

Deterministic feedforward-based generation of large optical coherent-state superposition

Michele N. Notarnicola^{1,*}, Marcin Jarzyna^{2,†} and Radim Filip^{1,‡}

¹Department of Optics, *Palacký University*, 17. Listopadu 12, 779 00 Olomouc, Czech Republic

²Centre for Quantum Optical Technologies, Centre of New Technologies, *University of Warsaw*, Banacha 2c, 02-097 Warszawa, Poland



(Received 12 January 2026; accepted 11 September 2026; published 19 October 2026)

Large optical coherent-state superpositions are essential to advance quantum sensing, quantum repeaters, and error-correction codes. We propose a deterministic feedforward protocol employing qubit-mode dispersive coupling, currently available in cavity quantum electrodynamics (QED). We show this single-mode protocol to outperform the advanced three-mode Gaussian-photon-number-resolving detector scheme in terms of both average fidelity and quantum non-Gaussian phase-space properties, and propose sensitivity to weak displacements of interference fringes as a feasible and conclusive witness of quantum interference. This approach combining QED with electro-optical feedforward is extendable to tailored states for applications and other platforms.

DOI: [10.1103/1944-5x1g](https://doi.org/10.1103/1944-5x1g)

I. INTRODUCTION

Generation of quantum superposition states is a challenging task in quantum information science as they provide resource for quantum metrology [1–3], computation [4–8], communication [9–13], and continuous-variable quantum error correction [14,15]. A special challenge arises for optical coherent-state superposition, hereafter referred to as *optical cat states*, namely, $|\psi_{\text{cat}}^{(\pm)}\rangle = (|\alpha\rangle \pm |-\alpha\rangle)/\mathcal{N}_{\pm}$, $\alpha \geq 0$, with $\mathcal{N}_{\pm} = \sqrt{2(1 \pm \exp(-2\alpha^2))}$ [16]. Conventional methods adopt probabilistic schemes based on either photon subtraction [17–20] or addition [21] that both exploit two- or three-mode Gaussian coupling between signal and ancillas followed by photon-number-resolving detectors (PNRDs), but, despite feasibility, are associated with two main limitations. First, they are efficient only for kitten generation, as they implement conditional projection over a finite superposition of Fock states embedded in a Gaussian envelope [22]. In turn, generation of cats with $\alpha \geq 2$ becomes challenging, requiring a large number of heralded photons to achieve high fidelity and, moreover, quantum non-Gaussian (QNG) phase-space properties, thus inducing a significant drop of the success rate [20]. Even if cat enlargement protocols have also been proposed, they come at the cost of further reduction of success probability [23], which can be partially compensated by high-quality quantum memories, whose realization is, however, currently challenging too. Second, QNG interference

properties of the resulting states are only partially reproducing the desired ones, as they only get finite stellar rank [24,25].

On the other hand, recent experimental progresses in optical cavity quantum electrodynamics (QED) have proved it as candidate for the task [27], with the advantage of deterministic generation and flexibility to implement different types of atom-light interaction [28–33]. Among all, dispersive coupling plays a promising role, as it can be effectively realized on traveling waves, thanks to reflection of an optical pulse from a cavity containing a single Λ -type atom in the limit of large cooperativity [28,30,34,35], thus providing a method directly competing with Gaussian coupling protocols. Moreover, given the existing limitations in cavity QED experiments, it is crucial to design feasible deterministic protocols that combine gates and operations already available with the state-of-the-art quantum optical technology and simultaneously offer scalable methods to extend the cavity QED capabilities, e.g., to target a wider class of QNG states for applications.

In this work, we propose a deterministic single-mode feedforward protocol based on dispersive coupling [Fig. 1(a)], obtained by embedding the single-atom cavity in Ref. [30] with Gaussian electro-optical feedforward, with the advantages to reach both high average fidelity and required QNG interference fringes in phase space. For a comparison, we briefly determine the performance achievable by Gaussian resources from nonlinear optics and highly efficient PNRDs by designing the heralded three-mode Gaussian-PNRD-based GP(2n) protocol [Fig. 1(b)], and show the dispersive scheme to outperform it for large cat generation in essential quantum features. The structure of the paper is the following. In Secs. II and III, we describe the proposed GP(2n) and dispersive protocols, and compare their performance in terms of fidelity and QNG features, respectively. Then, Sec. IV addresses the impact of imperfect dispersive coupling on the state preparation, while Sec. V presents displacement sensitivity as a suitable metric for cat-state quality. Finally, in Secs. VI and VII, we perform discussion about the obtained

