

Autonomous Quantum Shutter: Bogoliubov Particle Production, Backreaction, and Local–Global Separation in a Matter-Controlled Bosonic Lattice

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A change of field boundary conditions produces quanta when it mixes positive- and negative-frequency modes, but the boundary motion is usually prescribed externally. We construct a finite autonomous quantum shutter in which a quantum matter register controls one local bilinear of a one-dimensional bosonic lattice while the complete matter–field Hamiltonian remains fixed and time independent. The controller-gated term is $C = \eta\kappa q_\ell q_r$, and at full closure the open and closed field Hamiltonians obey the exact local identity $H_{\text{cl}} - H_{\text{op}} = \kappa q_\ell q_r$. An abrupt opening therefore gives a canonical Bogoliubov transformation with nonzero β , so the closed vacuum is squeezed in the open basis and a closed-basis one-quantum state becomes a one-quantum-added squeezed state. Exact sparse evolution and canonical matrix-product-state evolution show autonomous packet transmission between the static closed and open limits, controller backreaction, matter–field entanglement, and energy redistribution with total-energy drift below 4×10^{-13} in the canonical benchmark. After fixing outgoing modes independently from a prescribed smooth-history reference, far local field algebras can be close at the 10^{-4} – 10^{-6} level while the joint outgoing state remains distinguishable. A collective controller of size M approaches the prescribed classical switch with leading corrections proportional to $1/M$, in agreement with a spin-coherent fluctuation expansion. The controller register also has an exact embedding into ordinary matter qubits through frustration projectors and a unary clock sector. The construction is a regulated scalar/bosonic realization of standard dynamical-Casimir physics; it is not a derivation of electromagnetic $U(1)$ or of an unregulated infinite particle-number tail.

I. INTRODUCTION

Time-dependent boundaries and media can mix positive- and negative-frequency field modes. In quantum theory the corresponding in/out operators are related by a Bogoliubov transformation, and a nonzero anomalous block produces particles from the initial vacuum. This mechanism underlies the dynamical Casimir effect (DCE), from Moore’s variable-length cavity and the Davies–Fulling moving mirror to modern circuit realizations and reviews [1–7]. In most formulations, however, the mirror trajectory, refractive-index modulation, or boundary protocol is a prescribed classical function. The field is quantum, whereas the actuator is external.

There are good reasons to quantize the actuator. A finite boundary can exchange energy with the field, become entangled with it, and record which scattering history occurred. Quantum and optomechanical treatments of DCE backreaction have emphasized precisely these effects [8–10]. A related locality question has recently appeared in the theoretical “truncated photon” problem: a shuttered one-photon wave packet can have simple local descriptions on separated regions even though its global state contains a complicated particle-number structure [11].

Here we ask a narrower constructive question. Can a local finite quantum register autonomously open a shutter under a fixed Hamiltonian, while reproducing the standard Bogoliubov mechanism and retaining explicit controller backreaction? We answer yes for a regulated one-dimensional bosonic lattice. Each field site is a square-root ladder, the central boundary is a single edge–edge bilinear, and ordinary matter degrees of freedom gate that bilinear. No collapse and no externally inserted $H(t)$ are

required.

The main results are as follows. First, the closed/open difference is an exact local identity, and the sudden switch has an analytic nonzero Bogoliubov β block. Second, exact Fock-space calculations exhibit autonomous scattering, energy exchange, controller–field entanglement, and a reduced number distribution with support in every represented sector. Third, independent exact, Gaussian, and matrix-product-state solvers converge across the active cutoffs. Fourth, complete far-local observable algebras can agree closely with an independently prescribed smooth-history reference even when the global and joint outgoing states remain distinguishable. Fifth, the discrepancy between a finite quantum controller and its mean switching profile vanishes approximately as $1/M$ for a collective controller, and the coefficient is explained by a spin-coherent fluctuation expansion. Finally, the abstract controller is exactly embedded in a local matter-qubit clock sector controlled by frustration projectors.

The field-theoretic mechanism is not new: it is the DCE/Bogoliubov mechanism. The contribution is the autonomous matter–edge architecture, its controlled finite-regulator realization, and the quantitative local/global and quantum-to-classical analyses. All claims below are scoped to that model. In particular, the model is a scalar lattice regulator with a finite quantum actuator; it is not a derivation of electromagnetic gauge theory, a continuum mirror theorem, or a claim of unbounded particle number in any finite Hilbert space. The numerical claims are tied to explicit regulator checks and to the reproducibility archive described at the end of the manuscript.

