

Multipartite quantum correlated bright frequency combs

A. Bensemhoun,¹ S. Cassina², C. Gonzalez-Arciniegas,³ M. F. Melalkia^{1,*}, G. Patera⁵, J. Faugier-Tovar,⁶ Q. Wilmart,⁶ S. Olivier,⁶ A. Zavatta⁷, A. Martin¹, J. Etesse¹, L. Labonté,¹ O. Pfister^{3,4}, V. D'Auria,^{1,†} and S. Tanzilli¹

¹*Université Côte d'Azur, CNRS, Institut de Physique de Nice, 17 rue Julien Lauprêtre, 06200 Nice, France*

²*University of Insubria, Department of Science and High Technology, via Valleggio 11, 22100 Como, Italy*

³*University of Virginia, Physics Department, 382 McCormick Rd, Charlottesville, Virginia 22903, USA*

⁴*Charles L. Brown Department of Electrical and Computer Engineering, University of Virginia, 351 McCormick Road, Charlottesville, Virginia 22903, USA*

⁵*University of Lille, CNRS, UMR 8523-PhLAM-Physique des Lasers Atomes et Molécules, F-59000 Lille, France*

⁶*CEA-LETI, Université Grenoble Alpes, F-38000 Grenoble, France*

⁷*Istituto Nazionale di Ottica (CNR-INO), CNR, Largo Enrico Fermi 6, 50125 Firenze, Italy*



(Received 10 January 2025; accepted 16 July 2025; published 21 August 2025)

This experimental work demonstrates multipartite quantum correlation in bright frequency combs out of a microresonator integrated on silicon nitride operating above its oscillation threshold. Multipartite features, going beyond so far reported two-mode correlation, naturally arise due to a cascade of nonlinear optical processes, making a single-color laser pump sufficient to initiate their generation. Our results show the transition from two-mode to multipartite correlation, witnessed by noise reductions as low as -2.5 dB and -2 dB, respectively, compared to corresponding classical levels. A constant of the movement of the nonlinear interaction Hamiltonian is identified and used to assess the multipartite behavior. Reported demonstrations pave the way to next generation on-chip multipartite sources for quantum technologies applications.

DOI: [10.1103/cc69-5gq2](https://doi.org/10.1103/cc69-5gq2)

I. INTRODUCTION

Optical multipartite states that exploit continuous variable (CV) encoding are key resources for quantum information science [1], playing a crucial role in measurement-based optical quantum computing [2,3], multipartner secure quantum networks [4], and quantum metrology [5,6]. Future ambitious applications will require advanced quantum systems combining the ability to generate and/or manipulate entanglement over a high number of modes with devices' robustness and small footprints. Such scalability can be obtained by spatial [7], temporal [8–11], or spectral [12–15] multiplexing. Spectral multiplexing is particularly attractive as large numbers of entangled modes can be directly generated in a single nonlinear optical source, while eschewing the overhead from sequential processing [1]. In addition, thanks to their different colors, quantum correlated modes can be easily separated and addressed using standard optical components [1].

Another key property for multiple applications is the ability to produce bright quantum entangled beams. For example, in

quantum sensing, the ultimate quantum noise floor of optical interferometry decreases with the average photon number of the probing laser [16]. When the maximum usable light power is reached for whatever reason (record-high in LIGO, living-tissue vulnerability in biomedical imaging, embarked-power limitation on a drone), one can resort to single- or multimode squeezed light to further lower the detection noise. In most cases, however, this requires to displace a low-power squeezed state by reflecting it off a weakly transmitting beam splitter, through which a powerful laser beam is transmitted [17]. This arrangement is extremely wasteful of laser power and requires dedicated locking systems, making it evidently advantageous to generate bright squeezed light directly, with no displacement operations. More generally, working with bright quantum beams gives the possibility of implementing self-homodyne detections to measure the continuous variable properties of entangled modes and, in turn, to access the complete reconstruction of their states via quantum tomography [18,19]. Compared to standard homodyne, this technique does not need an external local oscillator beam at the signal frequency and is particularly convenient when multiple beams at different colors must be detected simultaneously.

There have been extensive explorations of bright bi- and multipartite quantum correlations and entanglement in free-space systems exploiting, in particular, second order nonlinear processes [$\chi^{(2)}$] [15,20–23]. In view of future practical applications, this experimental work aims at investigating the generation of bright multipartite states in miniaturized optical sources integrated on a silicon-based platform.

*Present address: Laboratoire Systèmes Lasers, Ecole Militaire Polytechnique, BP 17 Bordj El Bahri, 16046 Algiers, Algeria.

†Contact author: virginia.dauria@univ-cotedazur.fr

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

Integrated photonics on silicon (Si) and silicon nitride (SiN) is a go-to solution to reduce the physical size of optical systems while benefiting from cascaded optical components on CMOS compatible platforms [24–27]. Frequency combs generated by four wave mixing (FWM) in silicon-based microresonators (microcombs) advantageously bridge multipartite entanglement and integrated resources [28–31]. In below threshold devices pumped with a monochromatic input laser, the quantum state of low power microcombs has been proven to be made of many bipartite states independent of each other [32–34]. The use of multicolor input pumps has recently proved to engineer such bipartite states into multipartite CV correlations [35,36]. Bright frequency combs are efficiently obtained from above threshold microrings and disks [37]. Nevertheless, the complicated dynamics of the system, including the presence of bistability [38] and of morphing behavior [29], makes it not easy to transpose concepts proven in bulk systems into integrated ones ruled by FWM. As a matter of fact, reported works only show CV quantum correlation between two frequency modes out of microresonators pumped with a single input laser [19,39,40]. A signature of multimode features has been measured from second order photon correlation in solitons out of systems far above threshold [41].

This paper goes beyond two-color analysis to demonstrate experimentally that a multipartite behavior is naturally present in bright microcombs from above-threshold SiN microrings. Remarkably, the generation of the multipartite state is initiated by a single monochromatic pump at the device input thanks to a cascade of FWM processes. We predicted this behavior in a previous theoretical work [42]. CV multipartite dynamics is explored here by measuring quantum intensity correlation between up to four bright optical modes generated by an SiN microring. The analysis of intensity correlation makes it possible to observe that, when driving the system progressively farther above threshold, a transition takes place from a bipartite regime to a multipartite one at four modes and beyond. In our work, we go on to demonstrate a fundamental structure to these quantum correlations, epitomized by a multipartite phase invariant that we discover for any FWM interaction. This quantum physics was never observed in microresonators operating below threshold [35,36] because the cascade of FWM processes requires combs with bright frequency components, nor in above-threshold works [19,39,40] as its experimental evidence demands the investigation of more than two modes. Our results pave the way to a new operation regime for quantum microcombs as they push bright correlation in microresonators beyond one-to-one schemes reported so far [19,39,40]. By doing so, they explore their potential for high-dimensional state connectivity and its applications in quantum photonics.

