

Thresholded quantum sensing with a frustrated Kitaev trimer

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(Received 31 July 2025; accepted 9 February 2026; published 6 March 2026)

We investigate the response of a Ramsey interferometric quantum sensor based on a frustrated, three-spin system (a Kitaev trimer) to a classical time-dependent field (signal). The system eigenspectrum is symmetric about a critical point, $|\mathbf{b}| = 0$, with four of the spectral components varying approximately linearly with the magnetic field and four exhibiting a nonlinear dependence. Under the adiabatic approximation and for appropriate initial states, we show that the sensor's response to a zero-mean signal is such that below a threshold, $|\mathbf{b}| < b_{\text{th}}$, the sensor does not respond to the signal, whereas above the threshold, the sensor acts as a detector that the signal has occurred. This thresholded response is approximately omnidirectional. Moreover, when deployed in an entangled multisensor configuration, the sensor achieves sensitivity at the Heisenberg limit. Such detectors could be useful both as stand-alone units for signal detection above a noise threshold and in two- or three-dimensional arrays, analogous to a quantum bubble chamber, for applications such as particle track detection and long-baseline telescropy.

DOI: [10.1103/3z2c-2kkl](#)

Introduction. Quantum sensing, where a quantum state is prepared, allowed to interact with a signal of interest, and then measured, is used to extract information about the signal [1–3]. By using entangled input states across N sensors, one can enhance the sensitivity beyond the standard quantum limit, achieving the so-called Heisenberg limit [1,2,4–7]. This improves the standard error from $\Delta/\sqrt{N}$ to Δ/N , where Δ is the standard deviation of the signal and N is the number of measurements, yielding a $\sqrt{N}$ improvement in scaling.

Motivated by the critical behavior of quantum materials, where properties can change rapidly with large gradients near phase transitions [8–12], we introduce a minimal realization of a critical, thresholded quantum rectifying sensor, based on a frustrated three-spin Kitaev trimer (Fig. 1). Due to its frustrated nature, this structure exhibits behavior associated with a broken symmetry about a critical point (a crossing) in its spectrum. Similar frustrated trimers have been studied extensively at material scales both theoretically and experimentally using spin-1/2 lattices [13,14] and spin-1 lattices [15] in the presence of an external magnetic field.

Although quantum sensors are most commonly used for precise estimation of the mean phase encoded in the state of a sensor by a classical signal [16–26], they can also be designed

to probe higher-order moments, such as a signal's second-order moment [27,28], allowing for the study of quantum fluctuations. This type of rectifying sensor can be deployed as a detector that indicates *whether* a signal was sensed, rather than measuring the signal amplitude itself. One key challenge with quantum detection protocols is that, in the presence of a noisy background, the quantum detector can be rendered useless due to its sensitivity. This is particularly true for rectifying sensors, since the accrued phase increases rapidly. This Letter targets this issue. Here, we present a thresholded quantum rectifier that accumulates a phase proportional to a signal's second moment only above a threshold, $|\mathbf{b}| = b_{\text{th}}$. The threshold allows the sensor to selectively measure variance in excess of the threshold that, otherwise, could rapidly push the sensor beyond the monotonic region of response (up to the first fringe) causing it to give multivalued output and rendering it unusable.

This type of quantum sensor may be used independently or in a multisensor context with an entangled input state, achieving the Heisenberg limit. In an array configuration, such sensors may be used to measure particle tracks [29] in a similar way to the tracks measured in a bubble chamber. In photonics, this type of interferometry could be used to provide longer aperture times for quantum-enhanced long-baseline telescropy [30–32].

The thresholded rectifier. The thresholded rectifier (sketched in Fig. 1) is based on the physics of a frustrated Heisenberg-Kitaev trimer [33,34]. Without loss of generality (see the Supplemental Material [35]), the system may be described by the time-dependent Hamiltonian

$$H(t) = X_1 X_2 + Y_2 Y_3 + Z_1 Z_3 + \mathbf{b}(t) \cdot \boldsymbol{\sigma}. \quad (1)$$

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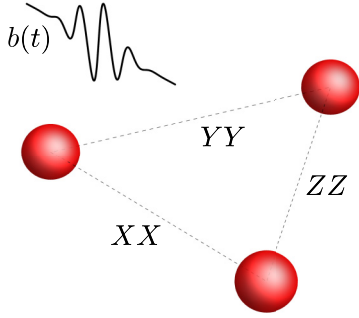


FIG. 1. Kitaev trimer sensor structure. The thresholded quantum rectifying sensor is based on a frustrated Kitaev trimer. Here, we illustrate a magnetic field impinging on the trimer from a particular direction (black waveform). Each vertex corresponds to a spin, and couplings along the edges encode interaction strengths.