II. MATTER-CONTROLLED EDGE-LADDER MODEL

We set $\hbar = 1$. The field consists of N bosonic edge modes. In a local dimension- d regulator,

$$a_{j,d}|n\rangle = \sqrt{n}|n-1\rangle, \quad n = 0, \dots, d-1, \quad (1)$$

with

$$[a_{j,d}, a_{j,d}^\dagger] = \mathbb{I} - d|d-1\rangle\langle d-1|. \quad (2)$$

Thus the oscillator algebra is exact away from the top rung and converges in the low-occupation sector. Define

$$q_j = \frac{a_j + a_j^\dagger}{\sqrt{2\omega_0}}, \quad p_j = -i\sqrt{\frac{\omega_0}{2}}(a_j - a_j^\dagger). \quad (3)$$

A. Open and closed field Hamiltonians

The open chain is

$$H_{\text{op}} = \frac{1}{2} \sum_{j=1}^N (p_j^2 + \omega_0^2 q_j^2) + \frac{\kappa}{2} (q_1^2 + q_N^2) + \frac{\kappa}{2} \sum_{j=1}^{N-1} (q_{j+1} - q_j)^2. \quad (4)$$

For even N , let $\ell = N/2$ and $r = \ell + 1$. The open central spring contributes

$$h_{\text{op}}^{(\ell r)} = \frac{\kappa}{2} (q_r - q_\ell)^2. \quad (5)$$

A fully closed central shutter replaces this spring by two local pins,

$$h_{\text{cl}}^{(\ell r)} = \frac{\kappa}{2} (q_\ell^2 + q_r^2), \quad (6)$$

so that

$$\boxed{H_{\text{cl}} - H_{\text{op}} = \kappa q_\ell q_r.} \quad (7)$$

We also use a finite barrier

$$C = \eta \kappa q_\ell q_r, \quad 0 \leq \eta \leq 1, \quad (8)$$

for which $\eta = 1$ is the exact link removal and $\eta < 1$ is a partial cancellation of the central coupling.

B. Autonomous controller

Let $\mathcal{H}_C$ be a finite controller Hilbert space, H_C its free Hamiltonian, and P_C an operator whose closed-sector eigenvalue is one. The complete Hamiltonian is

$$\boxed{H_{\text{tot}} = \mathbb{I}_C \otimes H_{\text{op}} + P_C \otimes C + H_C \otimes \mathbb{I}_F.} \quad (9)$$

It is fixed and time independent. Field energy need not be conserved separately, but H_{tot} is.

Two controller realizations are used. A propagating clock of dimension R has

$$H_C = -J \sum_{s=0}^{R-2} (|s+1\rangle\langle s| + |s\rangle\langle s+1|), \quad P_C = \sum_{s < s_*} |s\rangle\langle s|. \quad (10)$$

A wave packet initially supported in $s < s_*$ propagates into the open sector. A minimal smooth controller uses two matter qubits,

$$F = \frac{\mathbb{I} - Z_0 Z_1}{2}, \quad H_C = \frac{\Omega}{2} X_0, \quad (11)$$

with initial state $|01\rangle$. In the absence of field backreaction,

$$p_{\text{cl}}^{\text{free}}(t) = \langle F(t) \rangle = \cos^2\left(\frac{\Omega t}{2}\right). \quad (12)$$

This provides an independently fixed smooth-history reference,

$$H_{\text{ref}}(t) = H_{\text{op}} + p_{\text{cl}}^{\text{free}}(t)C, \quad (13)$$

without altering the autonomous evolution in Eq. (9).

C. Initial states and observables

The vacuum input is the closed-field ground state $|0_{\text{cl}}\rangle$. A localized incident packet is prepared as

$$|\psi_{\text{pkt}}(0)\rangle = \frac{B_f^\dagger |0_{\text{cl}}\rangle}{\|B_f^\dagger |0_{\text{cl}}\rangle\|}, \quad B_f^\dagger = \sum_m f_m b_m^\dagger, \quad (14)$$

where b_m are open-chain positive-frequency modes and f_m is a Gaussian wave packet. Field observables are evaluated as packet-minus-vacuum differences, which removes the squeezed background. If Δn_j is the local excess and the central sites are excluded, we define

$$\mathcal{T} = \frac{N_R}{N_L + N_R}, \quad N_{L/R} = \sum_{j \in L/R} \Delta n_j. \quad (15)$$

