

Four-photon subtraction from picosecond squeezed vacuum demonstrating Wigner negativity without loss correction

Mamoru Endo^{1,2,*}, Takefumi Nomura,¹ Tatsuki Sonoyama,¹ Kazuma Takahashi¹, Sachiko Takasu,³ Daiji Fukuda^{3,4}, Takahiro Kashiwazaki⁵, Asuka Inoue⁵, Takeshi Umeki,⁵ Rajveer Nehra^{1,6,7,8}, Petr Marek⁹, Radim Filip⁹, Kan Takase,^{1,2} Warit Asavanant,^{1,2} and Akira Furusawa^{2,1,†}

¹*Department of Applied Physics, School of Engineering, The University of Tokyo, 7-3-1 Hongo, Bunkyo, Tokyo 113-8656, Japan*

²*Optical Quantum Computing Research Team, RIKEN Center for Quantum Computing, 2-1 Hirosawa, Wako, Saitama 351-0198, Japan*

³*National Institute of Advanced Industrial Science and Technology, 1-1-1 Umezono, Tsukuba, Ibaraki 305-8563, Japan*

⁴*AIST-UTokyo Advanced Operando-Measurement Technology Open Innovation Laboratory,*

1-1-1 Umezono, Tsukuba, Ibaraki 305-8563, Japan

⁵*Device Technology Laboratories, NTT Inc., 3-1 Morinosato Wakamiya, Atsugi, Kanagawa 243-0198, Japan*

⁶*Department of Electrical and Computer Engineering, University of Massachusetts Amherst,*

100 Natural Resources Road, Amherst, Massachusetts 01003, USA

⁷*Department of Physics, University of Massachusetts Amherst, 100 Natural Resources Road, Amherst, Massachusetts 01003, USA*

⁸*College of Information and Computer Science, University of Massachusetts Amherst,*

100 Natural Resources Road, Amherst, Massachusetts 01003, USA

⁹*Department of Optics, Palacky University, 17 Listopadu 1192/12, Olomouc 77146, Czech Republic*



(Received 19 December 2025; revised 8 July 2026; accepted 13 July 2026; published 24 August 2026)

Multiphoton subtraction from a traveling squeezed vacuum is a standard route to non-Gaussian optical-state engineering. High state-generation rates call for ultrashort optical pulses driven by high-repetition-rate lasers, but combining higher-order subtraction with ultrashort pulsed light and full state reconstruction remains challenging. Here we demonstrate up to four-photon subtraction from 10-ps squeezed-vacuum wave packets generated by a broadband waveguide optical parametric amplifier. Using a Ti-Au transition-edge sensor for photon-number-resolved heralding and homodyne tomography on the signal mode, we reconstruct the generated states up to the four-photon-subtracted case. The reconstructed states exhibit clear parity-dependent structure and visible Wigner-function negativity without loss correction. These results extend the experimentally accessible regime of multiphoton subtraction in broadband traveling optical modes and establish a platform for non-Gaussian optical resource-state engineering in continuous-variable quantum optics.

DOI: [10.1103/PhysRevA.114.022618](https://doi.org/10.1103/PhysRevA.114.022618)

I. INTRODUCTION

Light has an exceptionally high frequency in the hundreds of terahertz. It serves not only as a fundamental carrier of information for ultrafast communication and data processing, but also as a promising platform for quantum technologies. Its ability to maintain quantum coherence under ambient conditions has positioned it as a key candidate for quantum computing [1–5], quantum sensing [6,7], and quantum communication [8]. The realization of ultrafast quantum processors by encoding quantum information into the continuous-variable (CV) quadrature-phase amplitudes of the electromagnetic field is a particularly promising direction for future research.

A major challenge in realizing fault-tolerant quantum computation, however, lies in the generation of logical qubits, which are necessary for error correction. In optical systems, the Gottesman-Kitaev-Preskill (GKP) qubit, characterized by a comblike wave-function structure in both quadratures, has emerged as a strong candidate for logical qubits [9]. Gottesman-Kitaev-Preskill qubits offer a significant advantage in that their fundamental operations, such as Clifford gates, can be implemented using only linear optics. Furthermore, universal quantum computation becomes achievable with these operations combined with ancilla GKP qubits [5,10,11]. Therefore, the successful realization of GKP qubits in optical systems is a key step toward fault-tolerant universal quantum computing. Several theoretical proposals have outlined methods to generate GKP qubits [12–16], with non-Gaussian resource states playing a central role. Among these, Schrödinger-cat states, that is, quantum superpositions of coherent states, have been widely studied as a versatile resource [17,18]. Different approaches have been proposed for preparing non-Gaussian resource states relevant to GKP-state preparation, including schemes that make use of multiple cat states as well as approaches based on multiphoton detection in integrated photonics. While proof-of-principle experiments

*Contact author: endo@ap.t.u-tokyo.ac.jp

†Contact author: akiraf@ap.t.u-tokyo.ac.jp

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI.

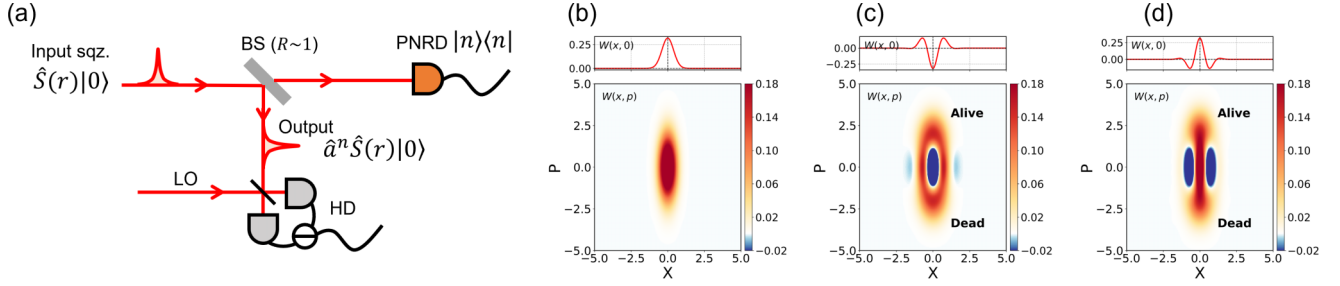


FIG. 1. (a) Conceptual diagrams of cat-state generation by photon subtraction from squeezed vacuum. Also plotted are the simulated Wigner functions $W(x, p)$ of (b) the input squeezed vacuum ($\hat{S}(0.576)|0\rangle$), that is, squeezing level of 5 dB and (c) the three- and (d) the four-photon-subtracted squeezed vacuum (proportional to $\hat{a}^3 \hat{S}(0.576)|0\rangle$ and $\hat{a}^4 \hat{S}(0.576)|0\rangle$ with $R = 0.78$, respectively).

have demonstrated the interference of two cat states [19], generating resource states with sufficiently large amplitudes and high rates remains experimentally demanding. Recent integrated-photonics demonstrations of GKP-like resource states using multiphoton detection [20] further underscore the growing interest in diverse optical routes to GKP-state preparation.

Photon subtraction [21] [Fig. 1(a)], a widely used method for generating Schrödinger-cat states in optical systems, is inherently probabilistic and thus subject to fundamental limitations. Nevertheless, it remains the most viable route to non-Gaussian state generation, as it requires only a deterministically generated squeezed-vacuum source, linear optics, and a photon-number-resolving detector, all of which are experimentally mature. In this approach, a portion of a squeezed-vacuum field is tapped by a beam splitter (BS), and conditional photon detection by a photon-number resolving detector (PNRD) in one output mode prepares a non-Gaussian state in the other mode that approximates a cat state. Single-photon subtraction has been studied extensively [22–25]; generating larger-amplitude states requires multiphoton subtraction, whose experimental realization remains challenging owing to limited photon-number-resolving performance, optical losses, phase instability, and rapidly decreasing count rates. Achieving high generation rates requires operation in picosecond or shorter optical wave packets together with high-repetition-rate laser systems. In this regime, the requirements for broadband operation and the management of optical losses become increasingly important, making it more difficult to preserve nonclassical features such as Wigner negativity. As a result, in the ultrashort-pulse regime, observations of Wigner-function negativity without loss correction have typically been limited to up to three-photon subtraction [26,27].

Here we demonstrate multiphoton subtraction by employing a broadband waveguide-based optical parametric amplifier [28] pumped with picosecond pulses to generate pulsed squeezed-vacuum fields with a wave-packet duration of about 10 ps. Using a high-speed, high-energy-resolution PNRD based on a Ti-Au superconducting transition-edge sensor (TES) [29], we successfully subtract up to four photons. The generation rate for three-photon subtraction was approximately 200 counts per second (cps), and even for four-photon subtraction, the rate was 1.5 cps. The resulting states exhibit clear nonclassical signatures, including Wigner-function negativity without loss correction and parity changes corresponding to the number of subtracted photons. Furthermore,

the experimentally obtained states are in close agreement with theoretical simulations. While multiphoton detection using TES has been demonstrated in various contexts, this work experimentally demonstrates four-photon subtraction in picosecond optical wave packets together with directly observed Wigner-function negativity without loss correction. These results highlight an experimentally accessible regime that combines ultrashort temporal modes with higher-order non-Gaussian operations, extending the capabilities of optical quantum state engineering beyond those achieved with spectrally narrowed or lower-rate systems.

The method established in this study is highly versatile and can be readily extended to next-generation quantum state engineering. By expanding the accessible regime of multiphoton subtraction, this work provides a basis for ultrafast non-Gaussian state preparation for applications in CV quantum information processing.

II. PHOTON SUBTRACTION FOR CAT-STATE GENERATION

A. Cat states

Coherent state superpositions hold particular significance in quantum optics. These states are often referred to as Schrödinger-cat states, in reference to the famous gedanken experiment [30]. In this analogy, the states of an “alive cat” and a “dead cat” correspond to the coherent states with amplitudes of α and $-\alpha$, respectively. The state vectors of the cat state are expressed as $|\pm \text{cat}\rangle \propto |\alpha\rangle \pm |-\alpha\rangle$, where $|+\text{cat}\rangle$ means even cat and $|-\text{cat}\rangle$ means odd cat. While GKP states offer stronger error-correction capabilities than the two-component cat states demonstrated here, their realization and practical use remain experimentally demanding. In contrast, Schrödinger-cat states are more readily generated and handled in current experimental platforms [31]. Owing to these properties, cat states serve as an important intermediate resource.

B. Photon subtraction

Approximate cat states can be generated by photon subtraction [21,26] as shown in Fig. 1(a). A squeezed-vacuum state $\hat{S}(r)|0\rangle$ is deterministically generated via an optical parametric process and split into two modes by a high-reflectivity BS with reflectivity of R (usually $R \sim 1$), where $\hat{S}(r)$ and r are a squeezing operator and a squeezing parameter, respectively.