Here, $\mathbf{b}(t) = (b_x(t), b_y(t), b_z(t))$ is a time-dependent classical external field and $\boldsymbol{\sigma} = (\sum_i X_i, \sum_i Y_i, \sum_i Z_i)$ is a vector of the collective operators that mediate the interaction between the classical field and the quantum sensor, where $\{X_i, Y_i, Z_i\}$ correspond to the Pauli matrices acting on the i th spin. The b -dependent spectrum arising from this Hamiltonian, for a field incident along one of the cardinal axes, is shown in Fig. 2.

In the following, we assume—supported in detail in the Supplemental Material [35]—that the sensor dynamics are well described by the adiabatic approximation [36]. Starting from an initial state $|\phi_0\rangle$, the probability of the final state

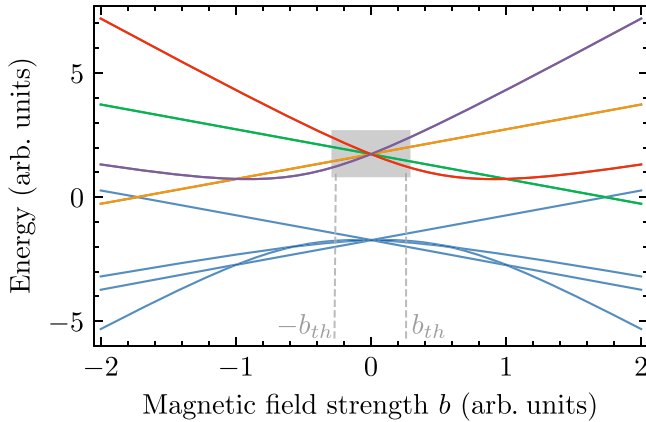


FIG. 2. Kitaev trimer eigenspectrum. We plot the eigenenergies λ_i as a function of the magnetic field amplitude b , with the field aligned along a cardinal axis. The eigenenergies associated with a standard sensor (green and gold lines) correspond to an interferometer with support on two linear branches. In contrast, the mousetrap sensor is associated with either the (asymmetric) gold-red pair or the green-purple pair of eigenenergies. Within the gray-shaded region, where the spectrum remains approximately linear, the rectifier behaves like a standard sensor. However, when $|b| > b_{th}$, curvature of the eigenspectrum leads to an additional phase proportional to the second moment of the signal, $b^2(t)$. Thus, the sensor acts as a rectifier (i.e., becomes sensitive to the signal's second moment) only above the threshold b_{th} .

remaining the initial state is given by

$$P[\phi_0](t) = |\langle \phi_0 | U(t) | \phi_0 \rangle|^2 = \left| \langle \phi_0 | \Psi_0 \exp \left[-i \int_{t_0}^t d\tau \Lambda(\tau) \right] \Psi_0^\dagger | \phi_0 \rangle \right|^2, \quad (2)$$

where $\Lambda(t)$ is a diagonal matrix describing the eigenenergies of $H(t)$ and Ψ_0 is a matrix describing the eigenstates of $H(t=0)$. To use the sensor as a Ramsey interferometer, we prepare the initial state $|\phi_0\rangle$ such that $\Psi_0^\dagger |\phi_0\rangle = |+\rangle$, where $|+\rangle$ denotes the superposition of two eigenstates, $|+\rangle = (|\lambda_p\rangle + |\lambda_q\rangle)/\sqrt{2}$. Under this protocol, the probability of remaining in the initial state is given by

$$P[\phi_0](t) = \cos^2(\chi(t)), \quad (3)$$

where we have defined the accumulated phase

$$\chi(t) \equiv \frac{1}{2} \int_{t_0}^t (\lambda_p(\tau) - \lambda_q(\tau)) d\tau, \quad (4)$$

with $\lambda_j(t)$ denoting the instantaneous eigenenergy associated with eigenstate $|\lambda_j(t)\rangle$. A detailed derivation is provided in the Supplemental Material [35].