We also monitor controller entropy, sector-resolved energy, reduced number distributions, and the complete local distinguishability

$$D_A = \frac{1}{2} \|\rho_A^{\text{auto}} - \rho_A^{\text{ref}}\|_1. \quad (16)$$

III. SUDDEN OPENING: ANALYTIC BOGOLIUBOV BENCHMARK

In the untruncated quadratic theory, write

$$V_{\text{cl}} U_{\text{cl}} = U_{\text{cl}} \omega^2, \quad V_{\text{op}} U_{\text{op}} = U_{\text{op}} \Omega^2, \quad S = U_{\text{op}}^T U_{\text{cl}}. \quad (17)$$

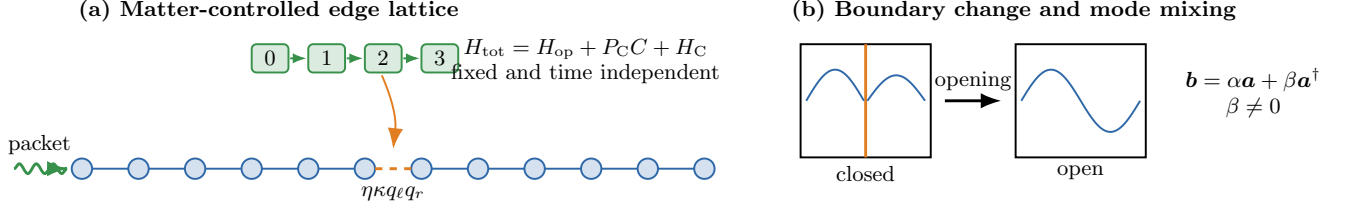


FIG. 1. Model and mechanism. (a) A finite matter/controller register gates one local edge–edge bilinear in a square-root bosonic lattice. (b) Opening changes the positive-frequency structure; the old and new mode operators are related by a Bogoliubov transformation with an anomalous block.

Continuity of $\mathbf{q}$ and $\mathbf{p}$ at an abrupt opening gives

$$\mathbf{b} = \alpha \mathbf{a} + \beta \mathbf{a}^\dagger, \quad (18)$$

with

$$\alpha = \frac{1}{2} \left(\Omega^{1/2} S \omega^{-1/2} + \Omega^{-1/2} S \omega^{1/2} \right), \quad (19)$$

$$\beta = \frac{1}{2} \left(\Omega^{1/2} S \omega^{-1/2} - \Omega^{-1/2} S \omega^{1/2} \right). \quad (20)$$

The canonical identities are

$$\alpha \alpha^\dagger - \beta \beta^\dagger = \mathbb{I}, \quad \alpha \beta^T - \beta \alpha^T = 0. \quad (21)$$

Numerically, their normalized residuals are below 2.2×10^{-15} in the lattice benchmarks.

The closed vacuum is a multimode squeezed state in the open basis,

$$|0_{\text{cl}}\rangle = \mathcal{N} \exp\left(\frac{1}{2} \mathbf{b}^{\dagger T} Z \mathbf{b}\right) |0_{\text{op}}\rangle, \quad Z = (\alpha^\dagger)^{-1} \beta^T, \quad (22)$$

and a normalized closed-basis one-quantum wave packet becomes a one-quantum-added squeezed state. Its open-basis mean number is

$$\langle N_b \rangle_{1_f} = \text{Tr}(\beta \beta^\dagger) + 1 + 2 \|\beta f^*\|^2. \quad (23)$$

The terms are vacuum production, the injected quantum, and stimulated pair production.

For a perfect central cut in a continuum finite box, the coefficients can be obtained in closed form. For odd open-mode index m and a left-half closed mode n ,

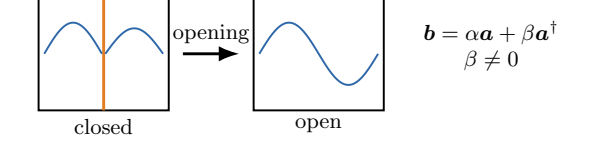
$$\beta_{mn}^L = (-1)^{n+(m-1)/2} \frac{2\sqrt{n}}{\pi\sqrt{m}(m+2n)}, \quad (24)$$

while the degenerate channel $m = 2n$ has nonzero overlap but $\beta = 0$. With the symmetric cutoff used here,

$$N_{\text{vac}}(K) = \frac{1}{2\pi^2} (\log K)^2 + O(\log K). \quad (25)$$

The derivation is given in Appendix A. Equation (25) is a regulator statement for the ideal abrupt cut; all autonomous simulations below use finite lattices and, where stated, finite barriers or smooth openings.

