily be obtained by mixing the P function (corresponding to $s_1 = 1$) with the proper amount of

Gaussian noise. The relation between the characteristic functions corresponding to different parameters s_1 and s_2 reads as

$$\mathcal{C}^{(s_2)}(\boldsymbol{\beta}) = \mathcal{C}^{(s_1)}(\boldsymbol{\beta}) \exp\left(\frac{s_2 - s_1}{2} \sum_n |\beta_n|^2\right). \quad (\text{A10})$$

It is valid for any $s_1, s_2 \in [-1, 1]$, contrary to the analogous relation in Eq. (A9) for the QPDs. We applied Eq. (A10) to calculate the symmetrically ordered characteristic function $C_S \equiv \mathcal{C}^{(0)}$ from the normally ordered characteristic function $C_N \equiv \mathcal{C}^{(1)}$, given by Eq. (7), for two-mode states.

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