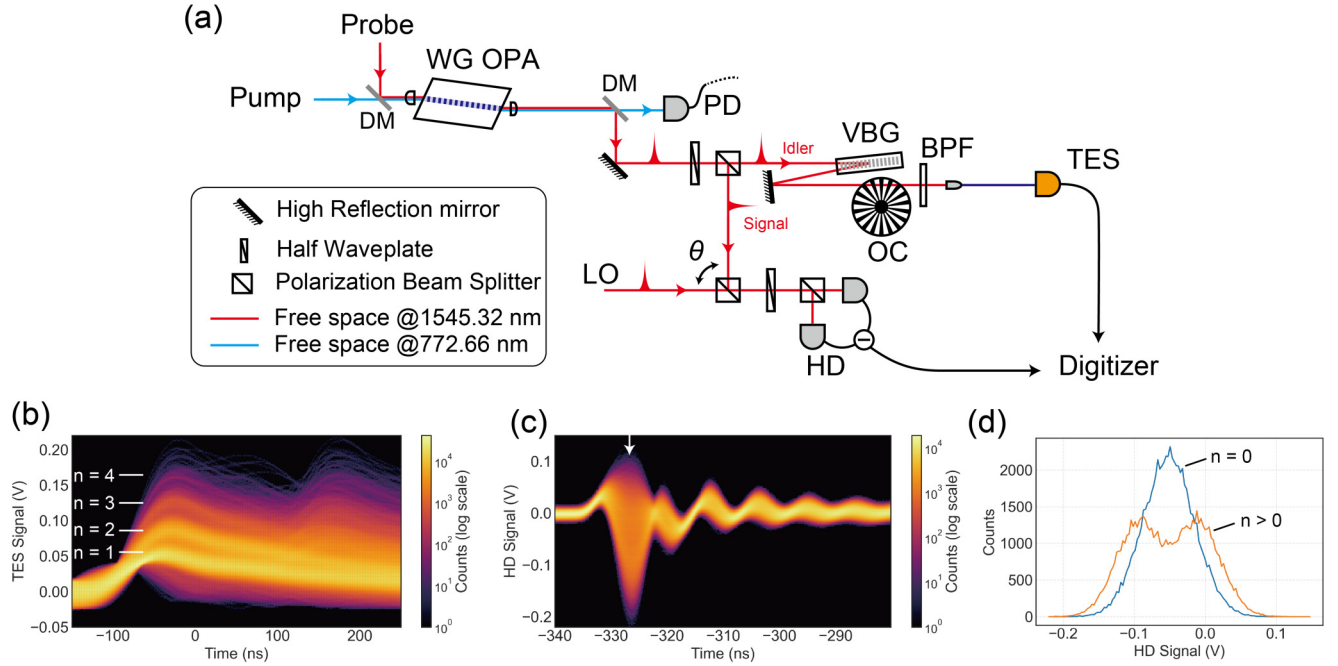


FIG. 2. (a) Experimental apparatus for the four-photon subtraction: DM, dichroic mirror; WG OPA, waveguide optical parametric amplifier; PD, photodetector; VBG, volume Bragg grating; BPF, optical bandpass filter; OC, optical chopper; TES, Ti-Au transition-edge sensor; LO, local oscillator; and HD, homodyne detector. (b) Heatmap of TES signals over approximately 70 000 frames, where signals corresponding to $n = 1, 2, 3, 4$ photons are observed. (c) Heatmap of HD signals at a fixed LO phase over approximately 70 000 frames when photon detection was performed using the TES ($n > 0$). The quadrature-phase amplitudes were calculated based on the voltage values at the timestamps indicated by the downward arrow in the figure. (d) Voltage histograms of the HD signal at the downward arrow (the blue curve is for $n = 0$ and the orange for $n > 0$).

The transmitted mode is referred to as the idler mode, and the reflected mode as the signal mode.

Next photon-number-resolving detection is performed on the idler mode by a PNRD. When n photons are detected, the signal mode collapses into the state $\hat{a}^n \hat{S}(r)|0\rangle$ ($\hat{a}$ is an annihilation operator), which serves as an approximation of a cat state. Notably, the parity (even or odd) of the generated cat state matches that of n , as the squeezed-vacuum state consists of even-photon-number superpositions. This approach is referred to as photon subtraction, since the conditional detection on the tapped mode implements the subtraction of photons from the squeezed-vacuum state. Figures 1(b)–1(d) show the relevant Wigner function defined as

$$W(x, p) = \frac{1}{\pi \hbar} \int_{-\infty}^{\infty} e^{2ip'x/\hbar} \langle p + p' | \rho | p - p' \rangle dp', \quad (1)$$

where ρ is the density matrix of the state. Figure 1(b) shows the Wigner function of a squeezed-vacuum state with $r = 0.576$, and the Wigner function exhibits no negative values. In Figs. 1(c) and 1(d), the Wigner functions of the simulated state generated by tapping the initial squeezed vacuum with a beam splitter ($R = 0.78$) and detecting three and four photons in the idler mode are presented. The Wigner function exhibits negative values that are not achievable for the input squeezed state, and the interference fringes in $W(x, 0)$ near the origin, together with the sign of W at the origin, reflect the quantum coherence and parity of the cat state. Achieving such negativity without loss correction is a key experimental goal, as it signifies a high degree of nonclassicality.

III. EXPERIMENTAL APPARATUS

The experimental apparatus is illustrated in Fig. 2(a). For more details (detailed experimental setup, phase-locking scheme, and loss budget), refer to Appendixes A–D and our previous paper [27]. A 10-mm-long waveguide-type periodically poled lithium niobate crystal (WG OPA) was pumped by a laser with a central wavelength of 772.66 nm, a pulse duration of 10 ps, and a repetition rate of 5 MHz, generating a pulsed squeezed-vacuum field at a central wavelength of 1545.32 nm.

Since this experiment is phase sensitive, a probe beam with a wavelength of 1545.32 nm was introduced coaxially through a dichroic mirror to measure the parametric gain in the waveguide via a photodiode, thereby referencing the phase of the squeezed-vacuum field to the probe beam. The squeezed vacuum was split into two optical paths by a BS composed of a half waveplate and a polarization beam splitter. When measuring the input squeezed-vacuum state, the reflectivity was set to $R = 1$, and for the photon subtraction experiment, the reflectivity was set to $R = 0.81$. The transmitted light is referred to as the idler, while the reflected light is referred to as the signal.

A volume Bragg grating (VBG) used in the idler path has a central wavelength of 1545.32 nm and a bandwidth of approximately 30 GHz. This narrowband filtering is used to select a portion of the broadband (terahertz-scale) squeezed vacuum, thereby defining the temporal mode of the heralded signal state to be on the order of 10 ps. While the VBG provides

TABLE I. Quantitative characterization of the generated states. For each quantum state, the tomography was performed using the number of data points listed in column 2. The Wigner negativity is the minimum value of the reconstructed Wigner function without loss correction; the bracketed value is that of the ideal (lossless) squeezed cat target at the same phase-space point. Column 4 gives the deepest (maximal) negativity attained by the ideal target. Here F_{sim} is the fidelity to the simulated state ($F = \text{Tr}[(\sqrt{\rho_{\text{expt}}}\rho_{\text{sim}}\sqrt{\rho_{\text{expt}}})^{1/2}]$, with ρ_{sim} and ρ_{expt} the simulated and experimental density matrices) and F_{cat} is the fidelity to an ideal squeezed cat state optimized over the cat amplitude $|\alpha|$ and squeezing (parity fixed by n ; the optimal $|\alpha|$ is listed). The $\sqrt{\lambda_{\text{max}}}$ is given by the largest eigenvalue λ_{max} of ρ_{expt} , and the purity is $\text{Tr}(\rho_{\text{expt}}^2)$. Uncertainties represent the standard deviation 1σ estimated from bootstrap samples.

State	Number of data points	Wigner negativity	Ideal W_{min}	F_{sim}	F_{cat}	$\sqrt{\lambda_{\text{max}}}$	$ \alpha $	Purity
Input squeezed	600k			0.9881(4)				0.856(2)
$n = 0$	600k			0.9968(4)				0.824(2)
$n = 1$	60k	$W_1(0, 0) = -0.017(4) [-0.318]$	-0.318	0.973(3)	0.713(4)	0.715(4)	≈ 0	0.449(3)
$n = 2$	60k	$W_2(-0.9, 0) = -0.008(2) [-0.062]$ $W_2(1.0, 0) = -0.004(3) [-0.042]$	-0.078	0.984(3)	0.803(4)	0.804(4)	0.97(1)	0.524(5)
$n = 3$	60k	$W_3(0, 0) = -0.019(5) [-0.318]$	-0.318	0.980(2)	0.711(4)	0.713(4)	1.18(3)	0.456(3)
$n = 4$	28k	$W_4(-0.8, 0) = -0.008(5) [-0.093]$ $W_4(0.5, 0) = -0.005(6) [-0.141]$	-0.164	0.951(3)	0.698(7)	0.712(6)	1.34(2)	0.428(5)

a narrow spectral bandwidth and a high diffraction efficiency (98%), its extinction ratio is limited to approximately 30 dB. To further suppress unwanted spectral components, we introduced an additional dielectric multilayer bandpass filter (BPF) with a bandwidth of approximately 200 GHz and a transmission of about 98%. Then the idler was coupled into a single-mode fiber and detected by a Ti-Au TES-type PNRD installed in an adiabatic demagnetization refrigerator, and the TES was cooled down to 280 mK.

The signal of the TES was read out using a superconducting quantum interference device and acquired by a high-speed digitizer. The overall idler detection efficiency, including fiber coupling, spectral filtering (VBG and BPF), the TES efficiency (89.2%), and the dominant joint-spectral-amplitude-mode-mismatch contribution, is approximately 40%. The 10.8% listed in Table II refers to the loss of the TES alone.

On the signal side, homodyne detection was performed by combining the signal light with a local oscillator (LO) beam and using a homemade balanced photodetector consisting of a pair of high-efficiency photodiodes and a low-noise transimpedance amplifier. The temporal waveform of the LO beam was optimized using a waveform shaper. The phase of the homodyne measurement can be determined by the

interference signal between the probe and the LO beam, and it was locked to any desired angle using a waveguide phase modulator. The results of the homodyne measurement were acquired by the same digitizer.

The digitizer measured the signal waveforms from both the TES and the homodyne detector (HD), and the signals were analyzed by the field-programable gate array (FPGA) embedded in the digitizer. Figure 2(b) shows a heatmap of TES signals over approximately 70 000 frames, where signals corresponding to $n = 1, 2, 3, 4$ photons are observed. Figure 2(c) shows a heatmap of HD signals over approximately 70 000 frames when photon detection was performed. The quadrature-phase amplitudes were calculated based on the voltage values at the timestamps indicated by the downward arrows. Examples of voltage histograms are shown in Fig. 2(d), where blue and orange lines represent the HD signals, when $n = 0$ and $n > 0$, respectively. The HD signal, when the signal light was blocked, was also recorded as the shot-noise level, and this information was used to normalize the voltage values of the homodyne measurement. It should be noted that the data shown in Figs. 2(b)–2(d) were not used for actual tomography in the next section. All the homodyne data used for the tomography can be found in Appendix F.

TABLE II. Losses in the experimental setup. Values marked with an asterisk are estimated, while the others are directly measured.