II. MULTIPARTITE CORRELATION IN BRIGHT MICROCOMBS

A. Concepts and experimental setup

In above-threshold FWM, the generation of bright frequency combs is associated with rich dynamics [24,30,37,38]. Figure 1 shows some examples of primary microcombs out

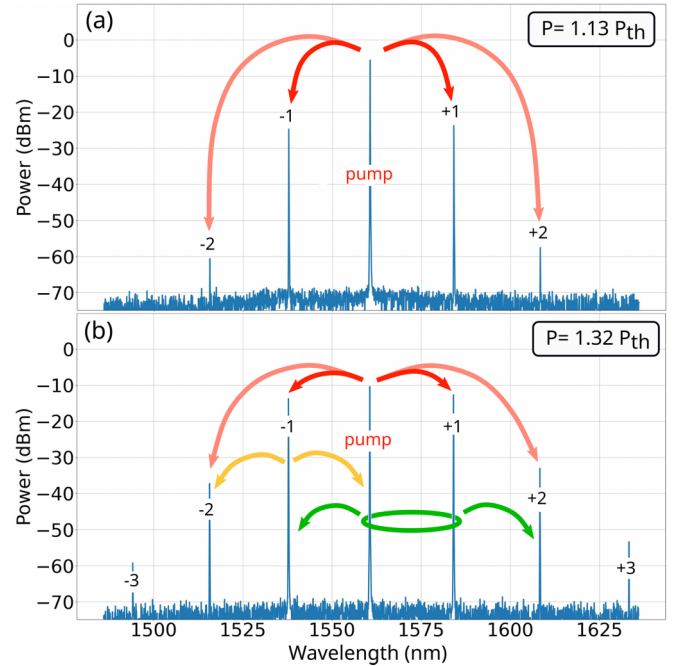


FIG. 1. Measured optical spectra of the comb generated by the SiN microring operating above the oscillation threshold for input pump powers of (a) $P = 1.13 P_{th}$ and (b) $P = 1.32 P_{th}$, where P_{th} is the pump corresponding to the threshold. The figure shows some of the multiple nonlinear processes that can simultaneously take place, including degenerate FWM of the input pump (red) or of a secondary mode (yellow), as well as nondegenerate FWM of two of the comb modes (green). Self- and cross-phase modulations are not shown to preserve the drawing legibility.

of the SiN microresonator, measured by an optical spectrum analyzer (OSA). The system is pumped with a continuous wave (CW) input laser and operates at moderate distance from the threshold, well below the soliton regime [37]; as shown in the figure, bright modes, including the input pump, are all in the telecom *c* band. In these conditions, different FWM processes can take place [42]. The single CW laser generates via degenerate FWM a primary comb of paired modes, symmetrically distributed with respect to the pump line [see modes ± 1 , ± 2 in Fig. 1(a)]. Such bright components provide, in turn, additional pumps for other degenerate or non-degenerate FWM processes [see Fig. 1(b)]. The coexistence of such simultaneous FWMs makes multipartite physics arise in a natural way when increasing the modes' powers. The higher the input pump power, P , the more number of modes, as well as their relative power, increase making their contribution to secondary interactions non-negligible anymore. A formal description of this behavior is provided in the analytical model described in Ref. [42].

The experimental setup used to generate and to investigate the multipartite state is sketched in Fig. 2. Before the microresonator, the setup is entirely made of plug-n-play telecom components based on optical fibers. The master source is a fiber-coupled CW laser (RIO-Orion) delivering 12 mW at 1560.053 nm with a linewidth of ≤ 5 kHz: its central wavelength can be tuned of ± 20 pm with minimal steps of 0.02 pm. The laser output is amplified by a Erbium Doped

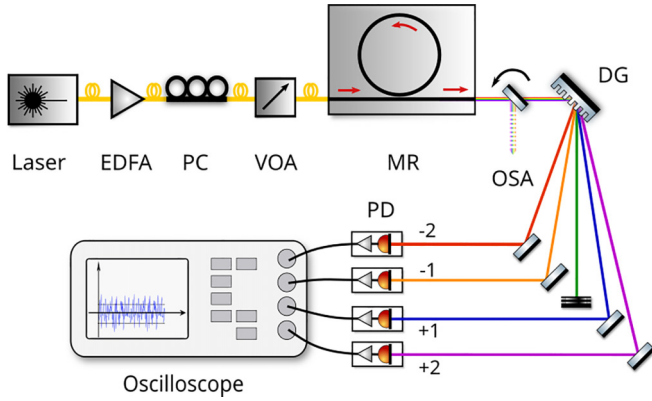


FIG. 2. Experimental setup for the measurement of intensity correlations at the output of an SiN microresonator (MR) above threshold. Modes ± 1 and ± 2 are separated by a grating (DG) thanks to their different colors and their intensities are individually detected (PD). Note that, in the actual experiment the diffraction grating works in transmission, and two mirrors for separated beams are required downstream of it. The corresponding photocurrents are evaluated in the Fourier transform domain and digitally combined as required to evaluate the correlation among different modes. Noise power spectrum of combined signals are used to retrieve nonclassical features. Noises are evaluated with respect to shot noise signals obtained by sending to the detectors laser beams whose powers match those of the modes under scrutiny.