For N entangled sensors with initial state, $|\Phi_+\rangle \equiv (|\lambda_p\rangle^{\otimes N} + |\lambda_q\rangle^{\otimes N})/\sqrt{2}$, the resulting response—derived in the Supplemental Material [35]—is given by

$$P[\Phi_+](t) = \cos^2(N\chi(t)). \quad (5)$$

This leads to a quadratic Heisenberg scaling of the quantum Fisher information, $\mathcal{F}_\chi = 4N^2$, with the variance of the accumulated phase lower bounded by the quantum Cramér-Rao bound of

$$(\Delta\chi(t))^2 \geq \frac{1}{4N^2}. \quad (6)$$

Results. We now analyze the behavior of the thresholded quantum rectifier, with eigenspectrum shown in Fig. 2, for the example of a b field in the z direction. The eigenenergies extend outward from two gapped critical points at energies of $\pm\sqrt{3}$ at $|b|=0$. Our focus will be on the upper critical point, where four spectral lines emerge. Two of these lines (shown in gold and green) are linear in the b field. If the initial state $|+\rangle$ has support only on the eigenstates corresponding to these spectral lines, the sensor behaves as a standard sensor that measures the mean of the b field, i.e.,

$$P[\phi_0](t) = \cos^2(N\chi_{std}(t)) = \cos^2\left(N \int_{t_0}^t c b(\tau) d\tau\right), \quad (7)$$

where c characterizes the slope of the linear dependence.

In contrast, if the initial state has support on either the green-purple or gold-red pairs of spectral lines depicted, the sensor response differs. For small field amplitudes, $|b| < b_{th}$, indicated by the gray-shaded region, the eigenspectrum remains approximately linear and the sensor effectively measures the mean. However, above this threshold, $|b| > b_{th}$, the linear symmetry of the eigenspectrum breaks to the left for the green-purple pair and to the right for the gold-red pair, and the eigenenergies become nonlinear in b . In this regime, the accumulated phase becomes approximately proportional to $b^2(t)$, as shown in the Supplemental Material [35]. The device therefore acts as a thresholded rectifier, detecting that

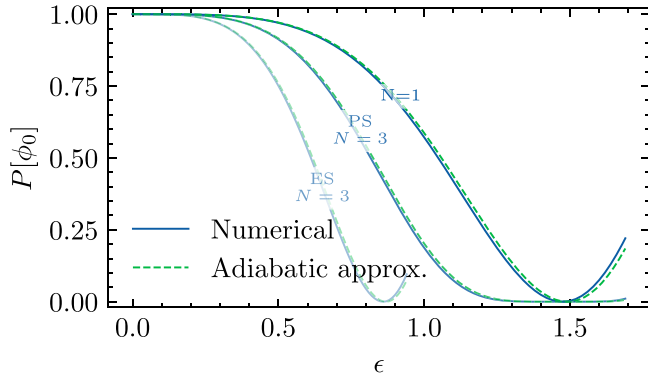


FIG. 3. Trimer mousetrap sensor response. The response probability $P[\phi_0]$ shown as a function of the signal amplitude ϵ for a zero-mean, oscillatory signal with Gaussian envelope, $b(t) = \epsilon \exp(-t^2/s) \sin(2\pi kt)$, where $k = 0.05$, $s = 20$, $t_0 = -15$, and $t = 15$. Results are shown for $N = 1$ and $N = 3$ to illustrate how the response scales with the number of sensor copies. Both product state (PS) and entangled state (ES) are used as inputs for $N = 3$. Note that a response is not incurred until ϵ is large enough (≈ 0.3) to be affected by the change in curvature of λ_i near the asymmetric wells in the spectrum, with basins centered at $b = \pm 0.8$ (see Fig. 2).

the field has exceeded b_{th} . These two pairs of spectral lines result in what we call *quantum mousetrap sensors*.

The response probability for the mousetrap sensor is $P[\phi_0] = \cos^2(N\chi_{\text{mt}}(t))$, where

$$\chi_{\text{mt}}(t) = \frac{1}{2} \int_{t_0}^t (\sqrt{3} + 2b(\tau) - \sqrt{3 - 4b(\tau) + 4b^2(\tau)}) d\tau. \quad (8)$$

In Fig. 3, we plot $P[\phi_0]$ for a zero-mean, oscillatory b field with Gaussian envelope. Three response probabilities are shown, one for $N = 1$, a single spin, one for $N = 3$ with product state input giving the standard quantum limit, and one for $N = 3$ with entangled state input giving the evident Heisenberg limit.

