Non-Gaussian entanglement distillation for continuous variables

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Entanglement distillation is an essential ingredient for long distance quantum communications1. In the continuous variable setting, Gaussian states play major roles in quantum teleportation, quantum cloning and quantum cryptography2. However, entanglement distillation from Gaussian states has not yet been demonstrated. It is made difficult by the no-go theorem stating that no Gaussian operation can distill Gaussian states3,4,5. Here we demonstrate the entanglement distillation from Gaussian states by using measurement-induced non-Gaussian operations, circumventing the fundamental restriction of the no-go theorem. We observed a gain of entanglement as a result of conditional local subtraction of a single photon or two photons from a two-mode Gaussian state. Furthermore we confirmed that two-photon subtraction also improves Gaussian-like entanglement as specified by the Einstein-Podolsky-Rosen (EPR) correlation. This distilled entanglement can be further employed to downstream applications such as high fidelity quantum teleportation6 and a loophole-free Bell test7.

Long distance quantum communications rely on the ability to faithfully distribute entanglement between distant locations. However, inevitable decoherence and the inability to amplify quantum signals hinder efforts to extend a quantum optical link to the practically large scale. To overcome this difficulty, entanglement distillation can be used - a protocol in which each distant party locally manipulates particles of less entangled pairs with the aid of classical communication to extract a smaller number of pairs of higher entanglement8. In discrete variable systems many distillation experiments have already been demonstrated9,10,11,12.

An alternative to the discrete variable system is the one described by continuous variables (CV), typically represented by the quadrature amplitudes of a light field. For CV quantum information, Gaussian states and Gaussian operations13 are of particular importance. They are readily available in the laboratory and serve as a complete framework for many quantum protocols2. The first experiments of CV entanglement distillation were reported in14,15. These two works rely on Gaussian operations. However, it was theoretically proven that Gaussian operations can never distill entanglement from Gaussian state inputs – this is known as the no-go theorem of Gaussian operations3,4,5. In14,15, the inputs had been subject to some specific classes of non-Gaussian noise, such as phase-diffusion14 or temporally varying attenuation15. In those cases, well established Gaussian technologies can be applied to distill the entanglement.

So far, there has been no demonstration of entanglement distillation with Gaussian inputs. This task essentially requires a new technology of non-Gaussian operations. Recent theories also revealed that this is a must to realize quantum speed-up of CV quantum information processing (QIP)16. Triggered by this new paradigm of non-Gaussian QIP, the research field extending to the non-Gaussian regime has rapidly developed17,18,19. The increase20 and preparation21 of entanglement from Gaussian inputs by non-local photon subtraction were also demonstrated. These are important steps towards the realization of entanglement distillation from Gaussian states.

Here we report on the entanglement distillation directly from CV Gaussian states by using local photon subtraction as non-Gaussian operations, circumventing the no-go restriction on Gaussian operations. A schematic of our experiment is depicted in Fig.1. A continuous wave squeezed vacuum is generated from an optical parametric oscillator (OPO) detailed elsewhere19. The initial Gaussian entangled state is prepared by splitting the squeezed vacuum by half at the first beam splitter and is distributed to the separate parties, Alice and Bob. This half-split squeezed vacuum with squeezing parameter r is effectively equivalent to the two-mode squeezed vacuum with r/2 (they are compatible by local unitary operations, see Appendix A2). At each site of Alice and Bob, a probabilistic non-Gaussian operation – photon subtraction – is performed. Specifically, a small part of the beam is picked off by a polarizing beam splitter with the variable reflectance R and sent through filtering cavities19 to an avalanche photodiode (APD) to detect a photon (Fig.1). Each photon detection at the APD heralds a local success of the photon subtraction attempt. Conditioned on the subtraction of a photon by a single party (single-photon subtraction) or the simultaneous subtraction of a photon by both parties (two-photon subtraction), Alice and Bob retain those two-mode states which have successfully had their entanglement increased. While the single-photon subtraction scheme will have a higher success rate, the two-photon subtraction scheme will give a more Gaussian-like final
state which is more readily applicable to further processing such as e.g. quantum teleportation.