Fiber Amplifier (EDFA), whose working point is optimized to reduce noise contributions. Downstream of the EDFA, residual amplified spontaneous emission is suppressed by a cascade of five wavelength demultiplexing stages (not represented), each providing a noise attenuation of -25 dB and centered around 1560 nm with a width of 100 GHz, well above the laser tunability range. After an in-line variable optical attenuator (VOA), the laser polarization is adjusted (PC) to match the transverse electric (TE) mode of the microresonator, whose propagation losses were measured to be the lowest. The laser light is edge coupled via a microlensed fiber into the microresonator and used as the input pump of the FWM process. Fiber-to-chip coupling losses are of about -3 dB. The microresonator is fabricated in-house at the CEA-LETI on Si_3N_4 [43,44]. It is a single-pass microring (MR, see Fig. 2) with a radius of $112\text{ }\mu\text{m}$ (corresponding to a free spectral range of 200 GHz) coupled to a straight waveguide set at a gap of 300 nm. The width and height of all waveguides are, respectively, $1.9\text{ }\mu\text{m}$ and 800 nm. The MR is isolated against mechanical and thermal instabilities and mounted on a thermoelectric cooler (TEC) that stabilizes its temperature at 68.5°C . The microresonator has an intrinsic $Q \sim 10^7$ with a cavity linewidth of the order of ~ 630 MHz and operates in the overcoupling regime, with a measured pump oscillation threshold $P_{\text{th}} = 53 \pm 3$ mW. Note that as in most works on nonlinear optical resonators pumped in CW regime [19,32–34,39,40], the laser linewidth is much smaller than the cavity resonance: such a narrow pump laser guarantees, in absence of secondary FWM effects, the generation of one-to-one correlations among symmetrical modes ($\pm 1, \pm 2, \dots$) [24,29]. For the scopes of our analysis, this allows us to be able to start from a bipartite regime and to follow its transformation when

secondary FWM processes become non-negligible. For each measurement point, the laser wavelength is adjusted to guarantee its perfect resonance with the cavity resonance labeled as $k=0$. No active locking to the cavity resonances is used on the pump laser.

Its output is collected in free space with a microscope objective. Since the chip facets are not provided with antireflection coating, the beams transmission, η_c , is reduced due to the Fresnell losses, which give $\eta_F = (\frac{n_{\text{SiN}} - n_{\text{air}}}{n_{\text{SiN}} + n_{\text{air}}})^2 \approx 0.90$, with n_{SiN} and n_{air} , the SiN and air refractive indexes. A low-loss diffraction grating splits the comb's optical components and sends them to separate photodiodes (PD). The grating works in transmission and has a groove density of ≈ 1000 grooves/mm. We measured for our grating an isolation between next modes' wavelengths is of 45 dB and a transmission $\eta_g \gtrsim 0.94$ for all modes of interest. A flip mirror before the grating is used to send the light to the OSA to check the MR's operation regime before any measurement. Downstream of the grating, each beam is sent to two dielectric mirrors and a lens that direct it toward one of the photodiode (not depicted in Fig. 2): the overall transmission of these optical components is $\eta_{\text{opt}} \approx 0.97$. Four modes (± 1 and ± 2) are simultaneously detected. Each photodiode mounts a head with an active surface of diameter $300\text{ }\mu\text{m}$ and a responsivity at the telecom wavelength of 1.1 A/W , corresponding to a measured detection efficiency $\eta_{\text{pd}} \approx 0.88$. Overall optical losses per mode lead to a detection efficiency $\eta = \eta_F \eta_g \eta_{\text{opt}} \eta_{\text{pd}} \approx 72\%$. To evaluate the intensity correlation, the PD photocurrents are acquired by a fast oscilloscope to be evaluated in the Fourier domain and digitally combined. Noise power spectra of combined signals are used to retrieve nonclassical features. All noise spectra are evaluated with respect to shot noise levels (SNL) that are obtained by sending to the detectors four laser beams whose mean intensities match those of the modes under scrutiny. Details on the data acquisition, digital treatment, and calibration with respect to the SNL are reported in the Appendix.

B. Quantifying multipartite intensity correlation

While different witnesses exist for multipartite correlations between field quadratures [1,35,36,42], a witness for multipartite intensity correlations is not present in the literature. We can, however, construct it by looking at the Hamiltonian symmetries. The overall system dynamics is described by the general FWM Hamiltonian, $\hat{H}_{\text{FWM}}$. It contains all processes by which two photons in frequency modes p and q annihilate to create paired photons in modes s and r [29,42]:

$$\hat{H}_{\text{FWM}} = \sum_{r,s,p,q} \delta_{r+s,p+q} \hat{a}_r^\dagger \hat{a}_s^\dagger \hat{a}_p \hat{a}_q, \quad (1)$$

where $\hat{a}_i$ and $\hat{a}_i^\dagger$ are the bosonic operators associated to intracavity modes $i = 0$ (the pump), $\pm 1, \pm 2, \dots$, and satisfy $[\hat{a}_i, \hat{a}_j^\dagger] = \delta_{i,j}$ and $[\hat{a}_i, \hat{a}_j] = 0$. The sum over all modes describes the fact that, above threshold, any pair of frequency modes can play the role of pumps for another FWM process, provided the energy conservation, expressed by the Kronecker delta, is respected (i.e., $\omega_r + \omega_s = \omega_p + \omega_q$, where ω is the mode optical frequency). Cross- and self-phase correlation terms are included in $\hat{H}_{\text{FWM}}$ [42]. The nonlinear interaction

coupling constant is set to 1 for simplicity. As detailed in Ref. [42], the Hamiltonian $\hat{H}_{\text{FWM}}$ allows constructing a set of linearized Langevin equations that link the quantum dynamics of each $\hat{a}_i$, to that of the other modes, via a coupling matrix whose elements depend on the classical amplitudes of the interacting fields, $\{A_i\}$: strong coupling are associated to high mean amplitudes. This formally translates the progressive appearance of multimode correlations when the fields at different frequency components become brighter.

Given $2M + 1$ interacting modes, it is easy to prove that the Hamiltonian of Eq. (1) admits as a constant of motion the observable associated with the Hermitian operator:

$$\hat{\mathcal{C}} = \sum_{k=-M}^{+M} k \hat{n}_k = \hat{n}_1 - \hat{n}_{-1} + 2(\hat{n}_2 - \hat{n}_{-2}) + \dots, \quad (2)$$

with $\hat{n}_k = \hat{a}_k^\dagger \hat{a}_k$ and $k \in \mathbb{N}$. In experiments, $\hat{\mathcal{C}}$ can directly be measured by simultaneously acquiring the intensities of all involved modes, $I_k \propto \hat{n}_k$, as functions of time. The operator $\hat{\mathcal{C}}$ commutes with the time independent Hamiltonian $\hat{H}_{\text{FWM}}$, meaning that the associated physical quantity is conserved throughout the system evolution. In the case of a nonseeded microresonator, apart from the input pump ($k = 0$), all modes $\pm k$ are initially in the vacuum state, leading to both $\langle \hat{\mathcal{C}} \rangle$ and $\text{Var}(\hat{\mathcal{C}}) = \langle \hat{\mathcal{C}}^2 \rangle - \langle \hat{\mathcal{C}} \rangle^2$ equal to 0. These values are expected to remain constant even after the FWM interaction, when the modes are populated by bright beams. In particular, the variance of $\hat{\mathcal{C}}$ contains both quantum noise associated with individual modes and covariance terms describing intermodal quantum correlation. It can be written as

