

Maximum-precision charging of multiqubit quantum batteriesDavide Rinaldi ^{1,*}, Radim Filip ², Dario Gerace ¹ and Giacomo Guarnieri ^{1,3}¹*Dipartimento di Fisica “A. Volta,” Università di Pavia, via Bassi 6, 27100 Pavia, Italy*²*Department of Optics, Faculty of Science, Palacký University, 17. listopadu 12, 77900 Olomouc, Czech Republic*³*INFN Sezione di Pavia, Via Agostino Bassi 6, I-27100 Pavia, Italy*

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Precision, robustness, and efficiency are central requirements for quantum technologies. We show that genuine quantum features combined with non-Gaussianity enable the simultaneous optimization of these properties in a quantum battery-charging process. Using a generalized Jaynes-Cummings interaction as a paradigmatic light-matter interaction model, we apply the Full Counting Statistics to characterize stochastic energy exchanges between a stack of qubits and a single-mode bosonic field. We demonstrate that a sequential charging protocol driven by a non-Gaussian quantum field yields high performance in charging precision, which remains maximal even under suboptimal operating conditions. Our results establish the use of non-Gaussian quantum states in battery charging as a robust route to a quantum precision advantage over protocols based on Gaussian states, achieved through the suppression of detrimental quantum fluctuations.

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The role of energy in Physics is related to the capability of a system to perform work. An energy-based perspective can be effectively adopted to study quantum devices, ranging from quantum information technologies [1–5] to thermal engines and refrigerators [6–13]. Recently, this viewpoint has fostered the emergence of a research field focused on the analysis of energy exchanges between quantum systems. The models developed in this context, where energy plays a central role, fall under the category of so-called quantum batteries [14–62]. These models are so defined [16] due to their classification of energy exchanges into three main processes: energy injection, energy storage, and on-demand energy delivery. This field is currently in the spotlight [14,15]: numerous theoretical works on quantum batteries, aimed at exploring the fundamental limits of quantum thermodynamics [28–42] and designing increasingly efficient and noise-resilient devices [43–53], are accompanied by experimental realizations [54] on superconducting apparatuses [55–58], trapped-ion platforms [59], quantum dots [60], and nuclear-spin systems [61].

However, dealing with energy exchanges at a quantum level means that energy fluctuations cannot be neglected: their magnitude can assume values comparable to the average energy exchanges themselves. A simple (and yet effective) method to quantify energy fluctuations is the Full Counting Statistics [63–66], with which process-dependent thermodynamical quantities can be properly examined. We employ such a technique to quantify the charging process in a Jaynes-Cummings [67–70] quantum battery, i.e., a battery

model consisting of a qubit (i.e., the battery) coupled to a single-mode harmonic oscillating field, typically confined in a resonating cavity (i.e., the charger), be it electromagnetic or mechanical.

In this Letter, we leverage on the concepts and tools of Quantum Thermodynamics to actually implement realistic quantum batteries with proven advantages in maximizing the charging precision. We specifically demonstrate that a multiqubit quantum battery can be operated at a low level of quantum fluctuations, crucial to envision a reliable quantum device, when a genuine quantum non-Gaussian state is employed as a charger. Such a quantum battery can then be employed to perform fine precision tasks, such as quantum state preparation [71–73] or energy flow monitoring in quantum devices [74–76].

Model and methods. The paradigmatic model for our quantum battery is based on the Jaynes-Cummings (JC) interaction [67–70]. Drawing inspiration from microwave cavity QED experiments [77], this model describes a qubit with a resonance frequency ω_{qub} , coupled to a single mode of a cavity harmonic oscillator (such as an electromagnetic or mechanical field operator), characterized by its frequency ω_{cav} , as depicted in Fig. 1. The qubit-cavity coupling is governed by the interaction strength g , which is assumed to remain constant over time. In the broadly available weak-coupling regime, where $g/\omega_0 \ll 1$, with $\omega_{\text{qub}} \simeq \omega_0 \simeq \omega_{\text{cav}}$, the counterrotating terms in the Rabi Hamiltonian can be neglected by applying the rotating-wave approximation (RWA) [67]: we quantify the error made by performing the RWA in the Supplemental Material (SM) [78]. To extend this framework to M noninteracting qubits coupled to the same harmonic oscillator, we introduce the following general Hamiltonian:

$$\hat{H}(t) = \hat{H}_{\text{free}} + \sum_{j=1}^M \hbar g A_j(t) (\hat{\sigma}_+^{(j)} \hat{a} + \hat{\sigma}_-^{(j)} \hat{a}^\dagger). \quad (1)$$