*Contact author: michelenicola.notarnicola@upol.cz

†Contact author: m.jarzyna@cent.uw.edu.pl

‡Contact author: filip@optics.upol.cz

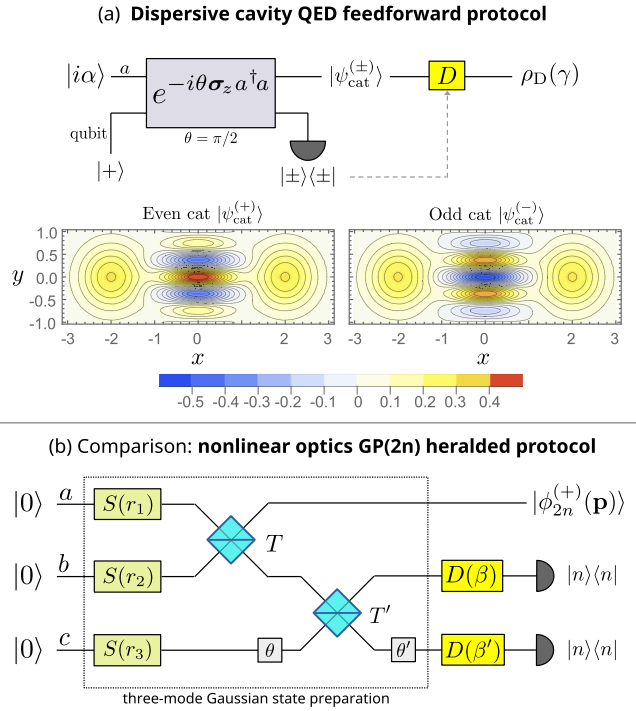


FIG. 1. (a) Dispersive coupling protocol for even-cat generation, based on qubit-light interaction and feedforward displacement conditioned on qubit measurement. The qubit measurement conditionally projects light to either the even or the odd cat state, whose Wigner functions $W_{\text{cat}}^{(\pm)}(x, y)$ are plotted in the panels above for $\alpha = 2$. If outcome “−” is retrieved, state $|\psi_{\text{cat}}^{(-)}\rangle$ is displaced along the y direction of phase space to make its quantum interference fringes overlap with those of the even cat, thus making the protocol deterministic. (b) Gaussian-PNRD-based [GP(2n)] protocol, performing conditional displacement-PNRD over an optimized three-mode Gaussian state. Extensions of this basic scheme, including displacements in state preparation and feedforward PNRD, are not helpful to improve quality of the output state [26].

results and draw some conclusions, respectively. Without loss of generality, we present the analysis by targeting the even cat $|\psi_{\text{cat}}^{(+)}\rangle$, whereas the odd-cat case is briefly discussed in Ref. [26].

II. THE GP(2n) AND DISPERSIVE PROTOCOLS

Upon generalization of photon addition and subtraction setups, we identify the scheme in Fig. 1(b) as the suitable configuration with the lowest number of modes for even-cat generation with Gaussian resources, hence referred to as the GP(2n) protocol. It is composed of an optimized Gaussian tritter, coupling vacuum modes a, b, c , decomposed by Bloch-Messiah theorem into a sequence of pure single-mode squeezers $S(r_j)$ and passive operations [36,37], after which we introduce displacements $D_b(\beta)D_c(\beta')$, recently proved as necessary for quantum state engineering [38], and projection by PNRDs onto Fock states $|n\rangle$, thus heralding $2n$ photons in total. To guarantee even parity of the output state, it is crucial to choose a balanced second beam splitter with π phase shifts, i.e., $T' = 1/2$, $\theta = \theta' = \pi$, and equal and opposite displacement amplitudes $\beta' = -\beta$, simplifying the tritter

structure to the set of free parameters $\mathbf{p} = (r_1, r_2, r_3, T, \beta)$, where we further assume $r_j \in \mathbb{R}$ as $\alpha \geq 0$. We also note that for the choices $(r_2 = r_3 = 0; \beta = 0)$ and $(r_2 = -r_1; r_3 = 0; T = 1/2)$, we retrieve the subtraction and addition protocols, respectively [26], justifying the GP(2n) as the more general lowest-mode Gaussian coupling scheme to achieve even parity. The $\mathbf{p}$ -dependent quantum state is obtained as $|\phi_{2n}^{(+)}(\mathbf{p})\rangle = S(\chi) \sum_{k=0}^n d_{2k} |2k\rangle$ (derived in Ref. [26]) that, by construction, gets finite stellar rank $\leq 2n$, limited by maximal errorless n -photon resolution of the PRNDs. Then, we optimize $\mathbf{p}$ to maximize fidelity with respect to the target state, namely

$$F_{\text{GP}}^{(2n)} = \max_{\mathbf{p}} |\langle \psi_{\text{cat}}^{(+)} | \phi_{2n}^{(+)}(\mathbf{p}) \rangle|^2. \quad (1)$$