$$\text{Var}(\hat{\mathcal{C}}) = \text{Var}(\hat{\mathcal{C}}_{BP}) + \text{Cov}(\hat{\mathcal{C}}_{MP}), \quad (3)$$

where

$$\begin{aligned} \text{Var}(\hat{\mathcal{C}}_{BP}) &= \sum_{k=0}^{+M} k^2 \text{Var}(\hat{n}_k - \hat{n}_{-k}) \\ \text{Cov}(\hat{\mathcal{C}}_{MP}) &= \sum_{\substack{k,l=-M \\ k \neq l}}^{+M} k l (\langle \hat{n}_k \hat{n}_l \rangle - \langle \hat{n}_k \rangle \langle \hat{n}_l \rangle). \end{aligned} \quad (4)$$

In the previous expressions, $\text{Var}(\hat{\mathcal{C}}_{BP})$ covers all bipartite correlations between paired symmetric modes $\pm k$, while $\text{Cov}(\hat{\mathcal{C}}_{MP})$ contains all correlations between nonsymmetric modes: such a covariance term is 0 for a state containing only bipartite correlations among symmetric modes, while it is in general non-negligible for multipartite systems. In the following, a $\text{Var}(\hat{\mathcal{C}})$ below the shot noise level is used to witness quantum intensity correlation involving all interacting modes. This choice allows introducing a measure of multipartite correlation in intensity, differing from other common correlation witnesses that are based on field quadratures [1]. Also note that $\hat{\mathcal{C}}$ represents one of the symmetries of the FWM Hamiltonian. Its associated evolution operator, $\mathcal{U}(\theta) = \prod_{k=-M}^M e^{i\theta k \hat{n}_k}$, is the operator that simultaneously phase shifts all modes k , each by an angle $k\theta$. H is invariant under $\mathcal{U}(\theta)$, $\forall \theta$.

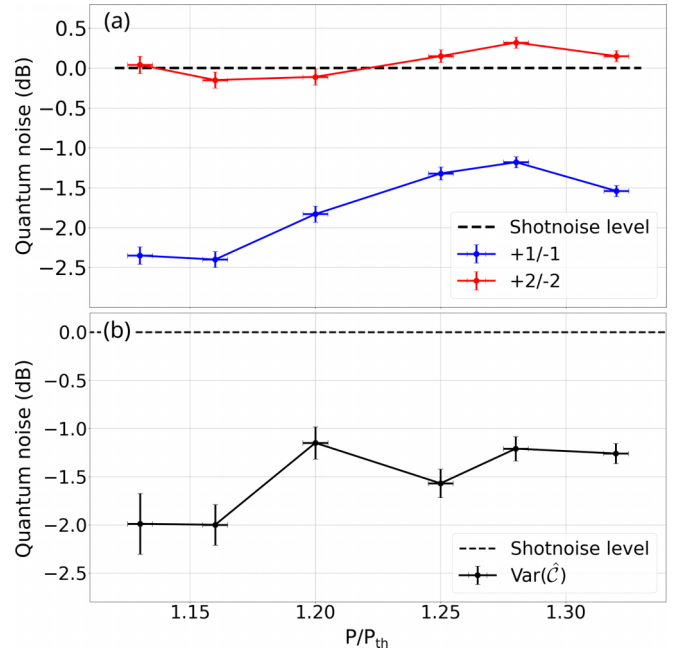


FIG. 3. (a) Quantum correlation between $+1/-1$ (blue) and $+2/-2$ (red) normalized to the shotnoise level as a function of the input pump power divided by the oscillation threshold pump, P_{th} . (b) Quantum noise normalized on the shotnoise level of the multipartite observable as a function of the input pump power. Experimental data are not corrected by quantum efficiency, only for residual electronic noise. All values refer to noise level evaluated at the same analysis frequency (≈ 4 MHz). For each of the curves the shot noise has been experimentally obtained by replacing each of the beams under scrutiny with a coherent beam of the same mean power. In all the figures, P is the input power coupled to the microresonator.

III. RESULTS

Figure 3 shows intensity correlation between modes ± 1 [(a), blue], modes ± 2 [(a), red], and the variance of $\hat{\mathcal{C}}$ on the ensemble of the four modes (b). Curves are plotted as functions of the input pump power, P , normalized to P_{th} . Bimodal intensity correlations are expressed in terms of the noise on their intensity difference that is proportional to $\text{Var}(\hat{n}_k - \hat{n}_{-k})$. This approach has been already adopted to study modes ± 1 in bright microcombs [39,40].

At low pump powers, the leading nonlinear process is the degenerate FWM of the input laser beam that produces symmetric modes [see Fig. 1(a)]. This leads to one-to-one correlation between modes ± 1 , that are, to first approximation, independent of ± 2 . This regime corresponds to the optimal intensity correlation of -2.5 dB at low P/P_{th} plotted in Fig. 3(a). When increasing P , the intensity correlation of modes ± 1 is degraded, eventually reaching ≈ -1 dB. Such a change can be associated with the appearance of cascaded FWM processes that link the dynamics of ± 1 to that of ± 2 , and then to that of ± 3 , and so on to higher-order modes [42]. Under these conditions, multipartite correlation arise and detecting only two modes corresponds to tracing out a part of the multipartite system; the measurement of $\text{Var}(\hat{n}_1 - \hat{n}_{-1})$ thus becomes insufficient to fully characterize intensity correlation available at the microring output. This

behavior is consistent theoretical results from two-mode correlations [42]. Figure 1(a) also shows the correlation between modes ± 2 . Although measured noise reductions are lower than for modes ± 1 , intensity correlations in modes ± 2 also show a decrease of correlation with the increase of the pump power. Lesser noise reductions in such modes are due to reduced signal-to-noise ratio when measuring weak intensity signals.