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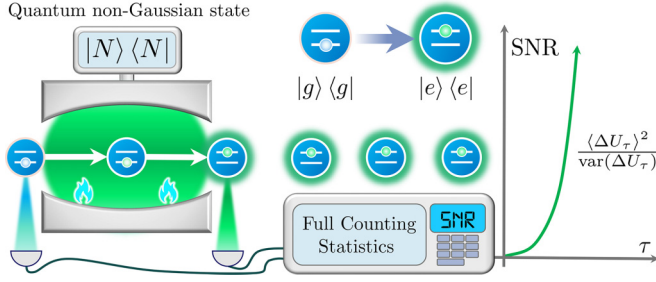


FIG. 1. Multiqubit JC quantum battery for precision-focused qubit charging with a Fock state cavity. The achieved charging advantage, quantified by the process-dependent SNR via the FCS, is both quantum non-Gaussian and operational, and it remains remarkably robust even beyond the ideal JC model and in realistic scenarios involving noise, dissipation, and parameter-tuning errors.

When the factor $A_j(t) = 0 \ \forall j \neq 1$ and $A_1(t) = 1 \ \forall t$, we recover the standard JC model [67,79–84], which can be analytically solved by computing the evolution operator $\hat{U}(t, 0) = \mathcal{T} e^{-\frac{i}{\hbar} \int_0^t \hat{H}(t') dt'}$ [67–70]. Here, $\hat{H}_{\text{free}} = \sum_{j=1}^M \hat{H}_{\text{qub}}^{(j)} + \hat{H}_{\text{cav}}$, $\hat{H}_{\text{qub}}^{(j)} = \hbar \omega_{\text{qub}} \frac{1}{2} \hat{\sigma}_z^{(j)}$, and $\hat{H}_{\text{cav}} = \hbar \omega_{\text{cav}} \hat{a}^\dagger \hat{a}$, in which $\hat{\sigma}_{x,y,z}^{(j)}$ denote the standard Pauli matrices for the j th qubit, $\hat{\sigma}_\pm^{(j)} = \frac{1}{2}(\hat{\sigma}_x^{(j)} \pm i\hat{\sigma}_y^{(j)})$, and $\hat{a}$ ($\hat{a}^\dagger$) is the annihilation (creation) operator of a single excitation in the harmonic single-mode field. The j th qubit state basis $\{|e\rangle_j, |g\rangle_j\}$, eigenvectors of $\hat{\sigma}_z^{(j)}$, defines the pure excited and ground state, respectively. We define the j th qubit state at time t as $\hat{\rho}_{\text{qub}}^{(j)}(t)$, and the cavity mode state as $\hat{\rho}_{\text{cav}}(t)$; in addition, at the initial time $t = 0$, the qubits and cavity mode states are assumed to be factorized: $\hat{\rho}_{\text{total}}(0) = \bigotimes_{j=1}^M \hat{\rho}_{\text{qub}}^{(j)}(0) \otimes \hat{\rho}_{\text{cav}}(0)$. The cavity field state is then initialized [85] in either a N -photon Fock state, $|N\rangle\langle N|$, or a coherent state $|\alpha\rangle\langle\alpha|$, for comparison, in which $|\alpha|^2$ is the average number of photons, or also in a quantum Gaussian state, e.g., a squeezed coherent state $|\tilde{\alpha}, \zeta\rangle\langle\tilde{\alpha}, \zeta|$, with ζ being the squeezing parameter [86]. The charging process consists of using the Hamiltonian interaction (1) with the charger to induce ground-to-excited state transitions in each qubit. The goal is to do so with the maximum precision (i.e., lowest variance) possible on the energy exchanged between the cavity and the qubit during a time window τ .