We also discovered that, for the present case study, maximizing fidelity is equivalent to maximizing negativity of the GP(2n) Wigner function $W_{\text{GP}}^{(2n)}(0, \bar{y})$ at the point $\bar{y}$, where the target Wigner function $W_{\text{cat}}^{(+)}$ takes its absolute minimum [39].

To overcome the heralded GP(2n) scheme, we propose the atom-light cavity QED protocol depicted in Fig. 1(a), where nonlinearity comes in the interaction and not only at detection stage. It extends the seminal idea first demonstrated by Hacker *et al.* in Ref. [30], where a qubit and optical mode interact via dispersive Hamiltonian $H_D = \theta \sigma_z a^\dagger a$, with $\theta \geq 0$ and $\sigma_z = |g\rangle\langle g| - |e\rangle\langle e|$ being the qubit Pauli z operator [40], such that the unitary $U_D = e^{-iH_D}$ implements a conditional rotation of a in either clockwise or anticlockwise direction according to the qubit state. Initially, the dispersive protocol applies a $\pi/2$ pulse, namely, $\theta = \pi/2$, onto coherent state of radiation $|i\alpha\rangle$, $\alpha \geq 0$, and atomic state $|+\rangle = (|g\rangle + |e\rangle)/\sqrt{2}$, to create the “atom-light entangled state” $|\Psi_D\rangle = (|g\rangle|\alpha\rangle + |e\rangle|-\alpha\rangle)/\sqrt{2}$ [30,35], with subsequent qubit measurement in the $\{|\pm\rangle\}$ basis, such that, if “+” is retrieved, light is projected exactly into the target state $|\psi_{\text{cat}}^{(+)}\rangle$, whereas in the opposite case we obtain the orthogonal odd cat.

Now, to make the protocol deterministic, a suitable strategy to “turn” the odd cat into the even one would be desirable. The problem has been addressed in circuit QED with different methods: (1) applying a subsequent conditional displacement gate $D(i\sigma_y\beta)$ that approximately disentangles the qubit and light degrees of freedom in state $|\Psi_D\rangle$ [41,42]; (2) inducing a nonadiabatic two-photon driving, available in Kerr nonlinear resonators [43,44]; and (3) designing reservoir engineering schemes that combine a classical [45] or two-photon [46,47] driving with a suitable two-photon dissipation process. We also remind that the Wigner function of a cat state exhibits a quantum interference pattern with Gaussian-shaped envelope between the two coherent-state peaks, with maximum (minimum) in the origin for the even (odd) state [see Fig. 1(a)]. Therefore, upon drawing inspiration on the displacement method in Ref. [45], we can make our cavity QED protocol deterministic by using quantum optics feedforward techniques [48–51]. That is, when outcome “−” is registered, we apply the feedforward displacement operation $D(\pm i\gamma)$, $\gamma \geq 0$, on the optical state, whose sign is randomly chosen with 50% probability, inducing either positive or negative shift of the interference fringes along the y direction of phase space. In this way, we displace the odd cat to let its fringes overlap with

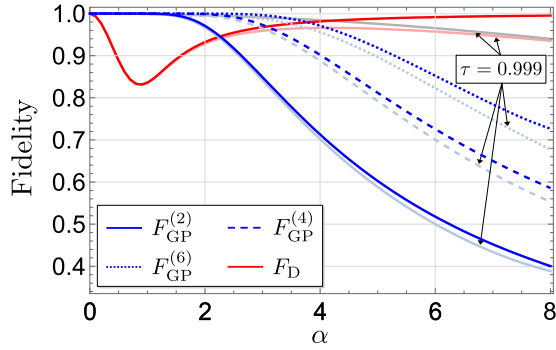


FIG. 2. Fidelities $F_{\text{GP}}^{(2n)}$ for $n = 1, 2, 3$ and F_D as a function of the target cat amplitude α . Dispersive interaction provides a deterministic scheme that outperforms $\text{GP}(2n)$ in the limit of large α ; in Ref. [26], we discuss a probabilistic version of the protocol beating $\text{GP}(2n)$ also for small α . Moreover, the $\text{GP}(2n)$ success probability for $\alpha \gg 1$ approaches $\approx 7.9\%$, 3.0% , and 1.6% for $n = 1, 2, 3$ [26]. The light-colored lines represent fidelities after optical losses with transmissivity $\tau \leq 1$, while the gray one refers to attenuation of the target cat $|\psi_{\text{cat}}^{(+)}\rangle$.