The mousetrap sensor response over the range $\epsilon \in (0, 1.49)$ is shown in Fig. 4. The plateau observed below the threshold at $\epsilon \sim 0.3$ in Fig. 3 can be understood by examining the integral, χ_{mt} (Fig. 4, inset of each panel), in Eq. (8). When the signal probes only the approximately linear region of the spectrum (highlighted by the gray box in Fig. 2 about the point $(0, 0)$, and depicted in Fig. 4, top-left panel), the zero-mean signal oscillates back and forth from positive to negative values, with $P[\phi_0]$ averaging to zero, giving no accumulated response. In contrast, for $\epsilon \gtrsim 0.3$, the eigenenergies deviate from linearity, and the signal accrues. Thus, the mousetrap sensor functions as a thresholded rectifier, sensitive to signal amplitudes in the approximate range $\epsilon \in [0.3, 1.5]$ with the maximum (first fringe) response near $\epsilon \sim 1.5$. This operating window is stated in dimensionless units; see the Supplemental Material [35] for the nondimensionalization that maps physical field amplitudes and pulse durations.

Although the adiabatic approximation can fail near degeneracies, the present regime is well described within the adiabatic treatment used here. Near the critical point, the spectrum becomes approximately linear and admits a closed-form

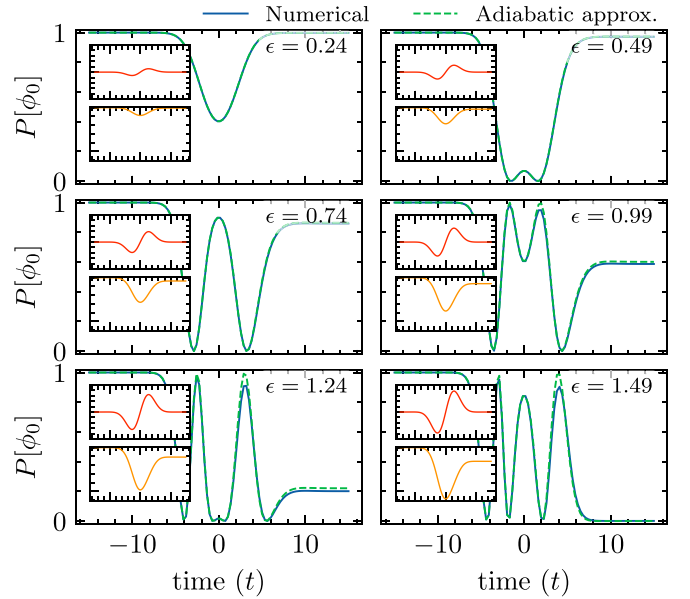


FIG. 4. Mousetrap sensor response. The probability $P[\phi_0]$ is shown as a function of time t for a range of signal amplitudes ϵ for a zero-mean, oscillatory signal with Gaussian envelope, $b(t) = \epsilon \exp(-t^2/s) \sin(2\pi kt)$, where $k = 0.05$, $s = 20$, $t_0 = -15$, and $t = 15$. Insets show the signal $b(t)$ (top left in red) and the accumulated phase $\chi_{\text{mt}}(t)$ (lower left in orange), both plotted on the same timescale as the main figure, with vertical ranges of $(-1, 1)$ and $(-6, 0)$, respectively. For small amplitudes ($\epsilon \lesssim 0.3$), the signal remains within the approximately linear region of the potential well, resulting in minimal phase accumulation and no integrated response. As ϵ exceeds the threshold near 0.3, portions of the signal probe the nonlinear regime, where symmetry breaking in the spectrum leads to phase accumulation. The sensor thus acts as a signal rectifier responding only when the signal amplitude crosses b_{th} , where it begins to integrate deviations from linearity.

adiabatic description. Consequently, the adiabatic approximation remains highly accurate for the mousetrap protocol, as verified numerically by comparison to full time-dependent simulations (see Figs. 3 and 4 and the Supplemental Material [35]). We have additionally shown that more realistic Ornstein-Uhlenbeck processes with appropriate bandwidth also exhibit mousetrap behavior (see the Supplemental Material [35]).