(b) Boundary change and mode mixing



IV. AUTONOMOUS SCATTERING, BACKREACTION, AND ROUTING

The first end-to-end benchmark uses the propagating clock in Eq. (10) with

$$N = 18, \quad n_{\text{max}} = 5, \quad \omega_0 = 3, \quad \kappa = 1.2, \quad \eta = 0.95, \quad (26)$$

$$R = 10, \quad J = 0.14, \quad s_* = 3, \quad T = 24. \quad (27)$$

The field basis contains all occupation configurations satisfying $\sum_j n_j \leq n_{\text{max}}$; its dimension is 33649, and the joint dimension is 336490.

The controller begins with $\langle P_C \rangle = 0.999994$. For a vacuum input it ends at 0.143485, while the incident packet leaves it at 0.216360. The packet therefore backreacts on the shutter and makes the closed sector more likely. The corresponding transmissions are

$$\begin{aligned} \mathcal{T}_{\text{closed}} &= 0.0074146, & \mathcal{T}_{\text{auto}} &= 0.7143166, \\ \mathcal{T}_{\text{open}} &= 0.9739391. \end{aligned} \quad (28)$$

The autonomous output lies between the static controls rather than being a basis relabeling.

The packet produces controller–field entropy

$$S_{\text{C:F}}^{\text{pkt}} = 0.671381 \text{ nats}, \quad (29)$$

compared with 8.27×10^{-4} nats for the vacuum. Conditioning the same global state on the final controller sector gives a right-side fraction 0.197895 in the closed branch and 0.835306 in the open branch. The controller therefore functions as a quantum router and carries which-history information.

Energy is redistributed rather than created. In the packet run,

$$\begin{aligned} \Delta E_F &= -3.2467 \times 10^{-3}, & \Delta E_C &= +3.6678 \times 10^{-3}, \\ \Delta E_{\text{int}} &= -4.2113 \times 10^{-4}. \end{aligned} \quad (30)$$

whose sum is 1.56×10^{-13} . The maximum total-energy drift over the trajectory is 3.98×10^{-13} . The sign of an individual sector exchange is protocol dependent; the invariant statement is autonomous exchange under a conserved total Hamiltonian.