those of the even state, and eventually obtain the mixture:

$$\rho_D(\gamma) = p_+ |\psi_{\text{cat}}^{(+)}\rangle \langle \psi_{\text{cat}}^{(+)}| + p_- \left(\frac{|\psi'(\gamma)\rangle \langle \psi'(\gamma)| + |\psi'(-\gamma)\rangle \langle \psi'(-\gamma)|}{2} \right), \quad (2)$$

with $|\psi'(\gamma)\rangle = D(i\gamma)|\psi_{\text{cat}}^{(-)}\rangle$ and $p_{\pm} = (1 \pm e^{-2\alpha^2})/2$ being the probability to detect the atom in the $|\pm\rangle$ state [26]. The optimal displacement γ is then chosen to maximize fidelity:

$$F_D = \max_{\gamma} \langle \psi_{\text{cat}}^{(+)} | \rho_D(\gamma) | \psi_{\text{cat}}^{(+)} \rangle. \quad (3)$$

Following the same approach, the dispersive protocol can be adopted to deterministically prepare the odd cat $|\psi_{\text{cat}}^{(-)}\rangle$ with analogous performance to the even-cat case (see the Supplemental Material [26]).

III. RESULTS

Plots of fidelity and cuts of the Wigner function along the y axis for both the $\text{GP}(2n)$ protocol with $n = 1, 2, 3$ and the dispersive scheme are reported in Figs. 2 and 3, respectively. Due to the finite stellar rank, the $\text{GP}(2n)$ performs well only in the kitten regime, while, for large enough α , $F_{\text{GP}}^{(2n)}$ becomes a decreasing function, and the interference pattern of the Wigner function exhibits only $2n$ negative dips, thus degrading the quality of the produced states for large α . In this regime, the success probability reaches $\approx 7.9\%$, 3.0% , and 1.6% for $n = 1, 2, 3$ [26]. On the contrary, the dispersive protocol achieves much better performance, as for $\alpha \gtrsim 0.8$, F_D becomes an increasing function that asymptotically approaches 1 as

$$F_D \approx 1 - \frac{\pi^2}{32\alpha^2} \quad \text{for } \alpha \gg 1, \quad (4)$$

where fidelity rises as a consequence of the feedforward displacement. Indeed, displacement of the odd-cat fringes makes states $D(\pm i\gamma)|\psi_{\text{cat}}^{(-)}\rangle$ and $|\psi_{\text{cat}}^{(+)}\rangle$ overlap, but worsens the correspondence between both their Gaussian envelopes and coherent-state peaks, so that, as the fringe spacing scales as $1/\alpha$ for $\alpha \gg 1$, the infidelity decreases as $1/\alpha^2$. Same asymptotic trend holds for the odd-cat generation [26]. Equation (4) provides comparable result to the protocol for microwave cat generation proposed in Refs. [41,42] (whose fidelity scales as $\approx 1 - \pi^2/64\alpha^2$), which, however, requires sequential iterations of dispersive couplings, whereas our proposal only needs a single light-atom interaction step, being assisted by Gaussian feedforward. Besides, we observe many oscillations of the corresponding Wigner function, jointly appearing after