To assess directional response, we numerically studied the mousetrap's behavior as a function of magnetic field incidence angle for fixed initial state (with a sensor initialized to eigenstates optimal for incidence along the z axis, but tested for all incidence directions in the octant). As shown in Fig. 5, while specific parameter values vary for different incidence angles, the mousetrap consistently exhibited the same response to b fields at all angles: a plateau in probability for small signal amplitudes ($\epsilon < 0.3$) followed by a significant response from $0.3 < \epsilon < 1.5$, consistent with Fig. 3. At higher amplitudes ($\epsilon \approx 0.7$), the response becomes increasingly sensitive to the incident angle; still, there is always a linear region around the critical point giving rise to the mousetrap thresholding effect. As a result, the quantum mousetrap can be used as an omnidirectional, thresholded rectifier and, hence, can be used in a two- or three-dimensional array to detect incident b fields

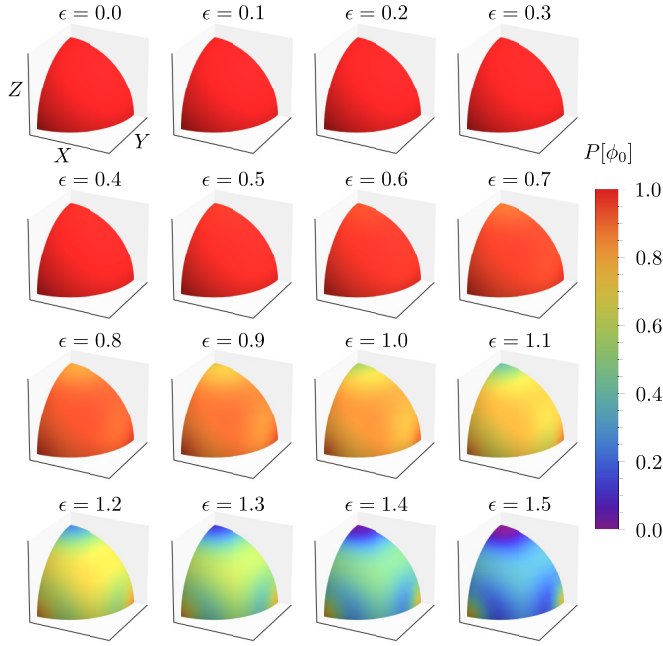


FIG. 5. Omnidirectional mousetrap threshold response. This figure shows the probability $P[\phi_0]$ of a mousetrap remaining in its initial state as a function of the incident direction of the b field, using the same Gaussian signal and parameters as in Figs. 3 and 4. Due to the eightfold symmetry of the response, only one octant of the sphere is shown. The response color mapping is indicated by the legend. At low amplitudes ($\epsilon \lesssim 0.3$), the sensor remains unresponsive across all directions (red), consistent with the plateau observed in Fig. 3. As ϵ increases, directional sensitivity begins to emerge around $\epsilon \approx 0.7$ (orange, yellow, green, blue), with significant anisotropic variations appearing in the range $0.7 < \epsilon < 1.5$, where the signal begins to probe the nonlinear regions of the spectrum.

in all directions. This frustration-induced omnidirectionality is a practical advantage of the trimer architecture, since the thresholded sensing characteristics persist under variations in the signal direction.

Many different combinations of the trimer eigenspectra may be used to construct quantum sensors. Of these, two are thresholded—the mousetrap and an additional sensor with

$$\chi_{\text{mt}}(t) = \frac{1}{2} \int_{t_0}^t (-2b(\tau) - \sqrt{3 - 4b(\tau) + 4b^2(\tau)} + \sqrt{3 + 4b(\tau) + 4b^2(\tau)}) d\tau, \quad (9)$$

which may be used as a thresholded sensor of the third moment of $b(t)$ (the skew, from purple and red eigenstates in Fig. 2).