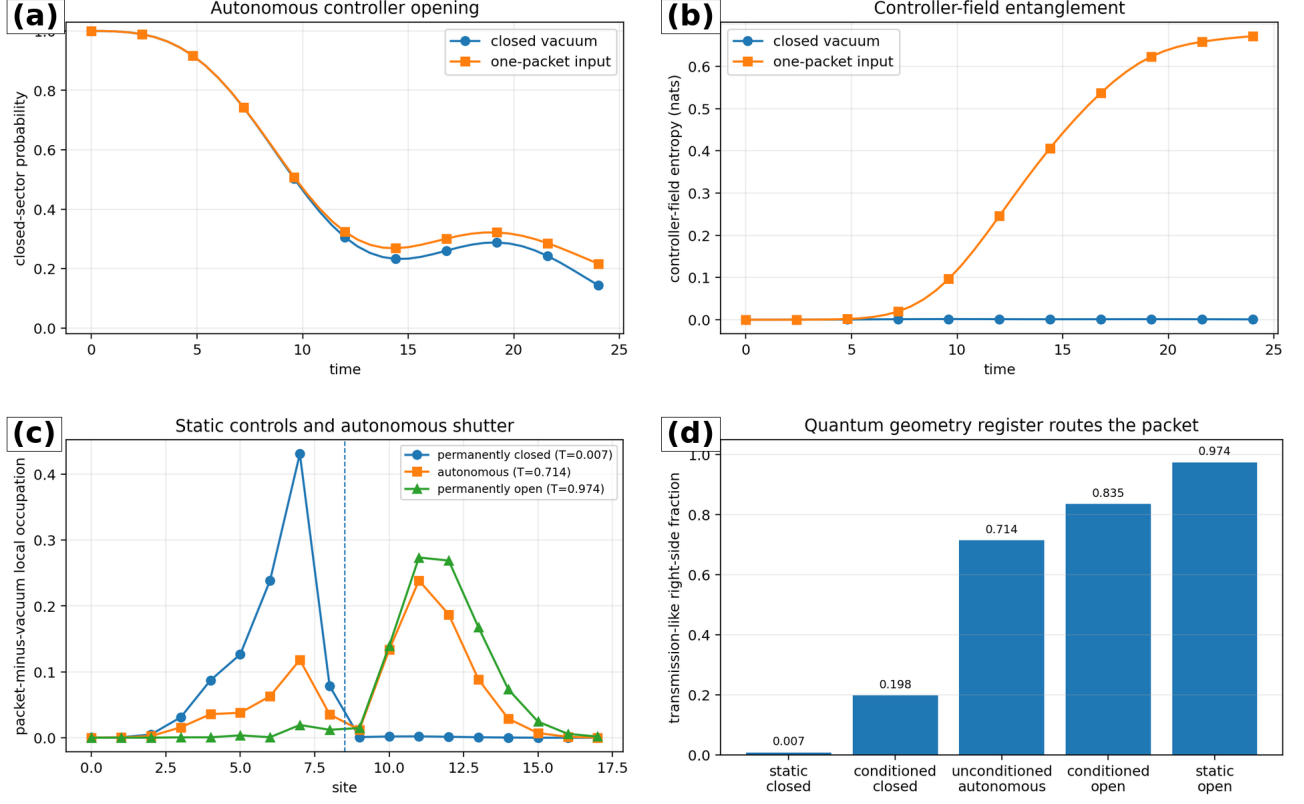


FIG. 2. Autonomous propagating-clock benchmark. (a) The controller leaves the closed sector under the fixed Hamiltonian; the incident packet changes its final closed probability. (b) The packet produces strong controller-field entanglement, whereas the vacuum does not. (c) Packet-minus-vacuum occupation for static closed, autonomous, and static open cases. (d) Conditioning the final field on the controller sector resolves reflected and transmitted branches.

The same run yields a right-half reduced number distribution

$$P_R(0) = 0.3011014, \quad P_R(1) = 0.6761192, \quad (31)$$

$$P_R(2) = 0.0065080, \quad P_R(3) = 0.0159481, \quad (32)$$

$$P_R(4) = 9.0307 \times 10^{-5}, \quad P_R(5) = 2.3290 \times 10^{-4}. \quad (33)$$

Every represented sector is positive and $P_R(n \geq 2) = 0.0227793$. This is the finite-cutoff precursor of an unbounded reduced distribution, not evidence for a literal infinite tail in the finite Hilbert space.

At fixed lattice spacing the transmission converges from 0.701053 at $N = 12$ to 0.714317 at $N = 18$. At $N = 14$, raising $n_{\max}$ from 5 to 6 changes $\mathcal{T}$ by 4.95×10^{-5} . A local-product calculation provides an independent regulator: in a six-site benchmark the global-sector and unrestricted local-product transmissions differ by 0.740%, while the excess open-mode number differs by 0.0783%. A unit-amplitude shift operator is a decisive negative control: its vacuum Bogoliubov-number relative error is 5.27, compared with 0.0177 for the square-root ladder in the same small benchmark.

V. NUMERICAL CONTROLS AND INDEPENDENT SOLVERS

Three representations are used. Small and medium systems are propagated by sparse Krylov evaluation of $e^{-iHt}|\psi\rangle$ [12]. Quadratic, untruncated reference dynamics are propagated symplectically. Long local-product chains are evolved as matrix product states using a second-order, center-moving canonical TEBD algorithm [13, 14].

The canonical gauge is essential. An earlier prototype truncated singular values outside a mixed-canonical form and produced nonmonotonic bond-dimension scans. The repaired algorithm moves the orthogonality center with each sweep, so every truncation acts on true Schmidt coefficients. At fixed local dimension $d = 3$ and $\Delta t = 0.2$, the transmission changes monotonically from 0.667742 at $\chi = 8$ to 0.667606 at $\chi = 32$, with accumulated discarded weight 2.34×10^{-9} . Canonical residuals are below 1.1×10^{-14} .

On an exact six-site, $d = 3$ benchmark without Schmidt truncation, the transmission error scales as $\Delta t^{2.0018}$ with $R^2 = 0.9999996$. For the 12-site benchmark, the $d = 5$ extrapolation $\mathcal{T}_{\Delta t \rightarrow 0} = 0.6691787$ differs from the independent global-occupation value 0.6690765 by 1.53×10^{-4} .