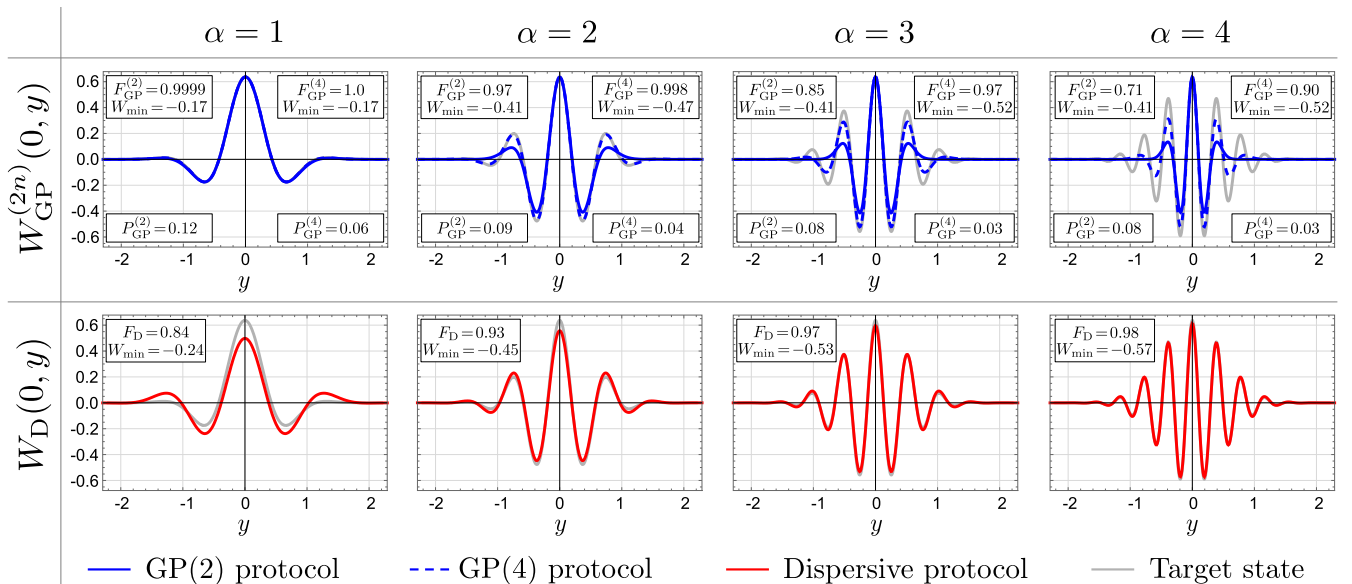


FIG. 3. Oscillations of the Wigner function along the y axis for the $\text{GP}(2n)$ protocol, with $n = 1, 2$ (top row), and the dispersive protocol (bottom row) for different values of the cat amplitude α . Boxes report the corresponding values of fidelity and Wigner-function minima.

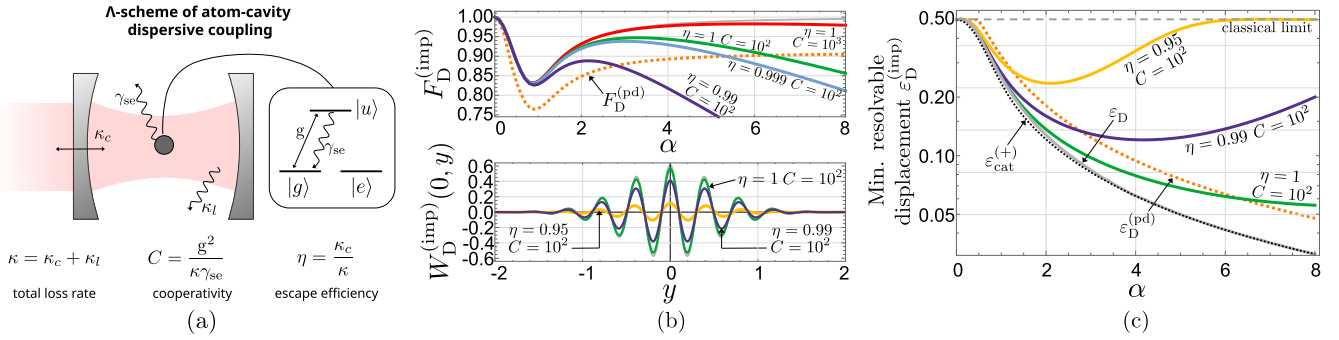


FIG. 4. (a) Cavity QED implementation of dispersive interaction, consisting of a cavity containing a Λ -type atom, coupled to impinging free-space field with rate κ_c and to scattering and loss modes with rate κ_l . The two atomic levels $|g\rangle$ and $|u\rangle$ are resonantly coupled to the cavity field with gain g , and transition $|g\rangle \rightarrow |u\rangle$ is also associated with spontaneous decay rate γ_{se} [26,30,34,35]. The system imperfections are conveniently described in terms of cooperativity C and escape efficiency $\eta \leq 1$: Ideal dispersive coupling is realized in the limits $C \gg 1$ and $\eta = 1$ (lossless cavity). (b) Fidelities $F_D^{(imp)}$ (top) of the imperfect dispersive coupling protocol and Wigner interference fringes $W_D^{(imp)}(0, y)$ with $\alpha = 4$ (bottom) for different values of C and η . The gray line corresponds to the ideal dispersive scheme; the dashed line refers to fidelity $F_D^{(pd)}$ in the presence of qubit phase damping of rate $\lambda = 0.2$ [26]. (c) Minimum resolvable displacements for ideal and imperfect dispersive schemes ε_D and $\varepsilon_D^{(imp)}$ as a function of α : The ideal protocol approaches the target state sensitivity $\varepsilon_{cat}^{(+)}$ for $\alpha \gg 1$. The dashed line refers to $\varepsilon_D^{(pd)}$ in the case of qubit phase damping of rate $\lambda = 0.2$ [26].