The distillation works since the local photon subtraction changes the non-local unfactorizable correlations of the initial state. To see this intuitively, let us describe the initial state. To see this intuitively, let us describe the initial state. To see this intuitively, let us describe the initial state. To see this intuitively, let us describe the initial state. To see this intuitively, let us describe the

\[
W(x_A, p_A, x_B, p_B) = W_0(x_+, p_+) W_0(x_-, p_-),
\]

where \(W\) is the Wigner function for the two-mode state, \(x_A, x_B, p_A, p_B\) are the quadrature amplitudes of modes \(A\) and \(B\), \(x_\pm = \frac{x_A + x_B}{\sqrt{2}}, p_\pm = \frac{p_A + p_B}{\sqrt{2}}\) and \(W_0\) are the Wigner functions for a (zero-, one-, or two-) photon subtracted squeezed vacuum and the vacuum respectively. For such a state, the scans of the homodyne measurements are necessary only for \(\theta_A = \theta_B\) and the experimental data of \(x_\pm\) is numerically obtainable from the measured \(x_A\) and \(x_B\) (Fig. 1). It should be stressed that although we assume the state factorization Eq. (1), it can be directly assessed by verifying experimentally whether the state of the g+h mode is indeed a pure vacuum state. For details, see Appendix A 1.

Examples of reconstructed Wigner functions obtained by the single- and two-photon subtractions, as well as the initial squeezed state are shown in Fig. 2a-c. The outputs of the homodyne detectors were sampled at 6 different phases of LO, namely \(\theta_A = \theta_B = 0, \pi/6, \pi/3, \pi/2, 2\pi/3, 5\pi/6\). We extract the measured values of the quadratures \(\hat{x}_A\) and \(\hat{x}_B\) by applying a mode function to the recorded traces \(18, 19, 22\). After calculating the corresponding values of \(\hat{x}_\pm\), we reconstruct the density matrices for the “+” and “−” modes by the conventional maximum likelihood estimation \(23\) without any correction of detection losses.

As shown in Fig. 2, for the “+” mode states we got almost perfectly pure vacuum states with more than 99% accuracy. We confirmed that this holds irrespective of the initial squeezing level. This experimental evidence justifies our tomography scheme based on the relation Eq. (1). On the other hand, for the “−” mode states we observed two different kinds of non-Gaussian state depending on whether single photon or two photons were subtracted (Fig. 2a and b). They respectively correspond to the odd and even Schrödinger cat state, i.e. \(|\alpha\rangle - |-\alpha\rangle\) and \(|\alpha\rangle + |-\alpha\rangle\) where \(|\alpha\rangle\) is a coherent state with coherent amplitude \(\alpha\). Having these reconstructed states we can use the relation Eq. (1) backwards to calculate the amount of entanglement shared by Alice and Bob. Specifically, we calculate the logarithmic negativity \(E_N\) which is a monotone measure of entanglement \(24\).

Fig. 2 shows the experimental negativities of the undistilled Gaussian states, the states distilled by single-photon subtraction with \(R = 5\%\), and by two-photon subtraction with \(R = 10\%\) as functions of the squeezing of the initial input states. When evaluating negativity, one must take care of its strong dependency on the size of the data set. We investigated the behavior of the negativity on the data size and deduced an extrapolative value corresponding to an infinitely large data set for each point in Fig. 2 (see Appendix B). Note that without this analysis, evaluation of negativity with finite sized data very likely goes into an overestimate of the negativity. As shown in the figure, over a wide range of the initial squeezing we got clear gains of entanglement.
A practical difference between the single- and two-photon subtracted scheme is on their rates of event detection. In the single-photon experiment the rate is around a few thousands per second, but in the two-photon experiment there are only a few events per second. So while for the former we can use hundreds of thousands of samples for the state reconstruction, for the latter we can only use a few tens of thousands limited by the long-term stability of the setup. In Fig. 2d the experimental logarithmic negativities for the single-photon subtracted states and the undistilled states are in very good agreement with theory, but ones for the two-photon subtraction are slightly below the theoretical predictions. This may be due to an uncontrollable drift of the system during a long period of the measurements.

As can be seen in Fig. 2d, in terms of the logarithmic negativity the two-photon subtracted scheme does not have an advantage over the single-photon subtracted scheme despite its significantly lower success rate. However the two-photon subtracted distillation transforms a two-mode Gaussian state into one relatively close to a Gaussian state (see Fig. 2b). Hence one would expect that states distilled by this scheme still possess a Gaussian-like property of entanglement. For Gaussian states, two-mode entanglement is usually specified in terms of the Einstein-Podolsky-Rosen (EPR) correlation quantified by $\langle (\Delta x_-)^2 \rangle \langle (\Delta p_+)^2 \rangle$. Since the “+” mode is always a vacuum state (see Eq. 1), we can focus on the degree of squeezing of the “−” mode as an equivalent measure. We carried out measurements of the
FIG. 3: Squeezed variances of $x_-$ (normalized by the vacuum level). For the undistilled states each point was calculated from the full-reconstructed density matrices common to the negativity measurements. For the distilled states each point was obtained by directly measuring variance of 1,600 - 5,000 samples at the most squeezed phase.