With the aim of quantitatively assessing the behavior of this energy exchange, and thus characterizing its fluctuations over the interaction time, we focus on the stochastic variable ΔU_τ [87]. We refer to this variable as the “qubit internal energy,” as it is directly related to the measurement outcomes of a specific observable: the Hamiltonian $\hat{H}_{\text{qub}} = \hbar \omega_{\text{qub}} \frac{1}{2} \hat{\sigma}_z$, evaluated at two distinct times. Here, we use theoretical techniques falling within the Full Counting Statistics (FCS) [63]. Within FCS, the statistics of ΔU_τ are fully described by a qubit-specific generating function $\mathcal{G}_\tau(\chi)$, such that $\langle \Delta U_\tau^n \rangle = (-i)^n \frac{\partial^n \mathcal{G}_\tau(\chi)}{\partial \chi^n} \big|_{\chi=0}$. $\mathcal{G}_\tau(\chi)$ can be experimentally measured [88] and compared to its theoretically derived form, whenever the underlying dynamical model is known. Specifically, the generating function is computed as $\mathcal{G}_\tau(\chi) = \text{Tr}\{\hat{\rho}_\chi(t)\}$, where $\hat{\rho}_\chi(t) = \hat{U}_\chi(\tau, 0) \hat{\rho}_{\text{JC}}(0) \hat{U}_\chi^\dagger(\tau, 0)$ and $\hat{U}_\chi(\tau, 0) = e^{i\frac{\chi}{2} \hat{H}_{\text{qub}}} \hat{U}(\tau, 0) e^{-i\frac{\chi}{2} \hat{H}_{\text{qub}}}$ [63,64]. The central

figure of merit for our discussion is the Signal-to-Noise Ratio (SNR), which can be extracted from the statistics of ΔU_τ , i.e., formally defined as

$$\text{SNR}(\Delta U_\tau) = \frac{\langle \Delta U_\tau \rangle^2}{\text{var}(\Delta U_\tau)}, \quad (2)$$

where $\text{var}(\Delta U_\tau) = \langle \Delta U_\tau^2 \rangle - \langle \Delta U_\tau \rangle^2$. As it is evident from Eq. (2), the SNR quantifies the signal amplitude (i.e., the mean exchanged energy) as compared to the magnitude of its fluctuation (i.e., its variance), and it therefore constitutes the best measure of the process precision. Within FCS, these *process*-dependent quantities allow to better capture the complexity of the fluctuating energy exchanges. This represents a remarkable conceptual leap from standard techniques, which employ *state*-dependent figures of merit focused only on the system state at a given time, although the latter have been applied in other thorough analyses, including Refs. [43,46,53]. For a comparison, we then compute the fidelity between the qubit state, $\hat{\rho}_{\text{qub}}(t)$, and the target state, $\hat{\rho}_{\text{target}} = |e\rangle\langle e|$, defined as

$$F(t) = \text{Tr}\left\{\left[\hat{\rho}_{\text{qub}}^{1/2}(t) \hat{\rho}_{\text{target}} \hat{\rho}_{\text{qub}}^{1/2}(t)\right]^{1/2}\right\}^2. \quad (3)$$

In the following, we show two different protocols allowing to charge more than a single qubit at a time, by using the harmonic oscillator state as a charger, with the aim of achieving the maximal charging precision: (i) sequential and (ii) parallel protocols, respectively.

Sequential charging protocol. Let us consider Eq. (1) with $A_j(t) = 1$ if $t \in [t_{j-1}, t_j]$, and $A_j(t) = 0$ otherwise. Now, the j th qubit interacts with the cavity during a time window $\tau_j = t_j - t_{j-1}$ (in general, $\tau_j \neq \tau_k$ for $j \neq k$, and for convenience $t_0 \equiv 0$). After the interaction, the i th qubit is immediately decoupled from the cavity, and at the same time the $(j+1)$ th qubit is coupled to the oscillator, analogously to the mechanism that occurs within a collisional-model description [89].