only one interaction step, with much better accordance with the target state; the minimum value of the Wigner function is equal to

$$W_D(0, \bar{y}) \approx W_{cat}^{(+)}(0, \bar{y}) + \frac{\pi}{8\alpha^2} \quad \text{for } \alpha \gg 1; \quad (5)$$

and associated distillable squeezing variance $V_D \approx V_{cat}^{(+)} - 1/8\alpha^4$ [26,52], demonstrating dispersive interaction as the most powerful for large cat generation. In fact, while the GP($2n$) conditionally projects a Gaussian state onto a finite superposition of Fock states of rank $\leq 2n$, the dispersive protocol exploits light-atom interaction to translate superposition created in the qubit to the optical mode, and experimental new part-Gaussian feedforward to achieve determinism, thus combining cavity QED experiments [28,30–35] with well-established electro-optical feedforward methods [48–51].

Figure 2 also shows robustness of the produced output states undergoing optical losses, modeled as a beam splitter of transmissivity $\tau \leq 1$, in which case the dispersive fidelity $F_D(\tau)$ loses the asymptotic behavior (4), becoming a decreasing function that for $\alpha \gg 1$ scales as

$$F_D(\tau) \approx e^{-\alpha^2(1-\sqrt{\tau})^2} \frac{1 + e^{-2\alpha^2(1-\tau)}}{2} \left(1 - \frac{\pi^2}{32\alpha^2}\right), \quad (6)$$

while Wigner negativities are exponentially reduced [26].

IV. ROBUSTNESS TO IMPERFECT DISPERSIVE COUPLING

The dispersive interaction required for the proposed protocol is implemented by the cavity QED setup in Fig. 4(a). To understand the difference to the GP($2n$) scheme (mainly limited by PNRDs), we investigate its robustness to imperfect coupling realization, considering realistic conditions of finite cooperativity $C < \infty$ and nonunit escape efficiency $1/2 < \eta < 1$ [30,34,53]. Both effects can be modeled as conditional optical losses with qubit-state-dependent transmissivity $\eta_{g(e)} \leq 1$, such that the ideally conditionally prepared states

$|\psi_{cat}^{(\pm)}\rangle$ decohere into the mixtures [26]:

$$\sigma_{\pm}^{(imp)} = \frac{1}{4p_{\pm}^{(imp)}} \{ |\sqrt{\eta_g}\alpha\rangle \langle \sqrt{\eta_g}\alpha| + |-\sqrt{\eta_e}\alpha\rangle \langle -\sqrt{\eta_e}\alpha| \pm e^{-\Gamma\alpha^2/2} (|\sqrt{\eta_g}\alpha\rangle \langle -\sqrt{\eta_e}\alpha| + \text{H.c.}) \}, \quad (7)$$

with $\eta_g = [1 - 2\eta/(1 + 4C)]^2$, $\eta_e = (1 - 2\eta)^2$, $\Gamma = 2 - \eta_g - \eta_e + 2\sqrt{(1 - \eta_e)(\eta' - \eta_g)}$, $\eta' = 1 - 16\eta C/(1 + 4C)^2$, and detection probabilities $p_{\pm}^{(imp)} = (1 \pm e^{-[\Gamma + (\sqrt{\eta_g} + \sqrt{\eta_e})^2]\alpha^2/2})/2$. The average output state becomes $\rho_D^{(imp)}(\gamma) = p_+^{(imp)}\sigma_+^{(imp)} + p_-^{(imp)}\sum_{k=\pm} D(k i\gamma)\sigma_-^{(imp)}D^\dagger(k i\gamma)/2$, with fidelity $F_D^{(imp)} = \max_{\gamma} \langle \psi_{cat}^{(+)} | \rho_D^{(imp)}(\gamma) | \psi_{cat}^{(+)} \rangle$. As depicted in Fig. 4(b), for large α , $F_D^{(imp)}$ deviates gradually from the ideal case and becomes a decreasing function, thus expectedly losing the asymptotic trend, such that for $\alpha, C \gg 1$ and $1 - \eta \ll 1$,