Category	Contribution	Loss (%)
Light source	Waveguide OPA	5*
Signal mode	Propagation loss	1
	Spatial mode mismatch	2
	Temporal mode mismatch	9*
	Photodiode inefficiency (HD)	2
	Electronic noise (HD)	2
Idler mode	Fiber coupling loss	10
	VBG	2
	BPF	2
	TES efficiency	10.8
	Mode mismatch (including joint-spectral-amplitude mismatch)	53*

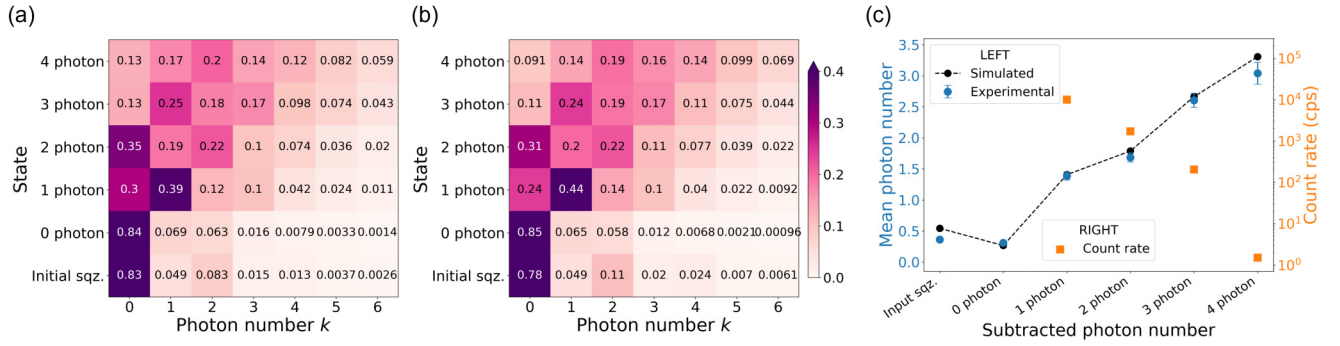


FIG. 3. (a) Experimentally reconstructed photon-number distributions $P(k)$ for the heralded states obtained after subtracting $n = 0-4$ photons from the squeezed vacuum. Each row corresponds to a state. The horizontal axis shows the photon number k in the heralded output state, and each cell gives the probability $P(k)$ for that state. (b) Simulated photon-number distributions $P(k)$. (c) Mean photon numbers of the heralded states (left axis) and the experimentally measured count rates (right axis). Blue circles with error bars show the experimental results, black circles the simulated values, and orange rectangles the measured count rates.

In this experiment, the probe beam used for phase referencing was chopped at a frequency of 500 Hz, and the phase locking of each part of the system was achieved using a sample-and-hold method. The phase was locked when the probe beam was ON, and the state was held when the probe beam was OFF. Photon detection with the TES and homodyne measurements were performed when the probe was OFF to prevent the classical probe light from interfering with the homodyne measurement. Additionally, an optical chopper was placed before the optical fiber to prevent the probe light from saturating the TES output.

IV. RESULTS

A. Homodyne measurement and quantum state tomography

We performed homodyne measurements at six discrete LO phases ($\theta = -45^\circ, -22.5^\circ, 0^\circ, 22.5^\circ, 45^\circ$, and 90° , where $\theta = 0^\circ$ and 90° correspond x and p , respectively). While continuous or more finely sampled phase scanning is in principle possible, in the present experiment, the number of conditioned events per phase becomes limited for higher-order photon-subtraction events due to their low heralding probability. We therefore used a discrete set of phases to ensure sufficient statistics at each phase for stable maximum-likelihood reconstruction. For each quantum state, the tomography was performed using the number of data points listed in Table I.

The phase spacing was chosen to be nonuniform. In particular, the sampling was made denser along the squeezed quadrature, where the Wigner-function negativity and interference fringes are most pronounced, in order to better resolve the non-Gaussian features of the reconstructed states.

The measured quadrature data were used to reconstruct the density matrix in the photon-number basis by the maximum-likelihood method [32], from which we obtained photon-number distributions, Wigner functions, and other derived quantities. The reconstructed states obtained from this discrete-phase dataset are in good agreement with the calibrated simulation model (see fidelities in Table I), indicating that the chosen phase sampling provides sufficient information for reliable state reconstruction under the present experimental conditions. The simulations in Figs. 3 and 4 are performed using the loss values listed in Table II.

B. Photon-number distributions and count rates

Figure 3(a) shows heatmaps derived from the photon-number distributions $P(k)$ for each state. In the absence of loss, the input squeezed-vacuum state consists solely of even-photon-number components. However, due to experimental losses, odd-photon-number components are also present. Specifically, the initial squeezed state is measured with $R = 1$, whereas the zero-photon state is measured with the tapping beam splitter ($R = 0.81$). This difference is clearly evidenced by the density-matrix elements: The two-photon component $P(2)$ decreases from 0.083 to 0.063, which quantitatively confirms the loss induced by the beam splitter. This trend is also reproduced in the numerical simulation shown in Fig. 3(b). Next, when one photon is subtracted, the photon-number distribution shifts. This indicates that applying the annihilation operator shifts the photon-number distribution by one $\hat{a}|n\rangle = \sqrt{n}|n-1\rangle$ ($n > 0$). Moreover, the original zero-photon component is truncated $\hat{a}|0\rangle = 0$, causing both the maximum and mean photon number of the cat state to increase as more photons are subtracted. This behavior is the same for cases where two or more photons are subtracted. Figure 3(b) shows heatmaps obtained by the simulation.

Figure 3(c) plots the simulated (black circles with dashed lines) and experimental (blue circles with error bars) mean photon numbers (on the left axis). The error bars for the experimental data represent the standard deviation 1σ calculated using the bootstrap method. As can be seen from the results, the experimental data agree very well with the simulation.

Figure 3(c) also plots generation rates on the right axis with orange rectangles. As evident from this figure, the generation rate decreases exponentially with the increase in photon number. In previous studies [26,27], the event rate for three-photon subtraction ranged from a few events per minute to a few events per second, rendering four-photon subtraction experiments impractical. In this study, leveraging a broadband squeezed light source, a PNRD, and other advanced experimental techniques, we achieved an event rate of 200 cps for three-photon and 1.5 cps even for four-photon subtraction.

The n -photon heralding rate scales as $R_n \propto f_{\text{rep}}(T \sinh^2 r \eta_i)^n$. Within the present scheme the rate can be improved along three independent directions. (a) The

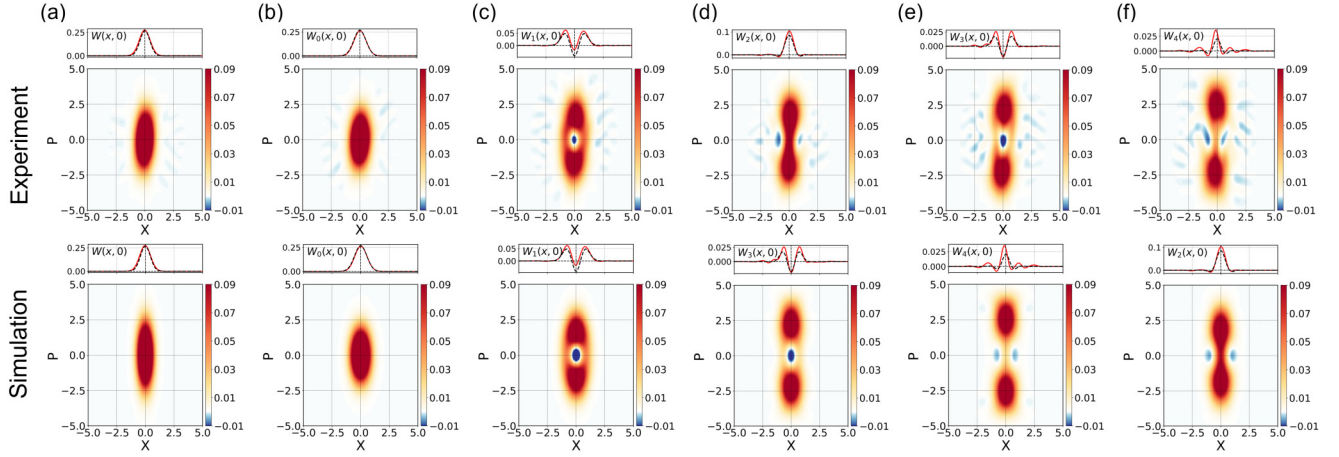


FIG. 4. Wigner-function representations of the (a) input squeezed-vacuum state and (b)–(f) states corresponding to zero- to four-photon subtraction. The top row shows the experimentally reconstructed Wigner functions (without loss correction), while the bottom row shows the corresponding simulated results. The top panels display the cross section $W(x, 0)$, where the red solid line denotes the experimental data and the black dashed line denotes the simulation.

repetition rate is currently 5 MHz, set by the TES timing of approximately 50 ns; since the source laser operates at 100 MHz, a faster detector would provide up to a 20-fold increase. (b) The dominant idler loss is the joint-spectral-amplitude-mode-mismatch contribution (53%, Table II); reducing it raises the idler efficiency η_i , which enters the rate as η_i^n . (c) The squeezing level can be increased, as the similar waveguide OPA has produced more than 10 dB [33]. For reference, optical quantum computation requires a logical clock at the megahertz level; because GKP synthesis from cat states is probabilistic and consumes several cats per qubit, the practically required cat-generation rate is well above the megahertz range, to be compared with the 200 cps achieved here for three-photon subtraction. We also note that complementary measurement schemes that are not constrained by a weak tap, such as generalized photon subtraction [25], provide a further route.

C. Wigner functions

The experimentally reconstructed and simulated Wigner functions are shown in Fig. 4. The top panels display the cross section $W(x, 0)$, where the red solid line denotes the experimental result and the black dashed line denotes the simulation. Figure 4(a) shows the input squeezed-vacuum state measured with the beam-splitter reflectivity set to $R = 1$, for which the Wigner function exhibits a Gaussian profile. In contrast, Figs. 4(b)–4(f) correspond to the cases where the beam-splitter reflectivity is set to $R = 0.81$ and zero to four photons are subtracted, respectively. The sign of the Wigner function at the origin indicates the parity of the state: positive for even n and negative for odd n . Remarkably, the negative values in the Wigner function persist even after four-photon subtractions, demonstrating the action of the annihilation operator four times without the need for loss correction.

For $n \geq 1$, the minimum values of the Wigner function were evaluated with statistical uncertainties estimated using a bootstrap resampling method. The extracted minimum values at representative phase-space points and the purities are

shown in Table I. In all photon-subtracted states, at least one region with negative values is observed. Ideally, for two- and four-photon subtraction, the Wigner function is expected to exhibit pairs of negative regions located symmetrically in phase space. In the experiment, this symmetry is partially broken due to experimental imperfections. In particular, for the four-photon-subtracted state ($n = 4$), the limited number of data points leads to larger tomographic uncertainties, and only a single statistically significant negative region is observed.

In addition to the persistence of Wigner-function negativity, we observe that the visibility of the interference fringes remains high even as the separation between the two peaks increases with photon subtraction. In particular, the cross sections $W(x, 0)$ show pronounced oscillatory features for higher photon numbers, indicating that the coherent interference between the two components of the cat state is preserved as the state amplitude grows. Although additional negative regions are not always resolved with statistical significance for higher-order subtraction, especially for $n = 4$, the maintained fringe visibility demonstrates that the cat state continues to grow without a substantial loss of coherence in the Wigner representation.

In general, higher-order photon-subtracted states are more sensitive to loss, and therefore metrics such as Wigner-function negativity and purity are typically expected to be optimal for single-photon subtraction. However, in the present results, the case of $n = 1$ does not exhibit the best performance. This is likely due to the photon-number-resolving detection with the TES, where electrical noise can trigger false detection events. As a result, the heralded $n = 1$ state is partially contaminated by contributions from the $n = 0$ state, effectively forming a mixed state.

To clarify what type of non-Gaussian states are produced, we additionally compare each reconstructed state with an ideal squeezed cat state of optimized amplitude and squeezing. As listed in Table I, the fidelities F_{cat} reach 0.71–0.80 and saturate the upper bound $\sqrt{\lambda_{\text{max}}}$ imposed by the state purity,

confirming that the generated states are squeezed cat states limited only by loss-induced mixedness.

For completeness, we also analyze the reconstructed states in the quadrature-basis density-matrix representation, which provides a complementary view of the coherence structure (see Appendix F).