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A and the vacuum states for modes B, C, and D:

\[
\hat{B}_{BD}(\theta)\hat{B}_{AC}(\theta)\hat{B}_{AB}(\pi/4) |\Psi\rangle_A |0\rangle_B |0\rangle_C |0\rangle_D \\
= \hat{B}_{AB}(\pi/4)\hat{B}_{CD}(\pi/4)\hat{B}_{AC}(\theta)\hat{B}_{BD}(\theta)\hat{B}_{CD}^\dagger(\pi/4) \\
\times |\Psi\rangle_A |0\rangle_B |0\rangle_C |0\rangle_D \\
= \hat{B}_{AB}(\pi/4)\hat{B}_{CD}(\pi/4)\hat{B}_{AC}(\theta) |\Psi\rangle_A |0\rangle_B |0\rangle_C |0\rangle_D \\
\]

This establishes the equivalency between the models shown in Fig. 4 and b again. Therefore the final output state of the protocol is identical to a half-split of the photon subtracted squeezed vacuum.

Let us return to the form of (A2) for simplicity. Then we immediately have

\[
\hat{B}^\dagger(\pi/4)|\psi_{out}\rangle_{AB} = \hat{a}^{-n_A+n_B}_{A}S_{A}(r)|0\rangle_{A} |0\rangle_{B},
\]

which means that if we let the conditional two-mode state be combined at a half-beam splitter it becomes disentangled and furthermore the outputs get separated as a photon subtracted squeezed vacuum and the vacuum.

Introducing new variables

\[
\hat{x}_\pm = \frac{\hat{x}_A \pm \hat{x}_B}{\sqrt{2}}, \quad \hat{p}_\pm = \frac{\hat{p}_A \pm \hat{p}_B}{\sqrt{2}}, \quad (A10)
\]

where \(\hat{x}_{A,B}\) and \(\hat{p}_{A,B}\) are the quadrature observables for Alice and Bob’s subsystem, then from the identity (A9), in terms of these variables the Wigner function of the output state can be written as

\[
W_{out}(x_A, p_A, x_B, p_B) = W_{c}(x_+, p_+)W_{s}(x_-, p_-), \quad (A11)
\]

where \(W_{c}\) and \(W_{s}\) are the Wigner functions of the vacuum state and the photon subtracted squeezed vacuum state respectively.

2. Local unitary equivalency between a half-split squeezed vacuum and a two-mode squeezed vacuum

A two-mode Gaussian state with zero local displacements is completely specified by its covariance matrix given by

\[
\begin{pmatrix}
\langle \hat{x}_A^2 \rangle & \frac{1}{2}\langle \hat{x}_A \hat{p}_A + \hat{p}_A \hat{x}_A \rangle \\
\frac{1}{2}\langle \hat{x}_A \hat{p}_A + \hat{p}_A \hat{x}_A \rangle & \langle \hat{p}_A^2 \rangle \\
\frac{1}{2}\langle \hat{x}_B \hat{p}_B + \hat{p}_B \hat{x}_B \rangle & \frac{1}{2}\langle \hat{x}_B \hat{p}_B + \hat{p}_B \hat{x}_B \rangle & \langle \hat{p}_B^2 \rangle \\
\frac{1}{2}\langle \hat{x}_A \hat{p}_B + \hat{p}_B \hat{x}_A \rangle & \frac{1}{2}\langle \hat{x}_B \hat{p}_A + \hat{p}_A \hat{x}_B \rangle & \langle \hat{p}_A^2 \rangle \\
\frac{1}{2}\langle \hat{x}_A \hat{p}_B + \hat{p}_B \hat{x}_A \rangle & \frac{1}{2}\langle \hat{x}_B \hat{p}_A + \hat{p}_A \hat{x}_B \rangle & \langle \hat{p}_B^2 \rangle \\
\end{pmatrix}
(A12)
\]

The covariance matrix of the half-split squeezed vacuum \(\langle \Psi_0 \rangle = \hat{B}_{AB}(\pi/4)S_{A}(r)|0\rangle_{A} |0\rangle_{B}\) is

\[
\begin{pmatrix}
\frac{e^{-r}}{2} \cosh(r) & 0 & \frac{e^{-r}}{2} \sinh(r) \\
0 & 2 & 0 \\
\frac{e^{-r}}{2} \sinh(r) & 0 & \frac{e^{-r}}{2} \cosh(r) \\
0 & 0 & 2
\end{pmatrix}
\]