$$F_D^{(imp)} \approx \frac{1 + e^{-2\alpha^2(\frac{2-\eta}{4C} + 1 - \eta)}}{2} \left(1 - \frac{\pi^2}{32\alpha^2}\right) - \frac{\pi^2(1 - \eta)}{64C} e^{-2\alpha^2(\frac{2-\eta}{4C} + 1 - \eta)} \left(1 - \frac{\pi^2}{16\alpha^2}\right). \quad (8)$$

Instead, Wigner negativities, which witness QNG features, are not erased for any α , albeit we observe α -dependent suppression of oscillations, such that for $\alpha, C \gg 1$ and $1 - \eta \ll 1$,

$$W_D^{(imp)}(0, \bar{y}) \approx W_{cat}^{(+)}(0, \bar{y}) + \frac{f_1(\alpha)}{2\pi} + \frac{\pi f_2(\alpha)}{8\alpha^2}, \quad (9)$$

with $f_1(\alpha) = 4 - e^{-2\alpha^2(1-\eta+\eta/4C)} - 3e^{-\alpha^2/2C}(1 - \pi^2(1 - \eta)/2C)$ and $f_2(\alpha) = -2 + 3e^{-\alpha^2/2C}(1 - 3\pi^2(1 - \eta)/4C)$, while the distillable squeezing variance increases with α , in faster way than negativity reduction [26].

With similar approach, in Ref. [26] we further address the impact of decoherence before and inside qubit detection, by considering amplitude and phase damping noise, which, as re-

ported in Fig. 4(b), induce only fidelity rescaling in the asymptotic limit and reduction of Wigner interference visibility proportional to the sole noise strength and independent of α .

V. DISPLACEMENT SENSITIVITY AS INTERFERENCE MEASURE

To further understand quality of the interference fringes, we test sensitivity of the previous quantum states to momentum displacement [54,55]. The problem can be recast in the framework of quantum estimation theory [56–60], where a weak displacement $D(i\varepsilon) = \exp(\sqrt{2}i\varepsilon\hat{x})$, with $\varepsilon \geq 0$ and $\hat{x} = (a + a^\dagger)/\sqrt{2}$ being the positionlike operator, is applied to a probe state. Considering homodyne detection of quadrature $\hat{y} = i(a^\dagger - a)/\sqrt{2}$, the Cramér-Rao bound determines the minimum resolvable displacement $\varepsilon_{\min} = 1/\sqrt{\mathbb{F}}$ per experimental run, where

$$\mathbb{F} = \int_{\mathbb{R}} dy \frac{1}{p(y + \sqrt{2}\varepsilon)} \left(\frac{d p(y + \sqrt{2}\varepsilon)}{d\varepsilon} \right)^2 \quad (10)$$

is the Fisher information (FI) associated with the momentumlike probability distribution $p(y)$ [26], which for cat-state probes beats the classical limit, obtained with coherent states and corresponding to vacuum fluctuations, $\varepsilon_{\min} = 1/2$. Figure 4(c) shows the $\varepsilon_{\min}$ numerically computed for both the ideal and imperfect dispersive coupling protocols, proving the ideal coupling case to approach the target cat for $\alpha \gtrsim 3$. Instead, imperfect coupling becomes detrimental in the large-amplitude regime, where $\varepsilon_D^{(\text{imp})}$ increases with α and eventually equals the classical limit, whereas for $\alpha \ll 1$ we have $\varepsilon_D^{(\text{imp})} \approx \varepsilon_D$. Therefore, error mitigation schemes [61,62] turn out to be essential also for high-quality deterministic state preparation. On the contrary, in the presence of qubit detection decoherence, $\varepsilon_{\min}$ remains a decreasing function for $\alpha \gg 1$, thus maintaining the quantum advantage. Analogous results are also obtained for the quantum FI, which provides the optimal quantum-limited sensitivity (see the Supplemental Material [26]).