D. Cat states as resources for GKP preparation

To assess the generated states as resources for GKP-state preparation, we consider one step of cat breeding [13], in which two of the generated cats are interfered on a balanced beam splitter and one output is homodyne projected, and quantify the bred state by the GKP effective squeezing s_{eff} [16,34] [$\Delta_q = \sqrt{-(2/|\alpha_q|^2) \ln |\langle \hat{D}(\alpha_q) \rangle|}$, $s_{\text{eff}} = -10 \log_{10} \Delta_q^2$ (s_{eff} in decibels) with $|\alpha_q| = \sqrt{2\pi}$ for the square-lattice code]. Using the even-parity cats generated here, one breeding step yields $s_{\text{eff}} \approx 4.7$ and 2.7 dB for $n = 2$ and 4 at the present loss (6.9 and 4.2 dB in the lossless case). The $n = 2$ amplitude is close to the single-step optimum ($|\alpha| \approx 1$), while the larger-amplitude $n = 4$ cat is better suited to larger lattices or to multistep breeding. These values confirm that the generated cats already act as GKP resources; the effective squeezing is presently limited by loss, and reducing the optical loss together with employing multistep breeding is the natural route toward the fault-tolerance regime, which is left for future work.

V. CONCLUSION

In this work we successfully demonstrated photon subtraction of up to four photons from a squeezed-vacuum state using a broadband OPA and a high-efficiency TES. Despite the system's sensitivity to residual optical losses and phase instability, the experimental results showed excellent agreement with simulations, confirming the nonclassical nature of the generated states in the picosecond optical wave packet through Wigner-function negativity, all obtained without any data correction.

This achievement extends the accessible regime of multi-photon subtraction in traveling optical modes and contributes to the generation of higher-amplitude non-Gaussian states, which are valuable resources for quantum information processing. The scalability of the experimental setup enables future extensions to photon addition [35–37], generalized photon subtraction [25], quantum state synthesizers [38], and GKP-state preparation schemes [16]. These techniques offer pathways toward the high-quality and high-rate generation of Schrödinger-cat states in the ultrashort wave packet and GKP-like resource states relevant to CV quantum technologies.

By advancing the generation of complex optical quantum states, this work supports future developments in quantum sensing, quantum communication, and high-speed optical quantum-state engineering.

ACKNOWLEDGMENTS

This work was partly supported by Japan Science and Technology Agency (Moonshot R&D, Grant No. JPMJMS2064, and PRESTO, Grant No. JPMJPR2254), the

UTokyo Foundation, and donations from Nichia Corporation. W.A. acknowledges funding from Japan Society for the Promotion of Science KAKENHI (Grant No. 23K13040). M.E., K. Takase, and W.A. acknowledge supports from the Research Foundation for Opto-Science and Technology. P.M. acknowledges Project No. 25-17472S of the Czech Science Foundation, the European Union's HORIZON Research and Innovation Actions under Grant Agreement No. 101080173 (CLUSTEC), and Project No. CZ.02.01.01/00/22_008/0004649 (QUEENTEC) of EU and the Czech Ministry of Education, Youth and Sport. R.F. acknowledges Project No. LUC25006 and the Quantera project CLUSSTAR (No. 8C24003) of MEYS, Czech Republic. Project CLUSSTAR has received funding from the European Union's Horizon 2020 Research and Innovation Programme under Grant Agreements No. 731473 and No. 101017733 (QuantERA). M.E. conceived the project and led the experiment. M.E. built the optical and electrical setup with support from T.N., T.S., K. Takahashi, and R.N. M.E. coded the control program and the FPGA program for data acquisition with support from T.N. M.E., P.M., and R.F. performed simulations and theory evaluation and analyzed the data. T.K., A.I., and T.U. provided the OPA used in the experiment. S.T. and D.F. provided the TES used in the experiment. M.E. wrote the manuscript with P.M., R.F., K. Takase, W.A., and A.F., with input from all the co-authors.

The authors declare no competing interests.

DATA AVAILABILITY

The reconstructed density matrices, Wigner functions obtained from the experiment, and the simulation codes used in this study are openly available [39]. The raw homodyne measurement data (quadrature-phase amplitudes) are available from the corresponding authors upon reasonable request.

APPENDIX A: DETAILS OF THE WG OPA

We use a periodically poled lithium niobate (PPLN) waveguide as the OPA for squeezed light generation. The core of the waveguide consists of ZnO-doped LN and is directly bonded to a lithium tantalate substrate. The periodic poling structure for quasi-phase-matching (QPM) is fabricated using an electrical poling method [40]. The poling cycle is approximately 18 μm , and the structure is designed as a type-0 PPLN. The waveguide is 10 mm long, and its core has a width of about 9 μm and a thickness of about 8 μm . The waveguide structure is fabricated using a mechanical dicing saw. More details on the fabrication can be found in Ref. [41]. Both the input and output end faces of the waveguide are coated with antireflective coatings for both 1545.32-nm and 772.66-nm light. The temperature of the waveguide is controlled to maintain the QPM condition.

APPENDIX B: DETAILS OF THE TI-AU TES

The TES device was fabricated at the superconducting quantum circuit fabrication facility at AIST. The superconducting TES film comprises bilayers of titanium and gold, deposited by dc magnetron sputtering to thicknesses of 20

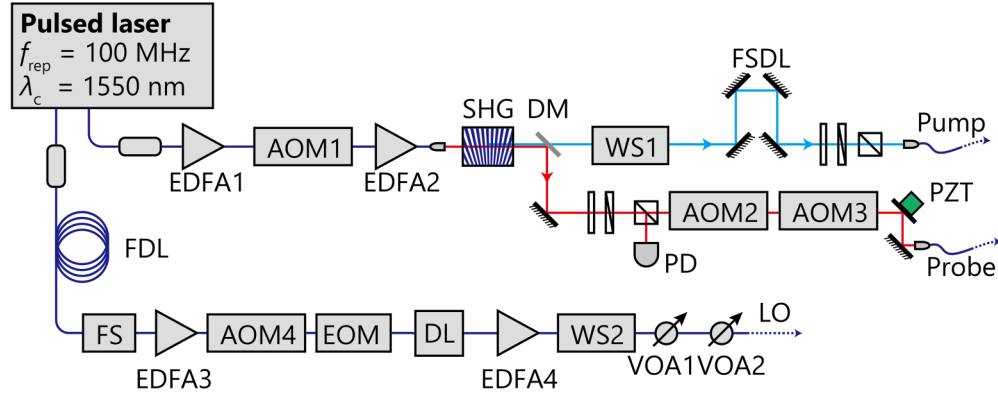


FIG. 5. Setup for the pump, probe, and LO light preparation: SHG, second-harmonic generation; FSDL, free-space delay line; PZT, piezo-electric transducer; FDL, fiber delay line; FS, fiber stretcher; EOM, electro-optic modulator; DL, delay line; and VOA, variable optical attenuator.

and 10 nm, respectively. To enhance detection efficiency, an optical cavity structure was adopted, in which the TES films are embedded between a high-reflection mirror and antireflection coatings. The TES is patterned into a square with dimensions of $5 \times 5 \mu\text{m}^2$. Superconducting niobium leads were fabricated on two opposite sides of the TES, and the resulting resistance changes are read via these leads using a superconducting quantum interference device-based current amplifier.

To deliver photons, a high-numerical-aperture (high-NA) optical fiber (UHNA7, Nufern) with a mode field diameter (MFD) of $3.2 \mu\text{m}$ was coupled to the TES device. The fiber-coupled TES module was installed in an adiabatic demagnetization refrigerator and cooled to below 280 mK. Inside the refrigerator, the high-NA optical fiber was spliced to an SMF28 fiber (Corning) to convert the MFD to $10 \mu\text{m}$, with an MFD conversion loss of less than 0.2 dB.

The critical temperature and normal resistance of the TES are 308.6 mK and 2.653Ω , respectively. The TES exhibits an energy resolution of $\Delta E = 0.176 \text{ eV}$, corresponding to an $E/\Delta E$ ratio of 4.54 for a single-photon energy of 0.8 eV at a wavelength of $1.5 \mu\text{m}$, thereby demonstrating sufficient capability to resolve individual photon states. The system detection efficiency of the TES was carefully calibrated using an evaluation system traceable to the National Metrology Institute of Japan at AIST, yielding an efficiency of 89.2% at a wavelength of $1.5 \mu\text{m}$. The decay time constant of the observed signals was 107 ns, enabling a repetition frequency of up to 5 MHz.

APPENDIX C: PUMP, PROBE, AND LO BEAM PREPARATION

As shown in Fig. 5, a femtosecond mode-locked fiber laser (M-Comb-ULN, Menlo Systems) with a center wavelength of 1550 nm, a spectral bandwidth of over 40 nm, and a repetition rate of 100 MHz was employed as the light source in this experiment. Since the timing resolution of the TES used in our experiment was limited to approximately 50 ns, we reduced the pulse repetition rate to 5 MHz via pulse picking with acousto-optic modulators (AOMs) and then applied bandpass filters to restrict the optical spectrum, thereby preparing a non-Gaussian state with a temporal width of about 10 ps.

Likewise, the probe light for monitoring the squeezed light phase during the parametric process and the LO light for homodyne detection were also pulse picked and spectrally filtered. Details are provided below.

1. Pump and probe lights

One of the multiple output ports of the mode-locked laser was used to generate both the pump and probe lights. All optical fibers used in this experiment were polarization-maintaining (PM) fibers. First, the spectrum of the output light was limited by an optical bandpass filter to a center wavelength of 1545.32 nm and a spectral width of 1.0 nm, and it was then amplified by a homemade erbium-doped fiber amplifier (EDFA1) using an Er80-8/125-PM fiber (nLight) as the gain medium. Next a fiber-coupled AOM (AOM1) [SGTF400-1550-1P(H), Chongqing Smart Science & Technology Development] driven at 400 MHz was used for pulse picking, thus lowering the repetition rate to 5 MHz. The driving and gating signals for this AOM were obtained by frequency dividing the repetition rate of the mode-locked laser [42]. The pulses were subsequently amplified again by EDFA2 (HPP-PMFA-22, PriTel) and then emitted into free space. The beam was then focused into a fan-out periodically poled MgO-doped near-stoichiometric LiTaO₃ crystal (fan-out PPMgSLT, OXIDE), generating second-harmonic light at a center wavelength of 772.66 nm and a spectral width of 0.5 nm. The second-harmonic and fundamental lights were separated by a dichroic mirror and used as the pump and probe beams, respectively. To suppress unwanted spectral components, we applied a homemade wave shaper (WS1). To match their timing, a free-space delay line was used for temporal adjustment, followed by a half waveplate (HWP) and a polarization beam splitter (PBS) to adjust the pump power. Subsequently, the pump light was coupled into a PM single-mode fiber. Part of the fundamental light was split off by a HWP and a PBS and directed to a photodiode for trigger signals, while the remaining fundamental light was used as the probe beam. The probe beam was modulated by two free-space AOMs (AOM2 and AOM3; AOM3080-1912, Gooch & Housego) at 500 Hz for the sample-and-hold phase-locking method. In addition, for heterodyne beat locking, AOM2 and