The measurement of $\hat{C}$ combines these results for the ensemble of the four modes. As expected, $\text{Var}(\hat{C}) < \text{SNL}$ for all considered input pump powers, and it goes down to -2 dB at low P . Remarkably, this stays true despite the fact that, for this microcomb, $\text{Var}(\hat{n}_2 - \hat{n}_{-2})$ stays high and, eventually, passes above its SNL. Interestingly, experimental results on $\text{Var}(\hat{C})$ can be compared with $\text{Var}(\hat{C}_{BP})$, i.e., with noise levels that are expected for $\hat{C}$ when the microcomb quantum state contains only bipartite correlations [see Eq. (3)]. For each P/P_{th}, these can be computed by injecting in Eq. (4) measured $\text{Var}(\hat{n}_1 - \hat{n}_{-1})$ and $\text{Var}(\hat{n}_2 - \hat{n}_{-2})$. Obtained $\text{Var}(\hat{C}_{BP})$ evolve from -1.6 dB at $P/P_{th} \sim 1.1$ to $\approx +1$ dB when $P/P_{th} \geq 1.25$, and are systematically higher than our measured $\text{Var}(\hat{C})$: the difference between $\text{Var}(\hat{C}_{BP})$ and $\text{Var}(\hat{C})$ confirm the presence of non-negligible covariance terms $\text{Cov}(\hat{C}_{MP})$ in Eq. (3) and, in turn, of cross correlations among nonsymmetric modes. This quantum feature was out of the reach of former experimental works on bright microcombs that were restricted to the analysis of modes ± 1 [19,39,40].

For the microcomb of Fig. 3(b), an abrupt decrease of $\text{Var}(\hat{C})$ can be observed when P/P_{th} goes from 1.15 to 1.20 (from -2 dB to ≈ -1 dB). Correspondingly, modes ± 3 start being non-negligible in the microcombs. Although weaker than modes ± 1 (typically -40 dB), ± 3 terms act as seeds to secondary FWM processes, thus connecting their dynamics to that of other modes. In this regard, note that the measurement of $\hat{C}$ with $M=4$ neglects contributions from symmetric modes ± 3 and higher; disregarding them is equivalent to trace-out information on a part of an entangled state of higher dimension. The increase of $\text{Var}(\hat{C})$ observed when $P/P_{th} \geq 1.20$, thus gives an indirect proof of a richer dynamics where more than four modes are quantum correlated.

To dig into the four-mode regime, a second set of measurements is performed at lower pump values, in a regime where modes ± 3 are expected to be negligible. Note that the microcomb under study in the low power regime is not the same as for measurements in Fig. 3. This is due to optical bistability in resonant FWM processes, i.e., to the existence of two stable stationary solutions in above-threshold microresonators [38]; depending on the chosen range of input powers, the system spontaneously emits one microcomb or the other. As explained, the modes' coupling and, thus, the actual levels of their correlation depend on the field classical amplitudes, $\{A_i\}$ (with $A_i \in \mathbb{C}$) that are expected to be different from a comb to the other [42]. Nevertheless the overall dynamic should follow a behavior analogous to the one of Fig. 3. This can be seen in Figs. 4(a) and 4(b), that show intensity correlation as functions of low P/P_{th} for symmetric modes; the behavior of $\text{Var}(\hat{n}_k - \hat{n}_{-k})$ qualitatively confirms what has already been observed for the microcomb of Fig. 3. Lower measured correlation levels are due to the microcomb emission that is not bright enough to make residual electronics noise negligible.

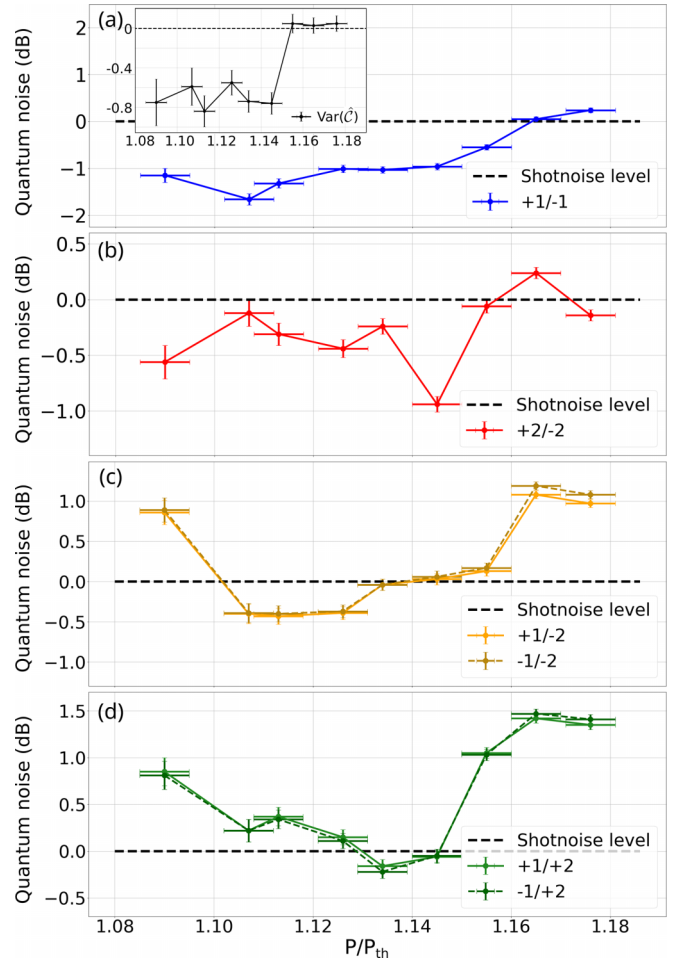


FIG. 4. Quantum correlation normalized to SNL as functions of the input pump power for modes $+1/-1$ (a), $+2/-2$ (b), and between nonsymmetric modes (c) and (d). The inset reports the corresponding $\text{Var}(\hat{C})$.

Figures 4(c) and 4(d) illustrate the dynamics of cross correlation between nonsymmetric modes in terms of $\text{Var}(\hat{n}_i - \hat{n}_j)$, with $i \neq j = \pm 1, \pm 2$. At very low powers, secondary FWM processes are not yet well established and quantum correlation only takes place between symmetric modes generated by degenerate FWM of the input pump. The reduction of $\text{Var}(\hat{n}_i - \hat{n}_j)$ observed when increasing P bears witness to the birth of multipartite correlation: the links between modes initially independent lead to noises below the classical level at intermediate values of P , as shown in Figs. 4(c) and 4(d). This regime can be considered as showing four-partite features. At $P/P_{th} \approx 1.15$, the growth of modes ± 3 reduces cross- as well as symmetric-mode correlation; this represents a confirmation of a six-partite dynamics for which neglecting information of modes above ± 2 corresponds to tracing out a part of a larger quantum correlated system. The differences between curves of Figs. 4(c) and 4(d) can be justified by optical losses and dispersion. Dispersion, in particular, affects the position of cavity resonance and can induce some asymmetries, especially experienced by modes $+2$ and -2 , that are the most spaced in terms of optical frequencies. At the same time, in CW regime, only modes that are distributed in a perfectly

symmetric way with respect to the input pump can satisfy the energy conservation for the initial degenerate FWM process: the mismatch of mode positions and cavity resonances affects the dynamics of the system and modify the strength of quantum correlations between symmetric and asymmetric modes [42].