VI. DISCUSSION

The above results qualify feedforward operations as a powerful strategy to improve the performance of cavity QED experiments, for twofold reason. First, while conditional operations based on optical homodyne detection are common [63–65], hybrid atom-light feedforward in cavity QED has been only preliminary demonstrated in Ref. [66] to implement a photon-photon gate in a low-photon regime, and its potential has not been fully unveiled. Our results prove atom-light feedforward as crucial to achieve determinism in large cat-state preparation, and provide a cornerstone toward further extension of the cavity QED capabilities. For instance, by sequential application of the proposed dispersive gate to two subsequent coherent pulses $|\alpha\rangle|\beta\rangle$, followed by atom detection in the $|\pm\rangle$ basis and double feedforward displacement, one can target the nonlocal cat states $|\alpha\rangle|\beta\rangle \pm |-\alpha\rangle|-\beta\rangle$ and achieve deterministic entanglement generation, being resource for quantum error correction [15].

Second, adopting electro-optical feedforward may simplify deterministic atom-light coupling protocols. In fact, the ex-

isting schemes for cat preparation require two subsequent conditional displacement gates $D(\sigma_x\alpha)D(i\sigma_y\beta)$ [67,68] that require multiple reflections of the optical beam to the cavity and precise control of the experimental setup in the realistic conditions of finite cooperativity and nonunit escape efficiency. In contrast, our protocol replaces gate $D(i\sigma_y\beta)$ with Gaussian feedforward to reduce the setup complexity to one single cavity reflection, obtaining comparable performance to circuit QED schemes [41,42]. Together with setup simplification, this choice also paves the way for distributed quantum experiments, where quantum gates can be obtained via nonlocal interaction of two remote subsystems [69]. While the current demonstrations of nonlocal gates have been only carried out in discrete-variable platforms [70–72], our results offer a powerful strategy to extend this methodology to hybrid platforms, with the near-future objective of remote hybrid quantum networks.

VII. CONCLUSIONS

In this work, we proposed a deterministic feedforward-based protocol using dispersive interaction with a qubit for cat-state generation, proving it to outperform the Gaussian-PNRD-based protocol for $\alpha \gg 1$. We demonstrate robustness of its QNG features against imperfect coupling and qubit decoherence, showing dispersive coupling as promising for the following reasons: determinism, large fidelity, high-visibility quantum interference obtained after a single interaction step, and flexible implementation by cavity reflection of coherent traveling beams. Our results also prove FI, already proposed as a quantum macroscopicity measure [73–78], as a reliable witness of cat-state interference. In fact, while fidelity and Wigner negativity are exponentially decreasing for $\alpha \gg 1$ in the presence of noise, the minimum resolvable displacement provides a more robust quantifier of large cat-state quality and reveals its nonclassicality also in the suboptimal, but feasible, case of homodyne detection. Moreover, our scheme can be straightforwardly extended to target a wider class of superposition states, e.g., (1) unbalanced cat states $|\psi\rangle \propto c_1|\alpha\rangle + c_2|-\alpha\rangle$ relevant for phase sensing, by preparing the initial qubit state into the superposition $c_1|g\rangle + c_2|e\rangle$; (2) squeezed cat states $S(r)|\psi_{\text{cat}}^{(\pm)}\rangle$, $r \in \mathbb{R}$, obtained with the same fidelity depicted in Fig. 2 by considering an input squeezed coherent state $S(-r)|i\alpha\rangle$ and the feedforward Gaussian operation $S(r)D(\gamma)S(-r) = D(e^r\gamma)$; (3) multiheaded cat states by subsequent beam reflection from array cavities; and (4) multimode and nonlocal cats. The proposed protocol can also be translated to other platforms, including microwaves [79–83], atomic ensembles [84,85], and mechanical modes [86–89], thus providing a crucial step to target exotic states of radiation for quantum information purposes.

ACKNOWLEDGMENTS

The authors thank J. Fiurášek for stimulating discussion about the GP(2n) scheme. M.N.N. acknowledges the QuantERA Project No. 8C24003 (CLUSSTAR) of the Ministry of Education, Youth and Sport (MEYS) of the Czech Republic. Project CLUSSTAR has received funding from the EU Horizon Programme under Grant Agreements No. 731473 and No. 101017733 (QuantERA). R.F. acknowledges the project 23-06308S of the Czech Science Foundation, the EU Horizon

Programme under Grant Agreements No. 101080173 (CLUSTEC) and No. 101257526 (SUPERSPIN) and the project CZ.02.01.010022_0080004649 (QUEENTEC) of the EU and the Czech MEYS.

DATA AVAILABILITY

The data generation codes supporting the findings of this manuscript are openly available on Zenodo [107].

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