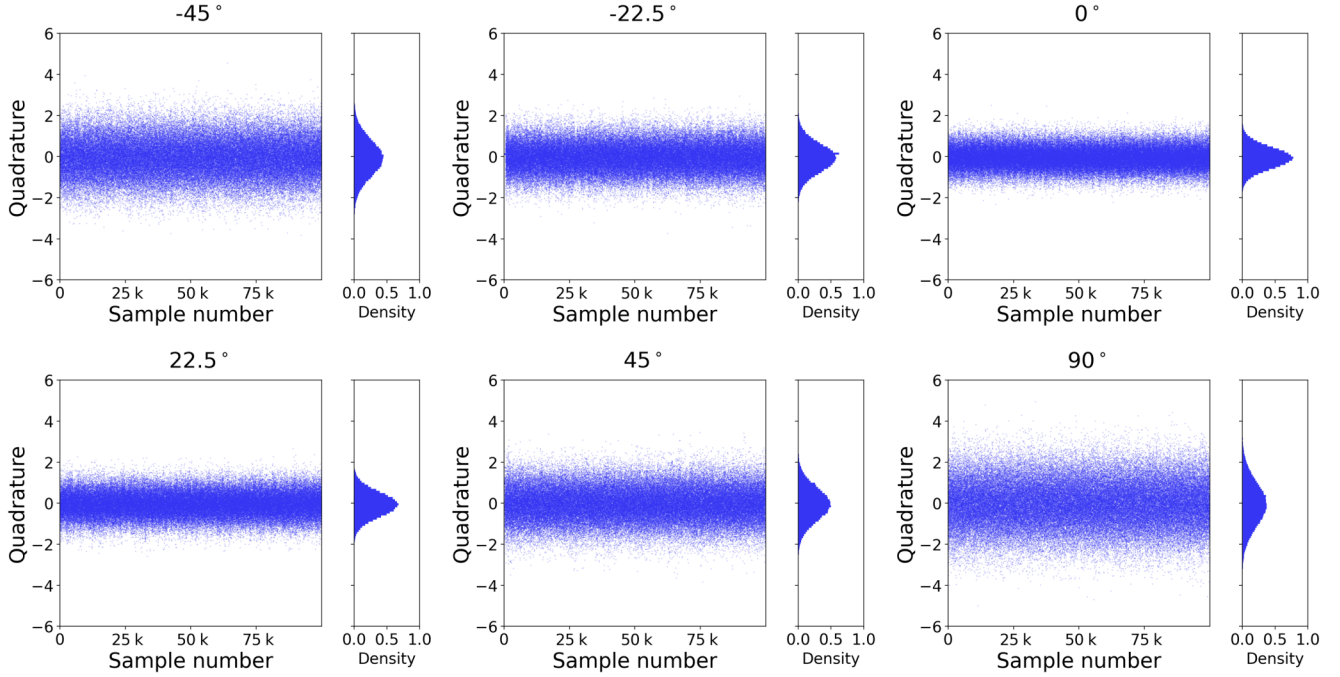


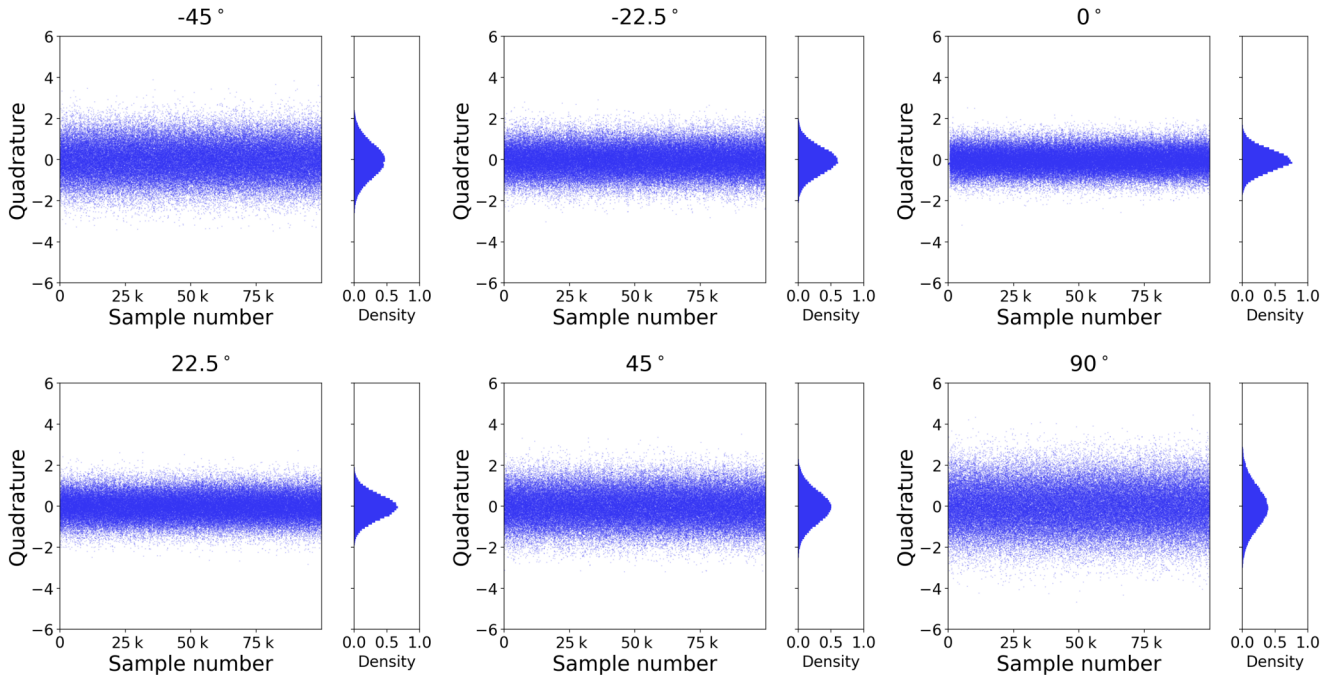
FIG. 6. Scatter plots and histograms of the measured quadrature data (input squeezed state).

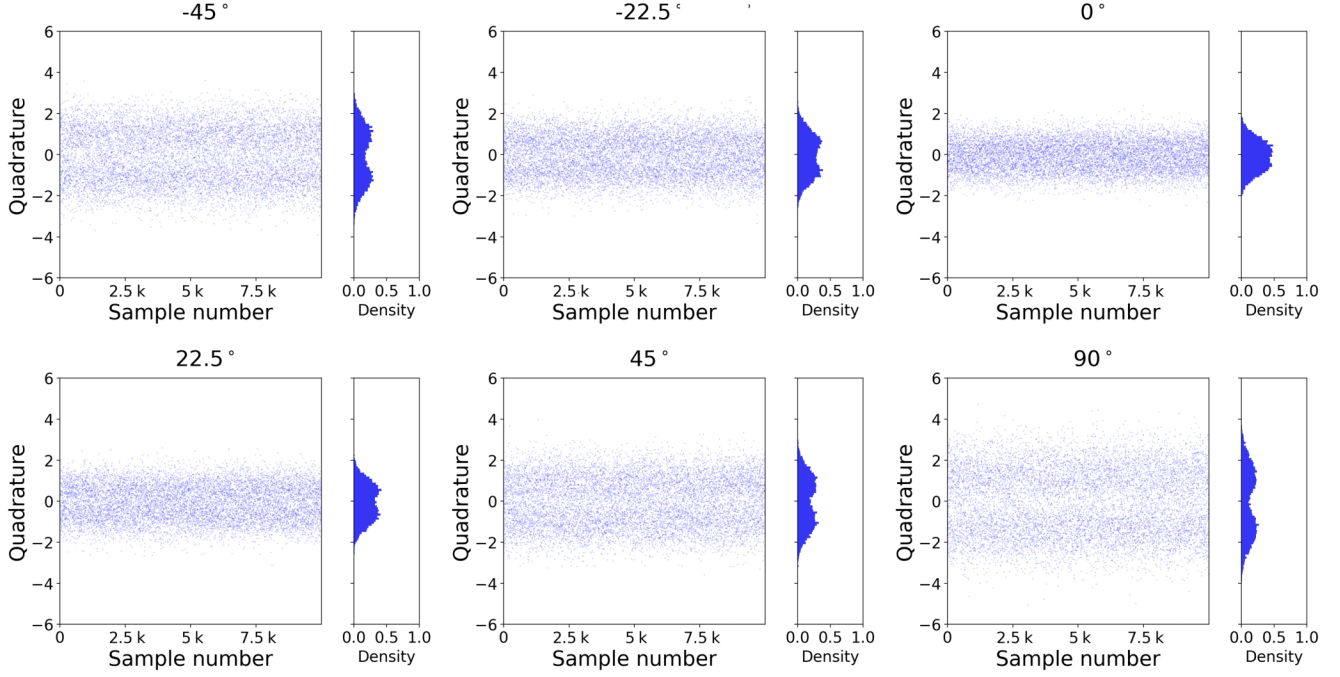
AOM3 were driven at -80.0 and 80.1 MHz, respectively, creating a difference frequency of 100 kHz. A mirror with a piezoelectric transducer was placed in the beam path for phase locking, and the beam was then coupled into a PM fiber.

2. LO light

The LO beam was obtained from a different output port of the mode-locked laser. In most pulsed-light experiments, the

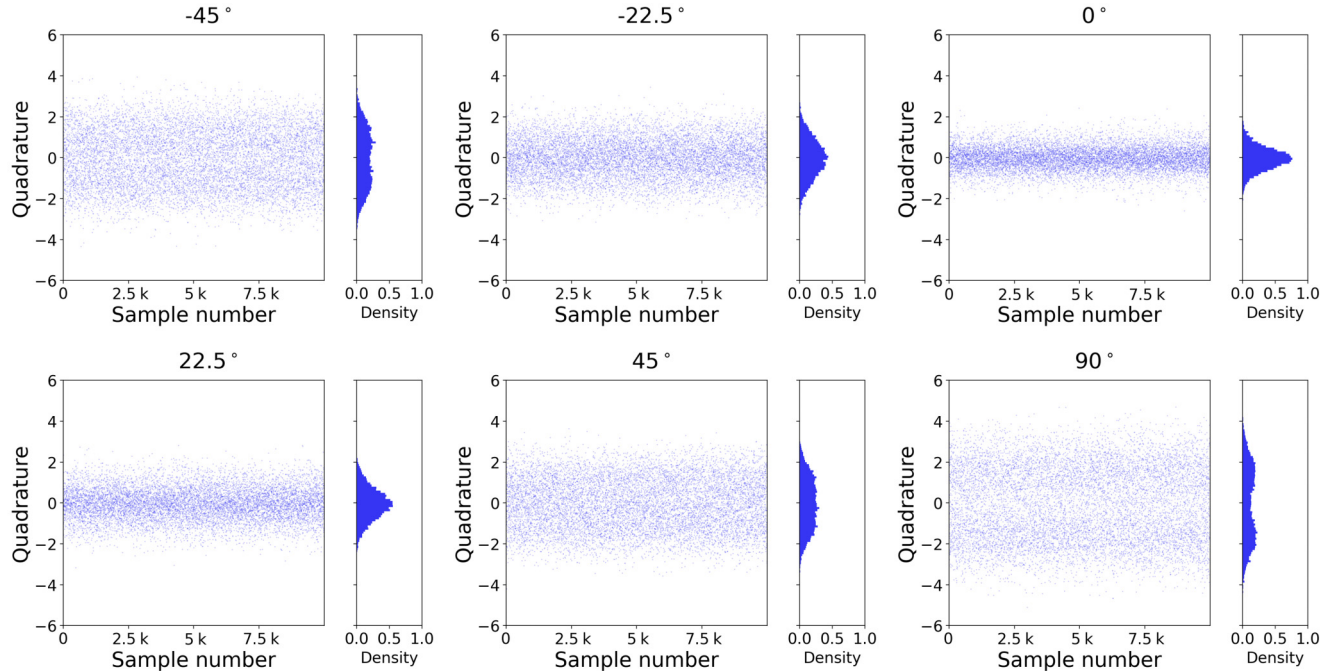
LO and pump beams are generated from the same port [26,43–45]; however, in that case the squeezed pulses after the OPA typically exhibit a shorter duration and broader bandwidth than the LO [26,46] and experience dispersive effects in the second-harmonic generation and OPA processes [47], making it difficult to realize high-level squeezing and large Wigner negativity without pulse shaping [26,46]. Here we used light taken from a separate output port and employed a commercial waveform shaper (WS2; WaveShaper 1000A, II-VI) to achieve more flexible waveform control of the LO beam.