Due to the peculiarity of the model in Eq. (1), the Hamiltonians of different qubits commute with each other. This allows us to apply the same formalism of a preparatory work [62], where the elementary case of a single qubit was discussed, to derive the SNR of the exchanged energy in a modular fashion, i.e., within each time window τ_j . Importantly, we now keep track of the cavity state after each interaction: the j th charging round begins with the j th qubit in a fresh ground state, while the cavity state is modified to account for the qubit measurement-induced backaction [78] after the interaction with the $(j-1)$ th qubit.

As we show in Fig. 2(a), the Fock state SNR on the energy ΔU_τ injected in each qubit can be arbitrarily increased, thus being higher than the SNR obtained via coherent or Gaussian states, provided that the model parameters can be tuned to achieve the desired working conditions, such as perfect resonance and the absence of external disturbances. For that reason, in an idealized situation it is possible to perfectly charge M qubits by initializing the cavity in a N -photon Fock state with $N = M$, a result confirmed by the fact that the fidelity between the j th qubit states $\hat{\rho}_{\text{qub}}^{(j)}(t)$ and $\hat{\rho}_{\text{target}} = |e\rangle\langle e|$ can always reach $F = 1$ for each charging round, contrary to what happens if the cavity is prepared in a classical or a Gaussian state. Furthermore, other realistic scenarios will