Performing local squeezing operations on both modes, it can be made to have symmetric variances:

\[
\begin{pmatrix}
\cosh(r) & 0 & \sinh(r) & 0 \\
0 & 2 & 0 & -\sinh(r) \\
\sinh(r) & 0 & 2 & 0 \\
0 & -\sinh(r) & 0 & 2
\end{pmatrix}
\]
This is identical to the covariance matrix of a two-mode squeezed state with squeezing parameter $\frac{r}{\sqrt{2}}$. So in terms of entanglement, the half-split squeezed vacuum with squeezing parameter $r$ is equivalent to the two-mode squeezed state with squeezing parameter $\frac{r}{\sqrt{2}}$. By using a state vector, this is described as

$$\hat{S}_A(-r/2)\hat{S}_B(-r/2)|\Psi_0\rangle = (1 - \lambda^2)\frac{1}{2}\sum_{n=0}^{\infty} \lambda^n |n\rangle_A |n\rangle_B,$$

(A13)

with $\lambda = \tanh r/2$.

3. Entanglement of the distilled states

Entropy of entanglement is defined as the von Neumann entropy of a reduced subsystem. Every bipartite pure state can be brought into the form of the Schmidt decomposition with Schmidt coefficients $c_n$:

$$\sum_{n=0}^{\infty} \sqrt{c_n} |n\rangle_A |n\rangle_B.$$

(A14)

From this, the entropy of the subsystem is calculated as

$$E = -\text{tr}(\rho_A \log(\rho_A)) = -\text{tr}(\rho_B \log(\rho_B))$$

(A15)

$$= -\sum_n c_n \log c_n.$$  

(A16)

From (A13) and the fact that local operations do not alter the amount of entanglement, we get the Schmidt coefficients of the half-split squeezed vacuum $|\Psi_0\rangle$ as

$$c_n^{(0)} = (1 - \lambda^2)\lambda^{2n}.$$  

(A17)

Let us denote the single-photon subtracted half-split squeezed vacuum as $|\Psi_1\rangle$.

$$|\Psi_1\rangle = \frac{1}{\sqrt{\mathcal{N}_1}}\hat{a}_A |\Psi_0\rangle,$$  

(A18)

$$\mathcal{N}_1 = \sinh^2 r/2.$$  

(A19)

From the argument in the last section, we have

$$\hat{S}_A(-r/2)\hat{S}_B(-r/2)|\Psi_1\rangle$$

$$= \frac{1}{\sqrt{\mathcal{N}_1}}\hat{S}_A(-r/2)\hat{S}_B(-r/2)\hat{a}_A |\Psi_0\rangle$$

(A20)

$$= \frac{1}{\sqrt{\mathcal{N}_1}}(\hat{a}_A \cosh(r/2) + \hat{a}_A^\dagger \sinh(r/2))$$

$$\times \hat{S}_A(-r/2)\hat{S}_B(-r/2)|\Psi_0\rangle$$

(A21)

$$= \frac{1}{\sqrt{\mathcal{N}_1}}(\hat{a}_A \cosh(r/2) + \hat{a}_A^\dagger \sinh(r/2))(1 - \lambda^2)\frac{1}{2}$$

$$\times \sum_{n=0}^{\infty} \lambda^n |n\rangle_A |n\rangle_B$$

(A22)

$$= \frac{(1 - \lambda^2)\frac{1}{2}}{\sqrt{\mathcal{N}_1}} \sum_{n=0}^{\infty} \lambda^n (\cosh(r/2)\sqrt{n} |n - 1\rangle_A$$

$$+ \sinh(r/2)\sqrt{n + 1} |n + 1\rangle_A) |n\rangle_B$$

(A23)

$$= \frac{(1 - \lambda^2)\frac{1}{2}}{\sqrt{\mathcal{N}_1}} \sum_{m=0}^{\infty} \sum_{n=0}^{\infty} A_{mn} |m\rangle_A |n\rangle_B,$$  

(A24)

where

$$A_{mn} = \lambda^n (\cosh(r/2)\sqrt{n} \delta_{m,n-1}$$

$$+ \sinh(r/2)\sqrt{n + 1} \delta_{m,n+1}).$$

(A25)

From the singular values of matrix $A_{mn}$ (let them be $\alpha_n$), the Schmidt coefficients for $|\Psi_1\rangle$ are given by

$$c_n^{(1)} = \frac{(1 - \lambda^2)\frac{1}{2}}{\sqrt{\mathcal{N}_1}} \alpha_n.$$  

(A26)

Similar to the single photon case, we obtain a representation of the two-photon subtracted half-split squeezed vacuum $|\Psi_2\rangle$ in the following form.