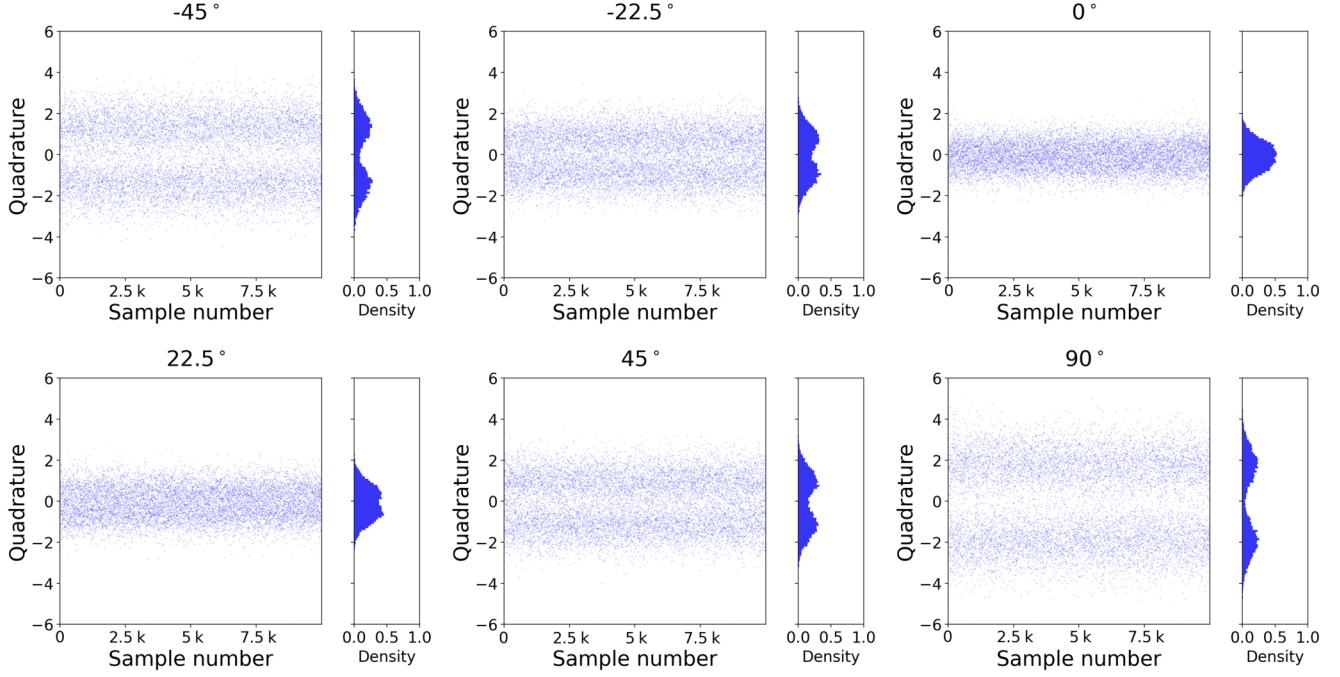
FIG. 7. Scatter plots and histograms of the measured quadrature data ($n = 0$).

FIG. 8. Scatter plots and histograms of the measured quadrature data ($n = 1$).

Specifically, the spectrum of the output beam was restricted by an optical bandpass filter (center wavelength of 1545.32 nm and bandwidth of 12 nm), and then the pulse was prechirped (stretched) using a 32-m fiber delay line to mitigate distortion caused by subsequent fiber amplification. A fiber stretcher using a cylindrical piezoelectric transducer (PT140.70, PI) was also inserted for long-term phase stabilization. The beam was then amplified by EDFA3 (EDFA100P, Thorlabs), after which it was pulse picked by a fiber-coupled AOM [AOM4; SGTF400-1550-1P(H), Chongqing Smart Science

& Technology Development] to reduce the repetition rate to 5 MHz. Subsequently, it was sent through a waveguide phase modulator (EOM; MPX-LN-0.1-00-P-P-FA-FA, iXblue) for phase stabilization in the feedback control loop. An optical delay line (DL; ODL-600-11-1550-8/125-P-60-3A3A-1-1-600, OZ optics) was used to fine-tune the optical path length. The beam was amplified again by EDFA4 (EDFA100P, Thorlabs), and then the spectral amplitude and phase were controlled by the WS2. Finally, the beam passed through variable optical attenuators (VOA1, MMVOA-1-1550-P-8/125-3A3A-1-0.5,

FIG. 9. Scatter plots and histograms of the measured quadrature data ($n = 2$).

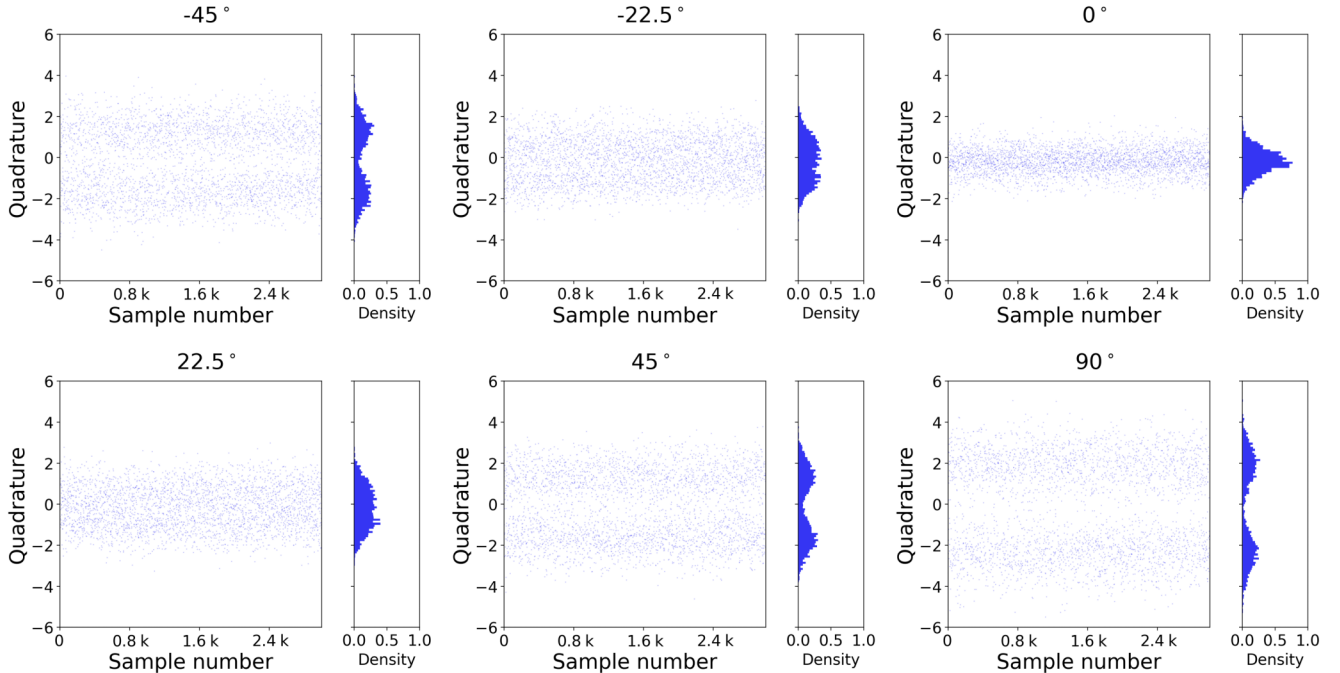
FIG. 10. Scatter plots and histograms of the measured quadrature data ($n = 3$).

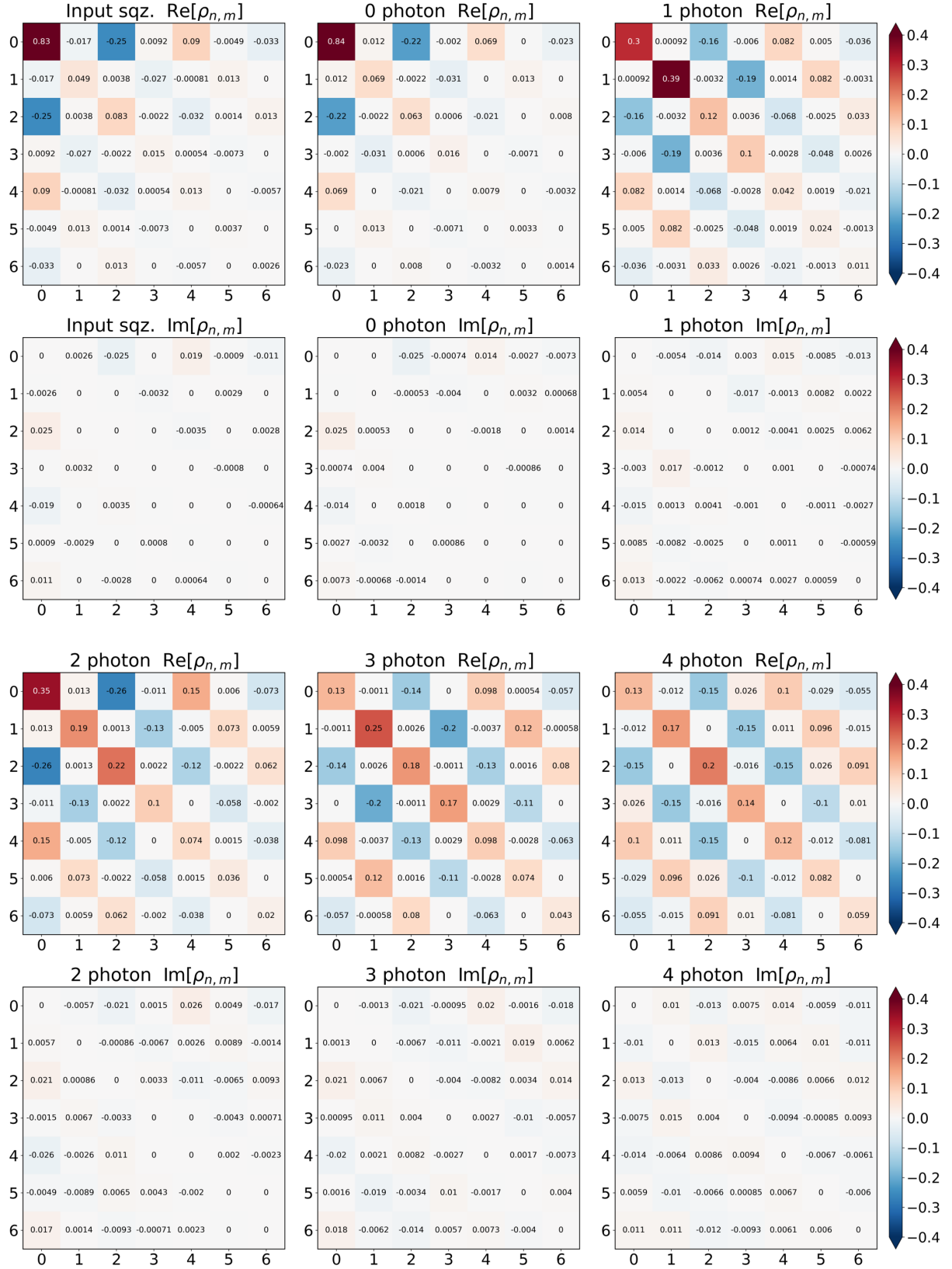
OZ optics; VOA2, NVOA-325125333-D0601, Agiltron) to stabilize and modulate the LO power and was emitted as free-space light through a collimator for homodyne detection. In this experiment, the WS2 was programed to produce a Gaussian waveform at a center wavelength of 1545.32 nm with a spectral width of 0.8 nm, to which a second-order dispersion of about -0.7 ps/nm was added to compensate for dispersion introduced by the optical elements. Autocorrelation measurements revealed an 8-ps pulse duration, corresponding to the transform-limited pulse width.