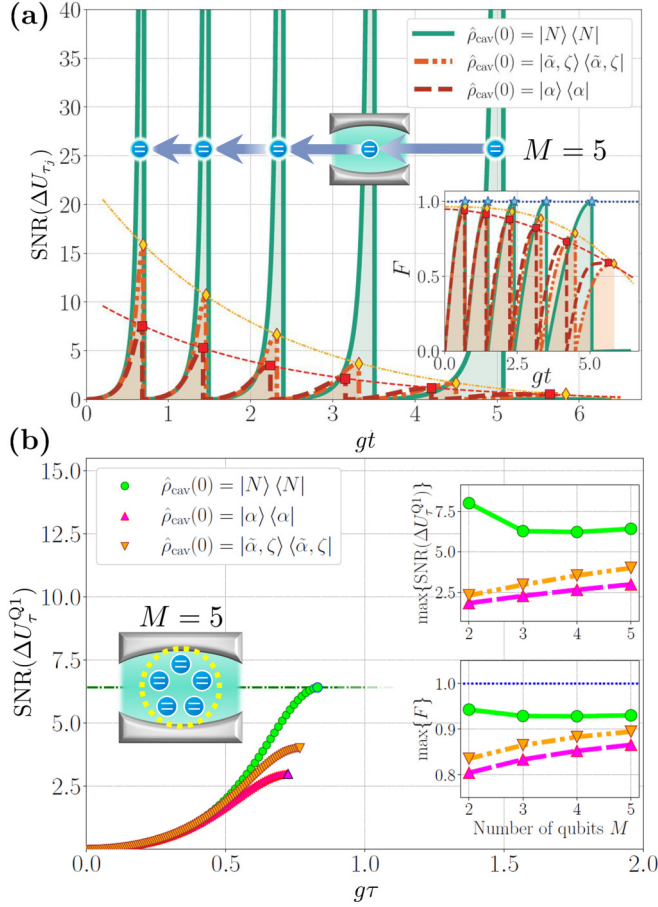


FIG. 2. (a) Sequential charging of $M = 5$ qubits, where $\text{SNR}(\Delta U_\tau)$ is calculated for each qubit charge. The fidelity $F = F(t)$ calculated between $\hat{\rho}_{\text{qub}}^{(j)}(t)$ and $\hat{\rho}_{\text{target}}^{(j)} = |e\rangle_j \langle e|_j$ is depicted in the inset. While in the presence of idealized conditions and perfect tuning of the interaction time τ the Fock state SNR peaks can be arbitrarily high, the ones related to the other states decrease exponentially in time. The fidelity reaches $F = 1$ only for the Fock state protocol; for the others, $F < 1$ and the peaks decrease in time as $F \propto -(gt)^\gamma$, with $\gamma > 0$. (b) Parallel charging of $M = 5$ qubits. The $\text{SNR}(\Delta U_\tau^{Q1})$ is related to the sole qubit 1, and it is obtained by numerically solving the TC model. Upper inset: maximum SNR for $M \in \{2, 3, 4, 5\}$ qubits. Lower inset: maximum fidelity F between $\hat{\rho}_{\text{qub}}^{(1)}(t)$ and $\hat{\rho}_{\text{target}}^{(1)} = |e\rangle_1 \langle e|_1$, varying the number of qubits. Simulation parameters: $\langle n \rangle = N = 5$, $\alpha = \sqrt{5}$ for the coherent state, $\zeta = 0.6$ and $\tilde{\alpha} \simeq 3.905$ for the squeezed coherent state, $g/\omega_{\text{qub}} = 10^{-2}$, and $\Delta\omega = 0$.

be discussed in what follows, showing that the quantum advantage persists even if the variance does not completely vanish.

Parallel charging protocol. Next, we analyze the parallel charging protocol, for which we resort to the Tavis-Cummings (TC) model [90], which describes multiple mutually noninteracting qubits, each simultaneously coupled to the single-mode field confined to the cavity. This model cannot be formulated as a sequence of excitation exchanges, it thus provides insights into the nature of the Fock precision advantage, which cannot be reduced to a mere swap of well-defined Fock states. The M -qubit TC Hamiltonian can be derived from Eq. (1) by setting

$A_j(t) = 1 \forall j, t$. Such a model can be analytically solved for $M \in \{1, \dots, 4\}$ qubits under the perfect resonance condition, namely $\Delta\omega = 0$ [90–92]. We analytically study the case for $M = 2$ as a proof of concept, and numerically solve the cases for $M \geq 3$.

Here, we primarily considered the energy exchange of a single qubit, as well as the fidelity between the single qubit states $\hat{\rho}_{\text{qub}}^{(j)}(t)$ and $\hat{\rho}_{\text{target}}^{(j)} = |e\rangle_j \langle e|_j$. Since the model is symmetric with respect to any permutation of the qubits, this is equivalent to study the fidelity between the M -qubit collective states $\hat{\rho}_{\text{qub}}^{(1,\dots,M)}(t) = \bigotimes_{j=1}^M \hat{\rho}_{\text{qub}}^{(j)}(t)$ and $\hat{\rho}_{\text{target}}^{(1,\dots,M)} = \bigotimes_{j=1}^M |e\rangle_j \langle e|_j$.

The statistics related to the energetics of the single qubit can be derived by rotating the evolution operator as $e^{i\frac{\chi}{2}\hat{O}}\hat{U}(\tau, 0)e^{-i\frac{\chi}{2}\hat{O}}$ with $\hat{O} = \hat{H}_{\text{qub}}^{(1)} = \hbar\omega_{\text{qub}}\frac{1}{2}\hat{\sigma}_z^{(1)}$ (further details are provided in SM [78]). On the numerical side, an effective way to numerically calculate the qubit energy statistics is to solve a differential equation for the transformed density matrix $\hat{\rho}_\chi(t)$ of the total system [64],

$$\frac{d}{dt}\hat{\rho}_\chi(t) = -\frac{i}{\hbar}[\hat{H}_{\text{TC},\chi}\hat{\rho}_\chi(t) - \hat{\rho}_\chi(t)\hat{H}_{\text{TC},-\chi}], \quad (4)$$