$$\hat{S}_A(-r/2)\hat{S}_B(-r/2)|\Psi_2\rangle$$

$$= \frac{1}{\sqrt{\mathcal{N}_2}}\hat{S}_A(-r/2)\hat{S}_B(-r/2)\hat{a}_A^\dagger \hat{a}_B |\Psi_0\rangle$$

(A27)

$$= \frac{(1 - \lambda^2)\frac{1}{2}}{\sqrt{\mathcal{N}_2}} \sum_{m=0}^{\infty} \sum_{n=0}^{\infty} B_{mn} |m\rangle_A |n\rangle_B,$$  

(A28)

where

$$\mathcal{N}_2 = 2 \sinh^4 r + \cosh^2 r \sinh^2 r$$

(A29)

$$B_{mn} = \cos^2(r/2)(m + 1)\lambda^{m+1} \delta_{m,n}$$

$$+ \sinh^2(r/2)\lambda^{m+1} \delta_{m+1,n}$$

$$+ \cosh(r/2) \sinh(r/2) \sqrt{(m + 1)(m + 2)} \lambda^{m+1} \delta_{m+2,n}$$

$$+ \cosh(r/2) \sinh(r/2) \sqrt{m(m - 1)} \lambda^{m+1} \delta_{m-2,n}$$

(A30)
Then the Schmidt coefficients $c_n^{(2)}$ for $|\Psi_2\rangle$ are given by the singular values $\beta_n$ of matrix $B_{mn}$:

$$c_n^{(2)} = \frac{(1 - \lambda_n^2)\beta_n^2}{N_2}$$  \hspace{1cm} (A31)

and we can readily calculate the entropy of entanglement for those states. Fig. 5 shows the entropy of entanglement for the single-photon subtracted states $|\Psi_1\rangle$, the two-photon subtracted states $|\Psi_2\rangle$ and the undistilled states $|\Psi_0\rangle$ as functions of squeezing parameter $r$. Note that in an actual experiment, we have several experimental imperfections and we end up with a mixed state output. In such cases, the pure state descriptions above no longer hold and the entropy of entanglement is not a good measure of entanglement, but Fig. 5 still outlines the general behavior of this protocol.

**APPENDIX B: CALCULATING THE LOGARITHMIC NEGATIVITY**

In principle, the logarithmic negativity $E_N$ can be directly calculated when we know the 2-mode entangled state, as the sum of the negative eigenvalues of the partially transposed density matrix, $(\rho_{AB})^{\text{TP}}$ [24]. In the experiment, however, we found that $E_N$ is very sensitive to statistical noise in the measurements. Smaller data sets lead to larger errors in the reconstructed density matrix elements which ultimately leads to larger calculated $E_N$ values. The intuitive understanding of this effect is the following: The absolute errors on each density matrix element due to statistical measurement noise are roughly the same. Hence, the relative errors are large for the high-photon number elements which are all close to zero - most likely their absolute values will increase due to the errors. But high photon numbers contribute a significant amount to the overall entanglement of the state, so in the end more noise will give seemingly higher entanglement.

We found empirically that the calculated negativity scales with the total data sample size $N$ as $E_N(N) = a + b/\sqrt{N}$. We interpret this as $E_N(\infty) = a$ being the “true” value that we would obtain in the asymptotic limit of very large data sample size, while the second term is the contribution from statistical noise. To obtain this value $E_N(\infty)$ from a given data set of $N_{\text{full}}$ samples, we perform multiple state reconstructions based on truncations of the full data set. Specifically, we partition the full set into $d$ subsets of $N_d \approx N_{\text{full}}/d$ samples each (with equal representation of all phase angles). The logarithmic negativity is then calculated from the reconstruction of all $d$ subsets, and we take the mean value of these to be an estimate for $E_N(N_d)$. We repeat the process for other numbers of partitions, $d$, and thereby get a plot as in Fig. 6 of the dependency of calculated entanglement on data sample size. A least-squares fit to $a + b/\sqrt{N}$ then gives us the asymptotic estimate $E_N(\infty)$. We have confirmed by simulated data that this approach does in fact give the correct value for $E_N$ within roughly ±2%.

![FIG. 6: Estimation of the logarithmic negativity for infinitely many data samples. Each point shows the average and standard deviation of the calculated logarithmic negativities, $E_N(N_d)$ (average taken over the $d$ subsets of the full data set), versus $1/\sqrt{N_d}$, where $N_d$ is the size of each subset. The points are well fitted by a line whose $y$-axis intersection (infinite data size) gives a good estimate for the true $E_N$ value (as confirmed by simulated data). The three plot series are from single-photon subtracted data sets with different initial squeezing levels, 2.3 (red), 3.2 (blue), and 4.2 (green) dB respectively.]

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