As a final remark, note that the behavior of symmetric and cross-mode correlation is well summarized by the $\text{Var}(\hat{C})$ evolution reported in the figure inset. Its steep change at higher power reproduces that of cross correlation and leads to values above the SNL at high powers. For this microcomb, in addition to the appearance of modes ± 3 , it is pertinent to note that the higher P is, the more probable the transition towards the comb corresponding to higher power becomes, thus potentially affecting the system quantum correlation.

IV. CONCLUSION

The presented experimental work demonstrates multipartite features in bright frequency combs out of an SiN microring exploited above its oscillation threshold. Multipartite features appear spontaneously due to a cascade of secondary FWM processes that exploit the primary comb out of a degenerate FWM of a monochromatic pump laser. The dynamics of quantum correlation in four bright modes is investigated by measuring their one-to-one correlation as well as in terms of the noise of the operator $\hat{C}$ that represents a symmetry of the FWM Hamiltonian and that combines all modes. Although experimental data refer to the first four modes appearing in the primary combs, all obtained results indicate that higher order connectivity can be expected when including in the analysis subsequent components that appear when working farther away from the threshold. Experimental evidence of multipartite correlation in bright microcombs and the study of their dynamics have never been performed before. Presented results, as well as the theoretical and experimental tools dedicated to their measurement, paves the way to the exploitation of bright microcombs for future quantum technology applications in CV multipartite quantum photonics, with a simple experimental configuration, where just one external CW pump laser is needed to feed the system. The experimental tools and concepts demonstrated in this work apply in a safe way to primary combs of any size. The actual number of components in the microcomb depends on the device geometry and performance and scales with the input pump [37], until the generation of the secondary comb is reached, eventually leading to a soliton regime. A theoretical [30] and experimental paper [41] have already discussed the presence of quantum correlation among the soliton frequency components. This result, combined with ours, shows that no strict limitation to the size of the correlated states are present in the shift from the primary comb to the soliton. Nevertheless, compared to other systems [26], the correlations of modes is predetermined by the classical emission comb: in this sense, our work thus opens to new theoretical and experimental investigations, in which the analysis of correlation between the bright modes is conducted with the perspective of engineering the quantum correlations as done for instance in below threshold systems [35,36].

ACKNOWLEDGMENTS

This work has been conducted within the framework of the project OPTIMAL granted by the European Union by means of the Fond Européen de développement régional (FEDER). The authors also acknowledge financial support from the Agence Nationale de la Recherche (ANR) through the projects SPHIFA (ANR-20-CE47-0012) and OQulus (ANR-22-PETQ-0013), and from Academy RISE of Université Côte d'Azur. V.D.A. thanks the Institut Universitaire de France for the support. O.P. acknowledges support from CNRS, Fondation Doebelin, and US NSF Grants No. ECCS-2219760 and No. PHY-2112867.

A.B. mounted the entire experimental setup and took care of data acquisition and analysis, with the help of S.C., M.F.M., and A.Z.; C.G.A. found the original theoretical tool for witnessing multipartite intensity correlation, with the help of O.P.; S.O. and Q.W. designed and fabricated the SiN integrated structures; V.D.A. and O.P. conducted the data analysis and interpretation, with the help of A.M., J.E., G.P., L.L., and S.T. All authors read, discussed, and contributed to the writing, reviewing, and editing of the manuscript. V.D.A., L.L., and S.T. coordinated and managed the project, ensuring its successful completion.

The authors declare no competing interests.

DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX: DATA ACQUISITION, TREATMENT, AND NORMALIZATION

Downstream of the grating, the beams are simultaneously detected by four photodiodes in-house built at INPHYNI and LENS-INO. The detection circuits allow to split the DC and AC components of photocurrents. The DC outputs of the photodiodes are directly proportional to the mean optical power of input beams. They are sent to an oscilloscope with a low bandwidth of 250 MHz, to monitor the individual powers of the four modes: these values are required to retrieve corresponding shot noise levels that will be used in the noise normalization. The AC amplification circuit is in a transimpedance configuration. We took care of balancing the performances of detectors measuring symmetric modes. Accordingly for the pair of detectors dedicated to modes ± 1 have an AC bandwidth at -3 dB of 69 MHz while the detectors for modes ± 2 have a AC bandwidth of 10 MHz. The choice of the analysis frequency, 4 MHz, is based on the individual photodiodes detection spectra and allows us to obtain the best signal-to-noise ratios with respect to the detectors' noise floors.

The detectors' photocurrents are sent to the four channels of a fast oscilloscope with a bandwidth of 4 GHz and acquired at a sampling rate of 5 GSamples/s: due to the involved bandwidths, to avoid the effect of fast electronic noise components, we used a low-pass filter option on the oscilloscope to cut off frequency above 20 MHz. For each measurement point, we made 50 subsequent acquisitions of

100 μ s. Acquired data are then treated with Python. Their fast Fourier transforms are computed with a Python FFT library, and then combined according to the desired observables. For each analysis frequency, Ω , digital data processing is used to compensate possible power unbalance and temporal delays between photocurrents associated to the detection of symmetrical modes $\pm k$. This is done by evaluating the difference $\Delta_k(\Omega) = (I_k(\Omega) - \alpha_k(\Omega)I_{-k}(\Omega))$ with $I_j(\Omega)$ the component at frequency Ω of the noise spectrum of the photocurrent associated to the detection of mode “j.” $\alpha_k(\Omega)$ is a complex parameter that is used for balancing the signals’ power balancing and synchronization. The mean value of α_k is given by the DC components of the photocurrents, that are directly proportional to the mean power of input beams. The phase of $\alpha_k(\Omega)$ account for delay in temporal domain and in principle depends on Ω ; its value is found by asking the algorithm to minimize $\Delta_k(\Omega)$. In the measurement of the observable $\hat{C}$, in addition to α_1 and α_2 , that are determined from $\Delta_1(\Omega)$ and $\Delta_2(\Omega)$, a parameter $\beta(\Omega)$ is used to comply with the overall delay between the two pairs of photodiodes. Based on the chosen analysis frequency, we implement the procedure for the determination of the three parameters at $\Omega = 4$ MHz. Once the digital parameters are set, for each measurement