APPENDIX D: LOSS BUDGET

Here we summarize the losses in the experimental setup in Table II, including those discussed in the main text. The values listed here are used in the numerical simulations.

Among these contributions, the loss associated with the joint-spectral-amplitude mismatch on the idler side is particularly significant. This originates from the imperfect mode overlap between the broadband signal-idler correlations and the spectral-temporal mode selected by the filtering process, which effectively reduces the heralding efficiency.

FIG. 11. Scatter plots and histograms of the measured quadrature data ($n = 4$).

FIG. 12. Experimentally obtained density matrices in the photon-number basis [$\text{Re}[\rho_{n,m}]$ and $\text{Im}[\rho_{n,m}]$].

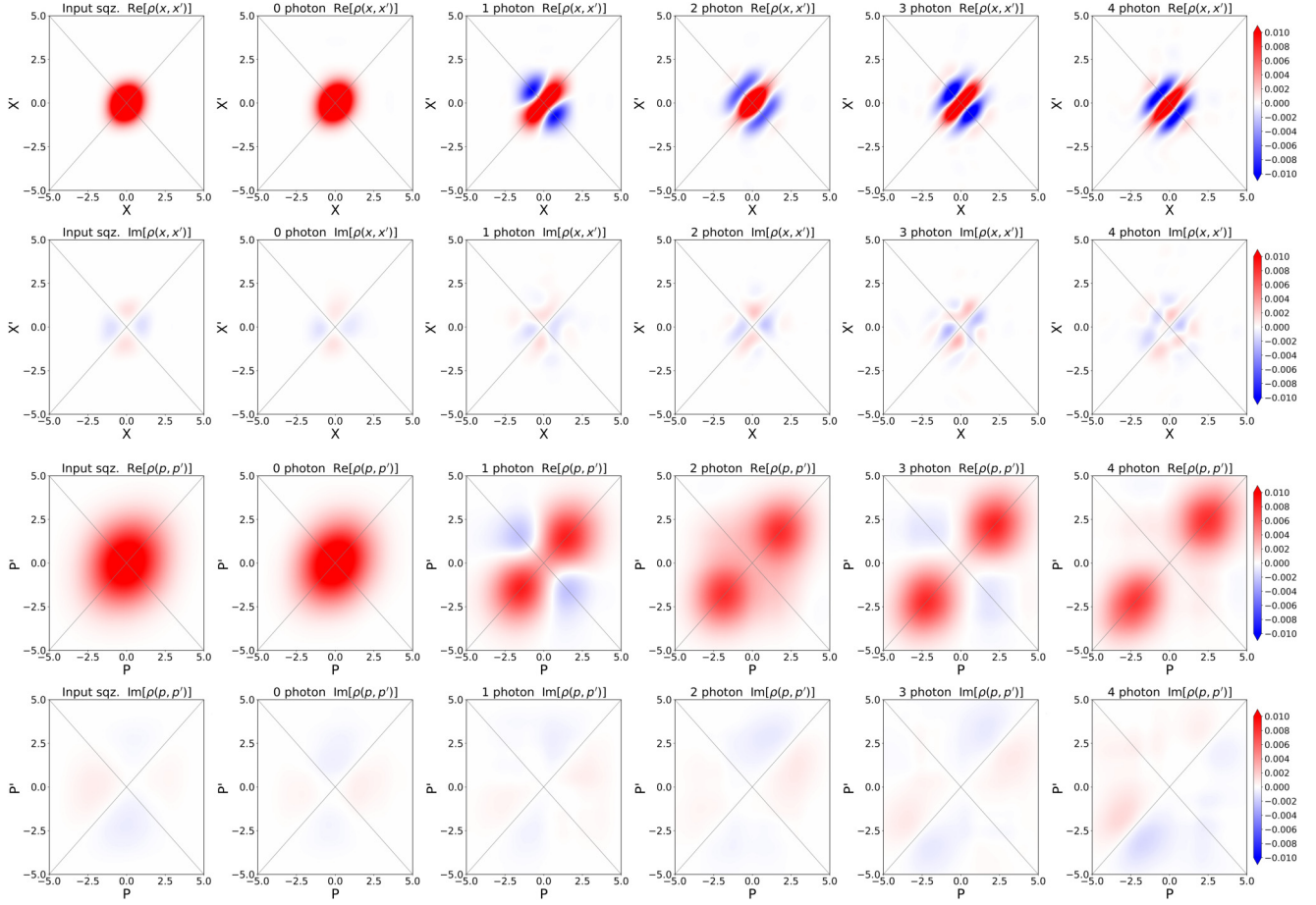


FIG. 13. Experimentally reconstructed density matrices in the position and momentum bases ($\text{Re}[\rho(x, x')]$, $\text{Im}[\rho(x, x')]$, $\text{Re}[\rho(p, p')]$, and $\text{Im}[\rho(p, p')]$). The columns correspond, from left to right, to the input squeezed-vacuum and the zero- to four-photon-subtracted states.

APPENDIX E: SIMULATIONS

The simulations were performed using *strawberryfields* [48,49]. The loss parameters were taken from the loss budget summarized in Table II. The input squeezed-vacuum state was assumed to have a squeezing level of 6.5 dB (before loss). Throughout this paper, we use units with $\hbar = 1$.

APPENDIX F: SUPPORTED DATA

1. Measured quadrature data

We summarized the measured quadrature-phase amplitudes obtained by homodyne detection into scatter plots and histograms, as shown in Figs. 6–11. Based on these results, we estimated the density matrix of the quantum state via a maximum-likelihood method without applying any loss correction.

2. Density-matrix representation

a. Density matrix in the photon-number bases

The density matrices reconstructed by homodyne tomography are presented in the photon-number basis in Fig. 12.

b. Density matrix in the quadrature bases

In the main text, we primarily discussed the reconstructed states using the Wigner function $W(x, p)$, which provides an intuitive visualization of the phase-space structure. Here we additionally present the density matrix in the quadrature basis $\rho(p, p')$, which offers a complementary perspective on the same states.

The density matrix in the momentum representation is obtained from the photon-number basis $\rho_{n,m}$ as

$$\rho(p, p') = \sum_{n,m} \rho_{n,m} \langle p|n\rangle \langle m|p'\rangle. \quad (\text{F1})$$

Figure 13 shows the real and imaginary parts of the density matrices in both position (x) and momentum (p) representations. While these representations contain the same information as the Wigner function, they highlight different aspects of the state, in particular the coherence between different quadrature components.

For the input squeezed-vacuum and the zero-photon subtraction case [Figs. 13(a) and 13(b)], the density matrices exhibit a single peak along the diagonal. The off-diagonal components are relatively narrow, reflecting a limited coherence length in the presence of loss.

In contrast, for one- to four-photon subtraction [Figs. 13(c)–13(f)], the diagonal components develop a

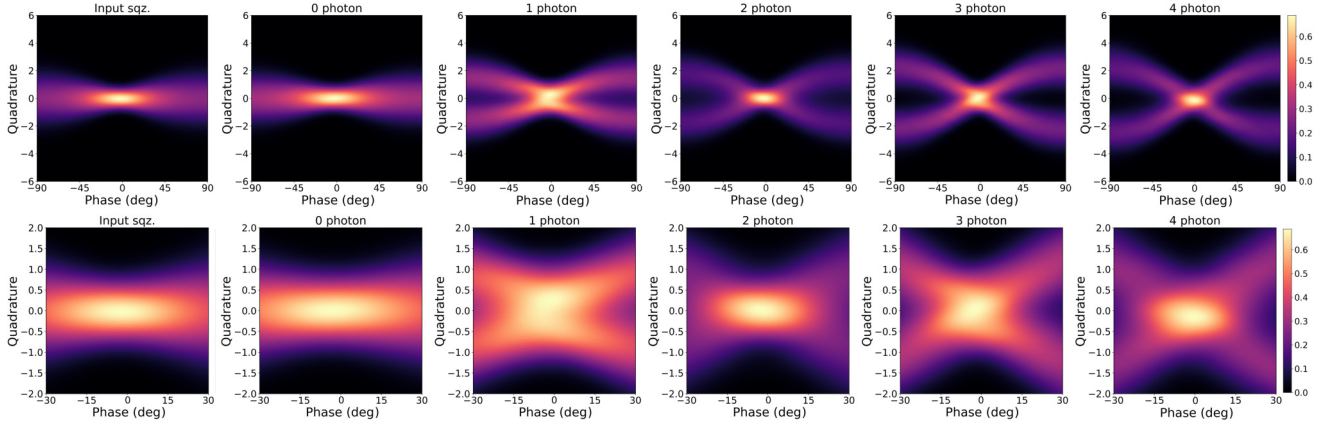


FIG. 14. Calculated marginal distributions from the experimental data. The columns correspond, from left to right, to the input squeezed vacuum and the zero- to four-photon-subtracted states. The top row displays data ranging from -90° to 90° , while the bottom row presents a magnified view of the range from -30° to 30° .

two-peak structure, corresponding to the two coherent-state components of a Schrödinger-cat state. Pronounced off-diagonal components are observed between these peaks, indicating coherent superpositions rather than statistical mixtures. The sign structure of these off-diagonal elements alternates depending on the number of subtracted photons, consistent with the expected parity of the generated states.

Overall, the density-matrix representation provides a complementary visualization of the coherence structure.

3. Marginal distributions

Figure 14 shows the marginal distributions calculated from the experimental data.