where $\hat{H}_{\text{TC},\chi}$ is the modified TC Hamiltonian $\hat{H}_{\text{TC},\chi} = \hat{H}_{\text{free}} + \hat{H}_{\text{int},\chi}^{(1)}$, with $\hat{H}_{\text{int},\chi}^{(1)} = e^{i\frac{\chi}{2}\hat{H}_{\text{qub}}^{(1)}}\hat{H}_{\text{int}}e^{-i\frac{\chi}{2}\hat{H}_{\text{qub}}^{(1)}}$, and $\hat{H}_{\text{int}} = \hbar g \sum_{j=1}^M (\hat{\sigma}_+^{(j)}\hat{a} + \hat{\sigma}_-^{(j)}\hat{a}^\dagger)$. As reported in Fig. 2(b) for the parallel charging of $M = 5$ qubits within a single time window τ , the advantage in the SNR offered by the Fock state cannot be arbitrarily enhanced by optimizing the model parameters, unlike what occurs with the sequential charging protocol. As well, the charging of a single qubit cannot be perfectly performed, since $F < 1$ even for the Fock state protocol. Furthermore, as M increases, the performance of the Fock state cavity becomes increasingly comparable to that of the coherent and Gaussian states. For these reasons, we conclude that the perfect charging of a stack of qubits with a single-mode field can only be achieved by employing both a Fock state reservoir and a sequential protocol.

Realistic scenarios. Two questions arise: Is the sequential protocol robust in realistic scenarios? If so, what conditions maximize its precision described by the process-dependent SNR (2)? To address this question, we investigate a broad variety of situations, accounting for the effects of noise, dissipation, and nonideal parameter choices. We also relax the JC model by introducing several kinds of generalizations. These results are thoroughly presented in the Supplemental Material [78,93–100]. Here, as a paradigmatic example of a source of nonideality, we discuss the impact of thermal photons, which can degrade the quantum properties of the cavity state. We do so by simulating a sequential charging of $M = 5$ qubits. For each charging process of the entire stack of qubits, we set the mean number of initial photons $\langle n \rangle$ in the cavity equal to a value among $\{1, 2, 3, 4, 5\}$. We then simulate the battery charging process for both Fock and Gaussian initial cavity states. In all cases, we consider the worst-case scenario of a noisy Fock cavity state compared to an ideal, noiseless Gaussian state. To do so, we introduce the following quantity, parametrized by the number of thermal photons, $\bar{n}_{\text{th}}$: $\text{SNR}_{\bar{n}_{\text{th}},\langle n \rangle} \equiv \frac{1}{\langle n \rangle} \sum_{j=1}^5 \max_{\tau_j} \{\text{SNR}_j(\bar{n}_{\text{th}}, \langle n \rangle)\}$. It

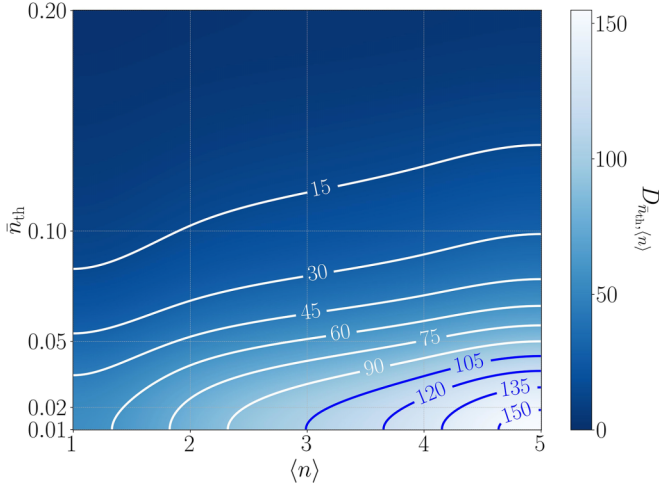


FIG. 3. Effect of the presence of $\bar{n}_{\text{th}}$ thermal photons on the difference $D_{\bar{n}_{\text{th}}, \langle n \rangle}$. To compare with phase-insensitive Fock states, we used a Gaussian coherent squeezed state $|\zeta, \bar{\alpha}\rangle\langle\zeta, \bar{\alpha}|$ with randomized phases [78,102]. Simulation parameters: $N = \langle n \rangle$ for the Fock state; ζ and $\bar{\alpha}$ are optimized to give the best performances for the squeezed coherent state; $g/\omega_{\text{qub}} = 10^{-2}$, and $\Delta\omega = 0$.