point, the program makes an average of data obtained over the 50 subsequent acquisitions and then compute the variance of obtained noise spectra. As a final step, correlation spectra are normalized to assess noise reduction (or increase) with respect to the shot noise associated to coherent beams. Shot noise traces come from the nonamplified laser, whose noise features have been carefully measured to exclude the presence of undesired classical noises within the detection bandwidths of our photodiodes. We calibrated the shot noise by sending on the photodiodes the laser and measuring the spectrum of noise of the corresponding photocurrent. To do so, the photocurrents undergo the same acquisition and digital data processing as for the signal beams and their noise features in Fourier domain are acquired. The shot noise data are acquired as functions of the laser mean powers, and the calibration is performed independently for each photodiode. For each of the investigated observables ($\hat{C}$ or $\hat{n}_k - \hat{n}_{-k}$), we computed the analytic expression of the quantum noise that would be obtained by replacing each of the modes $\pm 1, \pm 2$, with a coherent beam of the same optical power. To retrieve its actual experimental value, we used in the obtained formulas the shot noise levels that we got from the calibration: the obtained result is used to normalized the noises on $\text{Var}(\hat{C})$ or on $\text{Var}(\hat{n}_k - \hat{n}_{-k})$.

-
- [1] C. Fabre and N. Treps, Modes and states in quantum optics, *Rev. Mod. Phys.* **92**, 035005 (2020).
 - [2] O. Pfister, Continuous-variable quantum computing in the quantum optical frequency comb, *J. Phys. B* **53**, 012001 (2020).
 - [3] W. Asavanant and A. Furusawa, *Optical Quantum Computers: A Route to Practical* (AIP Publishing LLC, Melville, New York, 2022).
 - [4] M. Epping, H. Kampermann, C. Macchiavello, and D. Bruß, Multi-partite entanglement can speed up quantum key distribution in networks, *New J. Phys.* **19**, 093012 (2017).
 - [5] X. Guo, C. R. Breum, J. Borregaard, S. Izumi, M. V. Larsen, T. Gehring, M. Christandl, J. S. Neergaard-Nielsen, and U. L. Andersen, Distributed quantum sensing in a continuous-variable entangled network, *Nat. Phys.* **16**, 281 (2020).
 - [6] Z. Zhang, C. You, O. S. Magaña-Loaiza, R. Fickler, R. de J. León-Montiel, J. P. Torres, T. S. Humble, S. Liu, Y. Xia, and Q. Zhuang, Entanglement-based quantum information technology: a tutorial, *Adv. Opt. Photon.* **16**, 60 (2024).
 - [7] K. Zhang, W. Wang, S. Liu, X. Pan, J. Du, Y. Lou, S. Yu, S. Lv, N. Treps, C. Fabre, and J. Jing, Reconfigurable hexapartite entanglement by spatially multiplexed four-wave mixing processes, *Phys. Rev. Lett.* **124**, 090501 (2020).
 - [8] S. Yokoyama, R. Ukai, S. C. Armstrong, C. Sornphiphatphong, T. Kaji, S. Suzuki, J. Yoshikawa, H. Yonezawa, N. C. Menicucci, and A. Furusawa, Ultra-large-scale continuous-variable cluster states multiplexed in the time domain, *Nat. Photon.* **7**, 982 (2013).
 - [9] J.-i. Yoshikawa, S. Yokoyama, T. Kaji, C. Sornphiphatphong, Y. Shiozawa, K. Makino, and A. Furusawa, Invited article: Generation of one-million-mode continuous-variable cluster state by unlimited time-domain multiplexing, *APL Photon.* **1**, 060801 (2016).
 - [10] 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).
 - [11] 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).
 - [12] M. Pysher, Y. Miwa, R. Shahrokhshahi, R. Bloomer, and O. Pfister, Parallel generation of quadripartite cluster entanglement in the optical frequency comb, *Phys. Rev. Lett.* **107**, 030505 (2011).
 - [13] J. Roslund, R. M. De Araujo, S. Jiang, C. Fabre, and N. Treps, Wavelength-multiplexed quantum networks with ultrafast frequency combs, *Nat. Photon.* **8**, 109 (2014).
 - [14] M. Chen, N. C. Menicucci, and O. Pfister, Experimental realization of multipartite entanglement of 60 modes of a quantum optical frequency comb, *Phys. Rev. Lett.* **112**, 120505 (2014).
 - [15] F. A. S. Barbosa, A. S. Coelho, L. F. Muñoz-Martínez, L. Ortiz-Gutiérrez, A. S. Villar, P. Nussenzveig, and M. Martinelli, Hexapartite entanglement in an above-threshold optical parametric oscillator, *Phys. Rev. Lett.* **121**, 073601 (2018).
 - [16] C. M. Caves, Quantum-mechanical radiation-pressure fluctuations in an interferometer, *Phys. Rev. Lett.* **45**, 75 (1980).
 - [17] M. G. Paris, Displacement operator by beam splitter, *Phys. Lett. A* **217**, 78 (1996).
 - [18] F. A. Barbosa, A. S. Coelho, K. N. Cassemiro, P. Nussenzveig, C. Fabre, M. Martinelli, and A. S. Villar, Beyond spectral homodyne detection: Complete quantum measurement of spectral modes of light, *Phys. Rev. Lett.* **111**, 200402 (2013).
 - [19] R. A. Kögler, G. C. Rickli, R. R. Domenegueti, X. Ji, A. L. Gaeta, M. Lipson, M. Martinelli, and P. Nussenzveig, Quantum state tomography in a third-order integrated optical parametric oscillator, *Opt. Lett.* **49**, 3150 (2024).