- [1] S. Lloyd and S. L. Braunstein, Quantum computation over continuous variables, *Phys. Rev. Lett.* **82**, 1784 (1999).
- [2] S. Takeda and A. Furusawa, Toward large-scale fault-tolerant universal photonic quantum computing, *APL Photonics* **4**, 060902 (2019).
- [3] W. Asavanant, Y. Shiozawa, S. Yokoyama, B. Charoensombutamon, H. Emura, R. N. Alexander, S. Takeda, J. I. Yoshikawa, N. C. Menicucci, H. Yonezawa, and A. Furusawa, Generation of time-domain-multiplexed two-dimensional cluster state, *Science* **366**, 373 (2019).
- [4] M. V. Larsen, X. Guo, C. R. Breum, J. S. Neergaard-Nielsen, and U. L. Andersen, Deterministic generation of a two-dimensional cluster state, *Science* **366**, 369 (2019).
- [5] H. Aghaee Rad, T. Ainsworth, R. N. Alexander, B. Altieri, M. F. Askarani, R. Baby, L. Bianchi, B. Q. Baragiola, J. E. Bourassa, R. S. Chadwick, I. Charania, H. Chen, M. J. Collins, P. Contu, N. D'Arcy, G. Dauphinais, R. De Prins, D. Deschenes, I. Di Luch, S. Duque, *et al.*, Scaling and networking a modular photonic quantum computer, *Nature (London)* **638**, 912 (2025).
- [6] J. Südbeck, S. Steinlechner, M. Korobko, and R. Schnabel, Demonstration of interferometer enhancement through EPR entanglement, *Nat. Photonics* **14**, 240 (2020).
- [7] P. M. Anisimov, G. M. Raterman, A. Chiruvelli, W. N. Plick, S. D. Huver, H. Lee, and J. P. Dowling, Quantum metrology with two-mode squeezed vacuum: Parity detection beats the Heisenberg limit, *Phys. Rev. Lett.* **104**, 103602 (2010).
- [8] C. Couteau, S. Barz, T. Durt, T. Gerrits, J. Huwer, R. Prevedel, J. Rarity, A. Shields, and G. Weihs, Applications of single photons to quantum communication and computing, *Nat. Rev. Phys.* **5**, 326 (2023).
- [9] D. Gottesman, A. Kitaev, and J. Preskill, Encoding a qubit in an oscillator, *Phys. Rev. A* **64**, 012310 (2001).
- [10] B. Q. Baragiola, G. Pantaleoni, R. N. Alexander, A. Karanjai, and N. C. Menicucci, All-Gaussian universality and fault tolerance with the Gottesman-Kitaev-Preskill code, *Phys. Rev. Lett.* **123**, 200502 (2019).
- [11] S. Konno, W. Asavanant, K. Fukui, A. Sakaguchi, F. Hanamura, P. Marek, R. Filip, J.-i. Yoshikawa, and A. Furusawa, Non-Clifford gate on optical qubits by nonlinear feedforward, *Phys. Rev. Res.* **3**, 043026 (2021).
- [12] H. M. Vasconcelos, L. Sanz, and S. Glancy, All-optical generation of states for “Encoding a qubit in an oscillator”, *Opt. Lett.* **35**, 3261 (2010).
- [13] D. J. Weigand and B. M. Terhal, Generating grid states from Schrödinger-cat states without postselection, *Phys. Rev. A* **97**, 022341 (2018).
- [14] M. Eaton, R. Nehra, and O. Pfister, Non-Gaussian and Gottesman-Kitaev-Preskill state preparation by photon catalysis, *New J. Phys.* **21**, 113034 (2019).
- [15] I. Tzitrin, J. E. Bourassa, N. C. Menicucci, and K. K. Sabapathy, Progress towards practical qubit computation using approximate Gottesman-Kitaev-Preskill codes, *Phys. Rev. A* **101**, 032315 (2020).
- [16] K. Takase, K. Fukui, A. Kawasaki, W. Asavanant, M. Endo, J. Yoshikawa, P. van Loock, and A. Furusawa, Gottesman-Kitaev-Preskill qubit synthesizer for propagating light, *npj Quantum Inf.* **9**, 98 (2023).

- [17] P. T. Cochrane, G. J. Milburn, and W. J. Munro, Macroscopically distinct quantum-superposition states as a bosonic code for amplitude damping, *Phys. Rev. A* **59**, 2631 (1999).
- [18] T. C. Ralph, A. Gilchrist, G. J. Milburn, W. J. Munro, and S. Glancy, Quantum computation with optical coherent states, *Phys. Rev. A* **68**, 042319 (2003).
- [19] S. Konno, W. Asavanant, F. Hanamura, H. Nagayoshi, K. Fukui, A. Sakaguchi, R. Ide, F. China, M. Yabuno, S. Miki, H. Terai, K. Takase, M. Endo, P. Marek, R. Filip, P. van Loock, and A. Furusawa, Logical states for fault-tolerant quantum computation with propagating light, *Science* **383**, 289 (2024).
- [20] M. V. Larsen, J. E. Bourassa, S. Kocsis, J. F. Tasker, R. S. Chadwick, C. Gonzalez-Arciniegas, J. Hastrup, C. E. Lopetegui-Gonzalez, F. M. Miatto, A. Motamedi, R. Noro, G. Roeland, R. Baby, H. Chen, P. Contu, I. Di Luch, C. Drago, M. Giesbrecht, T. Grainger, I. Krasnokutskaya, *et al.*, Integrated photonic source of Gottesman-Kitaev-Preskill qubits, *Nature (London)* **642**, 587 (2025).
- [21] M. Dakna, T. Anhu, T. Opatrny, L. Knöll, and D. G. Welsch, Generating Schrödinger-cat-like states by means of conditional measurements on a beam splitter, *Phys. Rev. A* **55**, 3184 (1997).
- [22] A. Ourjoumtsev, R. Tualle-Broui, J. Laurat, and P. Grangier, Generating optical Schrödinger kittens for quantum information processing, *Science* **312**, 83 (2006).
- [23] J. S. Neergaard-Nielsen, Generation of single photons and Schrödinger kitten states of light, Ph.D. thesis, University of Copenhagen, 2008.
- [24] W. Asavanant, K. Nakashima, Y. Shiozawa, J. Yoshikawa, and A. Furusawa, Generation of highly pure Schrödinger's cat states and real-time quadrature measurements via optical filtering, *Opt. Express* **25**, 32227 (2017).
- [25] K. Takase, J.-i. Yoshikawa, W. Asavanant, M. Endo, and A. Furusawa, Generation of optical Schrödinger cat states by generalized photon subtraction, *Phys. Rev. A* **103**, 013710 (2021).
- [26] T. Gerrits, S. Glancy, T. S. Clement, B. Calkins, A. E. Lita, A. J. Miller, A. L. Migdall, S. W. Nam, R. P. Mirin, and E. Knill, Generation of optical coherent-state superpositions by number-resolved photon subtraction from the squeezed vacuum, *Phys. Rev. A* **82**, 031802(R) (2010).
- [27] M. Endo, R. He, T. Sonoyama, K. Takahashi, T. Kashiwazaki, T. Umeki, S. Takasu, K. Hattori, D. Fukuda, K. Fukui, K. Takase, W. Asavanant, P. Marek, R. Filip, and A. Furusawa, Non-Gaussian quantum state generation by multi-photon subtraction at the telecommunication wavelength, *Opt. Express* **31**, 12865 (2023).
- [28] T. Kashiwazaki, T. Yamashima, K. Enbutsu, T. Kazama, A. Inoue, K. Fukui, M. Endo, T. Umeki, and A. Furusawa, Over-8-dB squeezed light generation by a broadband waveguide optical parametric amplifier toward fault-tolerant ultra-fast quantum computers, *Appl. Phys. Lett.* **122**, 234003 (2023).
- [29] K. Hattori, T. Konno, Y. Miura, S. Takasu, and D. Fukuda, An optical transition-edge sensor with high energy resolution, *Supercond. Sci. Technol.* **35**, 095002 (2022).
- [30] E. Schrödinger, Die gegenwärtige situation in der quantenmechanik, *Naturwissenschaften* **23**, 844 (1935).
- [31] J. Hastrup and U. L. Andersen, All-optical cat-code quantum error correction, *Phys. Rev. Res.* **4**, 043065 (2022).
- [32] A. I. Lvovsky and M. G. Raymer, Continuous-variable optical quantum-state tomography, *Rev. Mod. Phys.* **81**, 299 (2009).
- [33] K. Hirota, T. Kashiwazaki, G. Ha, T. Yamashima, P. Jaturaphagorn, T. Suzuki, K. Takahashi, A. Kawasaki, A. Inoue, W. Asavanant, M. Endo, T. Umeki, and A. Furusawa, Generation of 10-dB squeezed light from a broadband waveguide optical parametric amplifier with improved phase locking method, *Opt. Express* **34**, 7958 (2026).
- [34] K. Duivenvoorden, B. M. Terhal, and D. Weigand, Single-mode displacement sensor, *Phys. Rev. A* **95**, 012305 (2017).
- [35] V. Parigi, A. Zavatta, M. Kim, and M. Bellini, Probing quantum commutation rules by addition and subtraction of single photons to/from a light field, *Science* **317**, 1890 (2007).
- [36] Y.-R. Chen, H.-Y. Hsieh, J. Ning, H.-C. Wu, H. L. Chen, Z.-H. Shi, P. Yang, O. Steuernagel, C.-M. Wu, and R.-K. Lee, Generation of heralded optical cat states by photon addition, *Phys. Rev. A* **110**, 023703 (2024).
- [37] J. Fadrný, M. Neset, M. Bielak, M. Ježek, J. Bílek, and J. Fiurášek, Experimental preparation of multiphoton-added coherent states of light, *npj Quantum Inf.* **10**, 89 (2024).
- [38] D. Su, C. R. Myers, and K. K. Sabapathy, Conversion of Gaussian states to non-Gaussian states using photon-number-resolving detectors, *Phys. Rev. A* **100**, 052301 (2019).
- [39] M. Endo, T. Nomura, T. Sonoyama, K. Takahashi, S. Takasu, D. Fukuda, T. Kashiwazaki, A. Inoue, T. Umeki, R. Nehra, P. Marek, R. Filip, K. Takase, W. Asavanant, and A. Furusawa, Data for four-photon subtraction from picosecond squeezed vacuum demonstrating Wigner negativity without loss correction [Dataset], Zenodo, 2026, <https://doi.org/10.5281/zenodo.19533547>.
- [40] T. Umeki, O. Tadanaga, and M. Asobe, Highly efficient wavelength converter using direct-bonded PPZnLN ridge waveguide, *IEEE J. Quantum Electron.* **46**, 1206 (2010).
- [41] T. Kashiwazaki, T. Yamashima, N. Takanashi, A. Inoue, T. Umeki, and A. Furusawa, Fabrication of low-loss quasi-single-mode PPLN waveguide and its application to a modularized broadband high-level squeezer, *Appl. Phys. Lett.* **119**, 251104 (2021).
- [42] O. de Vries, T. Saule, M. Plotner, F. Lücking, T. Eidam, A. Hoffmann, A. Klenke, S. Hadrich, J. Limpert, S. Holzberger, T. Schreiber, R. Eberhardt, I. Pupeza, and A. Tunnermann, Acousto-optic pulse picking scheme with carrier-frequency-to-pulse-repetition-rate synchronization, *Opt. Express* **23**, 19586 (2015).
- [43] A. Ourjoumtsev, H. Jeong, R. Tualle-Broui, and P. Grangier, Generation of optical 'Schrödinger cats' from photon number states, *Nature (London)* **448**, 784 (2007).
- [44] N. Namekata, Y. Takahashi, G. Fujii, D. Fukuda, S. Kurimura, and S. Inoue, Non-Gaussian operation based on photon subtraction using a photon-number-resolving detector at a telecommunications wavelength, *Nat. Photonics* **4**, 655 (2010).
- [45] M. Bouillard, G. Boucher, J. F. Ortas, B. Kanseri, and R. Tualle-Broui, High production rate of single-photon and two-photon Fock states for quantum state engineering, *Opt. Express* **27**, 3113 (2019).
- [46] Y. Eto, A. Koshio, A. Ohshiro, J. Sakurai, K. Horie, T. Hirano, and M. Sasaki, Efficient homodyne measurement of picosecond squeezed pulses with pulse shaping technique, *Opt. Lett.* **36**, 4653 (2011).

- [47] Y. Taguchi and Y. Ozeki, Time-domain analysis on the pulsed squeezed vacuum detected with picosecond pulses, [J. Opt. Soc. Am. B](#) **37**, 1535 (2020).
- [48] N. Killoran, J. Izaac, N. Quesada, V. Bergholm, M. Amy, and C. Weedbrook, Strawberry Fields: A software platform for photonic quantum computing, [Quantum](#) **3**, 129 (2019).
- [49] T. R. Bromley, J. M. Arrazola, S. Jahangiri, J. Izaac, N. Quesada, A. D. Gran, M. Schuld, J. Swinarton, Z. Zabaneh, and N. Killoran, Applications of near-term photonic quantum computers: Software and algorithms, [Quantum Sci. Technol.](#) **5**, 034010 (2020).