represents an averaged SNR, normalized with respect to the field energy. This ensures the fairness of our findings, as their accuracy is no longer limited by the resources (i.e., the initial number of photons) available to the experimenter. The figure of merit we use is defined as the difference between the quantity calculated for both the Fock and Gaussian cavity states:

$$D_{\bar{n}_{\text{th}}, \langle n \rangle} = \overline{\text{SNR}}_{\bar{n}_{\text{th}}, \langle n \rangle}^{\text{Fock}} - \overline{\text{SNR}}_{\bar{n}_{\text{th}}, \langle n \rangle}^{\text{Gaussian}}. \quad (5)$$

From the definition of $D_{\bar{n}_{\text{th}}, \langle n \rangle}$, it is evident that a precision advantage of the Fock state over the Gaussian state corresponds to the condition $D_{\bar{n}_{\text{th}}, \langle n \rangle} > 0$. Moreover, $D_{\bar{n}_{\text{th}}, \langle n \rangle}$ also accounts for situations in which a cavity state with $\langle n \rangle$ initial photons does not perfectly charge the first n qubits, but delivers the remaining energy to the subsequent ones.

To model the presence of thermal photons in the Fock state, we follow Ref. [101], where an N -photon Fock state exposed to a source of additive thermal noise for a short time is expressed in terms of the map $\mathcal{K}$ such that $\hat{\rho}_{\text{Fock}}^{\text{Th}} = \mathcal{K}(|N\rangle\langle N|) = \sum_{m=0}^{\infty} p_m(\bar{n}_{\text{th}})|m\rangle\langle m|$ (details are reported in SM [78]). From Fig. 3, we observe that the difference $D_{\bar{n}_{\text{th}}, \langle n \rangle}$ remains positive even for $\bar{n}_{\text{th}} = 0.2$ on average; for lower values, such as $\bar{n}_{\text{th}} = 0.02$, $D_{\bar{n}_{\text{th}}, \langle n \rangle}$ can become significantly large, ranging from $D_{\bar{n}_{\text{th}}, \langle n \rangle=1} \sim 50$ to $D_{\bar{n}_{\text{th}}, \langle n \rangle=5} \sim 150$. Therefore, the precision

advantage of the Fock state charging protocol proves robust against thermal noise, as long as the noise level remains below a critical threshold. Notably, the performance of the non-Gaussian quantum protocol improves with the mean photon number $\langle n \rangle$; furthermore, the tolerance to noise scales with the initial photon number N stored in the Fock state, reinforcing the protocol's resilience at higher photon populations.

Conclusions. We demonstrated that a multiqubit Jaynes-Cummings (JC) quantum battery can achieve maximally precise charging, i.e., high-precision performances even in presence of non-ideal working conditions, by using an optimized protocol and harnessing non-Gaussian quantum states. We employed the Full Counting Statistics (FCS) to compute the process-dependent Signal-to-Noise Ratio (SNR) associated with the energy injected into the battery, and compared sequential and parallel charging protocols of M qubits. Our results show that under ideal, noiseless conditions, a cavity prepared in a N -photon Fock state can perfectly charge $M = N$ qubits sequentially, whereas parallel charging fails to achieve perfect charging. We then focused on the sequential protocol, evaluating its robustness in realistic scenarios, here exemplified by the detrimental effects of spurious thermal photons. A more complete analysis and thorough discussions are developed in SM [78]. Despite the negative impact of the considered external limiting factors, we demonstrate that the maximum charging precision is achievable when using a genuine non-Gaussian quantum state for sequential qubit charging. The proposed protocol, albeit simple, is therefore proven to be resilient in realistic working conditions, thus being feasible for direct implementations. Such a result serves as a clear signature of the battery quantum nature, and suggests that this effect should be experimentally observable: it applies indeed to recent experiments in circuit QED [103], trapped ions [104–106], and quantum electromechanics [107].

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Data availability. The data that support the findings of this article are openly available [108].

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