- [20] A. Heidmann, R. J. Horowicz, S. Reynaud, E. Giacobino, C. Fabre, and G. Camy, Observation of quantum noise reduction on twin laser beams, *Phys. Rev. Lett.* **59**, 2555 (1987).
- [21] S. Feng and O. Pfister, Quantum interference of ultrastable twin optical beams, *Phys. Rev. Lett.* **92**, 203601 (2004).
- [22] A. S. Villar, L. S. Cruz, K. N. Cassemiro, M. Martinelli, and P. Nussenzveig, Generation of bright two-color continuous variable entanglement, *Phys. Rev. Lett.* **95**, 243603 (2005).
- [23] A. Coelho, F. Barbosa, K. N. Cassemiro, A. S. Villar, M. Martinelli, and P. Nussenzveig, Three-color entanglement, *Science* **326**, 823 (2009).
- [24] Y. K. Chembo, Kerr optical frequency combs: theory, applications and perspectives, *Nanophotonics* **5**, 214 (2016).
- [25] L. S. Madsen, F. Laudenbach, M. F. Askarani, F. Rortais, T. Vincent, J. F. F. Bulmer, F. M. Miatto, L. Neuhaus, L. G. Helt, M. J. Collins, A. E. Lita, T. Gerrits, S. W. Nam, V. D. Vaidya, M. Menotti, I. Dhand, Z. Vernon, N. Quesada, and J. Lavoie, Quantum computational advantage with a programmable photonic processor, *Nature (London)* **606**, 75 (2022).
- [26] L. Labonté, O. Alibert, V. D'Auria, F. Dautre, J. Etesse, G. Sauder, A. Martin, E. Picholle, and S. Tanzilli, Integrated photonics for quantum communications and metrology, *PRX Quantum* **5**, 010101 (2024).
- [27] Y. Zhang, M. Menotti, K. Tan, V. Vaidya, D. Mahler, L. Helt, L. Zatti, M. Liscidini, B. Morrison, and Z. Vernon, Squeezed light from a nanophotonic molecule, *Nat. Commun.* **12**, 2233 (2021).
- [28] Y. K. Chembo and N. Yu, Modal expansion approach to optical-frequency-comb generation with monolithic whispering-gallery-mode resonators, *Phys. Rev. A* **82**, 033801 (2010).
- [29] E. Gouzien, L. Labonté, J. Etesse, A. Zavatta, S. Tanzilli, V. D'Auria, and G. Patera, Hidden and detectable squeezing from microresonators, *Phys. Rev. Res.* **5**, 023178 (2023).
- [30] M. A. Guidry, D. M. Lukin, K. Y. Yang, and J. Vuckovic, Multimode squeezing in soliton crystal microcombs, *Optica* **10**, 694 (2023).
- [31] M. Sloan, A. Viola, M. Liscidini, and J. E. Sipe, High gain squeezing in lossy resonators: an asymptotic field approach, *Phys. Rev. A* **111**, 063502 (2025).
- [32] V. D. Vaidya, B. Morrison, L. Helt, R. Shahrokshahi, D. Mahler, M. Collins, K. Tan, J. Lavoie, A. Repington, M. Menotti *et al.*, Broadband quadrature-squeezed vacuum and nonclassical photon number correlations from a nanophotonic device, *Sci. Adv.* **6**, eaba9186 (2020).
- [33] Z. Yang, M. Jahanbozorgi, D. Jeong, S. Sun, O. Pfister, H. Lee, and X. Yi, A squeezed quantum microcomb on a chip, *Nat. Commun.* **12**, 4781 (2021).
- [34] M. Jahanbozorgi, Z. Yang, S. Sun, H. Chen, R. Liu, B. Wang, and X. Yi, Generation of squeezed quantum microcombs with silicon nitride integrated photonic circuits, *Optica* **10**, 1100 (2023).
- [35] Z. Wang, K. Li, Y. Wang, X. Zhou, Y. Cheng, B. Jing, F. Sun, J. Li, Z. Li, B. Wu, Q. Gong, Q. He, B.-B. Li, and Q.-F. Yang, Large-scale cluster quantum microcombs, *Light Sci. Appl.* **14**, 164 (2025).
- [36] X. Jia, C. Zhai, X. Zhu, C. You, Y. Cao, X. Zhang, Y. Zheng, Z. Fu, J. Mao, T. Dai, L. Chang, X. Su, Q. Gong, and J. Wang, Continuous-variable multipartite entanglement in an integrated microcomb, *Nature (London)* **639**, 329 (2025).
- [37] M. Karpov, M. H. Pfeiffer, H. Guo, W. Weng, J. Liu, and T. J. Kippenberg, Dynamics of soliton crystals in optical microresonators, *Nat. Phys.* **15**, 1071 (2019).
- [38] C. Godey, I. V. Balakireva, A. Coillet, and Y. K. Chembo, Stability analysis of the spatiotemporal lugiato-lefever model for kerr optical frequency combs in the anomalous and normal dispersion regimes, *Phys. Rev. A* **89**, 063814 (2014).
- [39] A. Dutt, K. Luke, S. Manipatruni, A. L. Gaeta, P. Nussenzveig, and M. Lipson, On-chip optical squeezing, *Phys. Rev. Appl.* **3**, 044005 (2015).
- [40] Y. Shen, P.-Y. Hsieh, S. K. Sridhar, S. Feldman, Y.-C. Chang, T. A. Smith, and A. Dutt, Strong nanophotonic quantum squeezing exceeding 3.5 db in a foundry-compatible kerr microresonator, *Optica* **12**, 302 (2025).
- [41] M. A. Guidry, D. M. Lukin, K. Y. Yang, R. Trivedi, and J. Vučković, Quantum optics of soliton microcombs, *Nat. Photon.* **16**, 52 (2022).
- [42] A. Bensemhoun, C. Gonzalez-Arciniegas, O. Pfister, L. Labonté, J. Etesse, A. Martin, S. Tanzilli, G. Patera, and V. d'Auria, Multipartite entanglement in bright frequency combs out of microresonators, *Phys. Lett. A* **493**, 129272 (2024).
- [43] H. E. Dirani, L. Youssef, C. Petit-Etienne, S. Kerdiles, P. Grosse, C. Monat, E. Pargon, and C. Sciancalepore, Ultralow-loss tightly confining Si₃N₄ waveguides and high-Q microresonators, *Opt. Express* **27**, 30726 (2019).
- [44] Q. Wilmart, S. Guerber, J. Faugier-Tovar, Y. Ibrahimi, C. Petit-Etienne, L. Youssef, C. Socquet-Clerc, A. Myko, K. Ribaud, F. Duport, E. Pargon, and F. V. Dijk, A device library for the ultra-low loss Si₃N₄ platform, in *Silicon Photonics XVII*, Vol. 12006, edited by G. T. Reed and A. P. Knights, International Society for Optics and Photonics (SPIE, Bellingham, Washington USA, 2022), p. 120060D.