Due to their complexity, quantum mousetraps can be difficult to implement. In the Supplemental Material [35], we outline an experimental protocol for realizing a mousetrap using a neutral-atom (Rydberg) platform. In brief, the protocol involves initializing a linear array of $3N$ qubits prepared in the standard Greenberg-Horne-Zeilinger (GHZ) state, $|\text{GHZ}\rangle = (|0\rangle^{\otimes 3N} + |1\rangle^{\otimes 3N})/\sqrt{2}$; rearranging the $3N$ qubits into trimers; and finally, quenching the system to $H(p=0) = H_{\text{init}}$ and adiabatically sweeping the parameter p to 1. This

prepares the desired GHZ-like state. Finally, as in most GHZ-based sensing protocols, the preparation sequence is reversed after the phase-acquisition stage so that the accumulated phase can be read out in the computational basis. This implementation perspective naturally suggests a modular route to scaling, in which the trimer acts as a minimal sensing unit that can be replicated to build an N -trimer sensor. Entanglement resources (e.g., GHZ-type states across trimers) then provide a direct route to quantum-enhanced scaling, with the main additional challenge being preparation of the required GHZ-like eigenstate superposition across trimers.

Discussion. The thresholded quantum rectifier and the additional sensors outlined in the Supplemental Material [35] may be combined to measure quantities of interest for the analysis of the input b field. For instance, we can use a standard sensor and a Landau-Zener sensor to measure the first ($\int_{t_0}^t b(\tau) d\tau = \arccos \sqrt{P_{\text{std}}[+]}$) and second ($\int_{t_0}^t b^2(\tau) d\tau = \arccos \sqrt{P_{\text{LZ}}[+]}$) moments, respectively, to postcompute the variance $\text{Var}[b(t)] = \int_{t_0}^t b^2(\tau) d\tau - |\int_{t_0}^t b(\tau) d\tau|^2$. Similarly, for a thresholded rectifier, if the mean is nonzero, it can be separately measured with a standard sensor and subtracted.

There has been recent interest in making use of quantum computers to postprocess quantum data acquired by quantum sensors [29,37,38]. Since they are coherent, mousetraps and other spin systems, such as those in the Supplemental Material [35], could be considered as quantum sensing modules that preprocess the quantum data before incorporation in larger quantum computational systems, such as quantum phase estimation [38,39]. From this point of view, they are a first example of quantum computation at the edge [38,40,41].

Acknowledgments. We thank Arnab Banerjee, Daniel Bowring, David Weld, Tyler Volkoff, and Guillem Müller-Rigat for helpful discussions. A.K.S. was a participant in the 2024 Quantum Computing Summer School at LANL. This work was supported by the U.S. Department of Energy (DOE) through a quantum computing program sponsored by the Los Alamos National Laboratory (LANL) Information Science & Technology Institute. This material is based upon work supported by the U.S. Department of Energy, Office of Science, National Quantum Information Science Research Centers, Quantum Science Center (QSC). A.T.S. was funded by the QSC during the conception of the project. A.T.S. and C.H.A. acknowledge support for a numerical feasibility study of nonlinear quantum sensors from LANL's LDRD Program under project 20248110CT-IMS. Subsequent numerical and analytical work was performed via the QSC. C.H.A. was also supported by the U.S. Department of Energy, Office of Science, Office of Advanced Scientific Computing Research, under the Computational Partnerships program. A.K.S. acknowledges support from the European Union's Horizon 2020 Research and Innovation Programme under the Marie Skłodowska-Curie Grant Agreement No. 847517; MCIN/AEI (PGC2018-0910.13039/501100011033, CEX2019-000910-S/10.13039/501100011033, Plan National STAMEENA PID2022-139099NB-I00, project funded by MCIN and by the "European Union NextGenerationEU/PRTR" (PRTR-C17.I1), FPI); Ministry for Digital Transformation and of Civil Service of the Spanish Government through the

QUANTUM ENIA project call—Quantum Spain project, and by the European Union through the Recovery, Transformation and Resilience Plan—NextGenerationEU within the framework of the Digital Spain 2026 Agenda; CEX2024-001490-S [MICIU/AEI/10.13039/501100011033]; Fundació Cellex; Fundació Mir-Puig; Generalitat de Catalunya (European Social Fund FEDER and CERCA program); Barcelona Supercomputing Center MareNostrum (FI-2023-3-0024); HORIZON-CL4-2022-QUANTUM-02-SGA PASQuanS2.1,

101113690, EU Horizon 2020 FET-OPEN OPTologic, Grant No. 899794; QU-ATTO, 101168628; EU Horizon Europe research and innovation program under Grant Agreement No. 101080086 NeQST.

This paper is approved by Los Alamos National Laboratory for universal release under designation LA-UR-25-24691.

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.

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