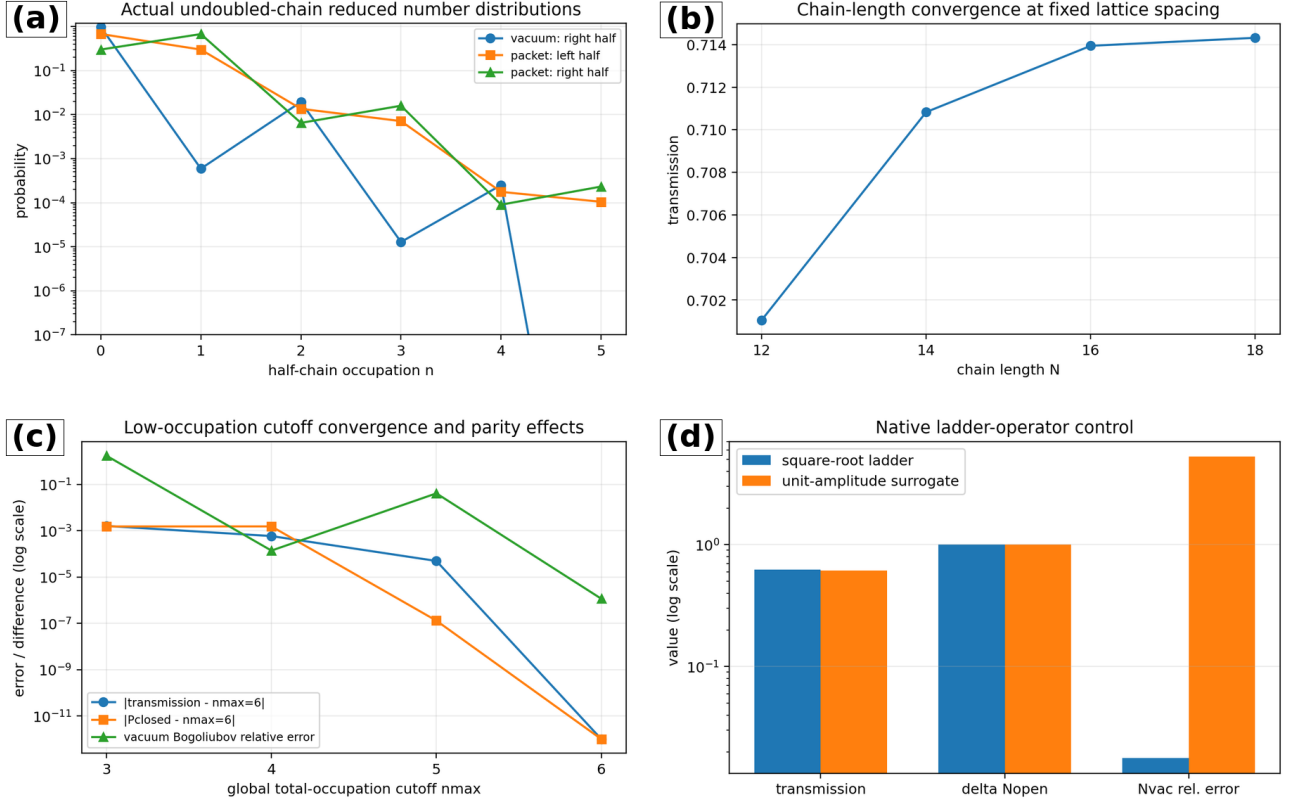


FIG. 3. Finite-regulator number structure and convergence. (a) Half-chain reduced number distributions in the actual undoubled autonomous model. (b) Transmission converges with chain length at fixed lattice spacing. (c) Transmission, controller probability, and vacuum Bogoliubov number converge with the global occupation cutoff; the squeezed-vacuum observable displays the expected parity effect.

relatively. In a separate $N = 8$, $d = 3$ full-product calculation, sparse evolution and canonical TEBD at $\Delta t = 0.05$ have packet fidelity 0.9999954 and transmission difference 1.31×10^{-3} ; halving the step reduces the latter by a factor 4.014.

Stable local-product calculations reach $N = 24$. At $N = 16$, $d = 4$, $\chi = 32$, and $\Delta t = 0.1$,

$$\mathcal{T}_{\text{MPS}} = 0.7785091, \quad \mathcal{T}_{\text{global}} = 0.7788985, \quad (34)$$

with accumulated discarded weight 1.74×10^{-9} . An exact entanglement budget on the 12-site reference has maximum entropy 0.696783 nats; $\chi = 33$ is sufficient to keep discarded Schmidt weight below 10^{-8} on every cut. These controls show that the observed effects are not artifacts of uncontrolled entanglement truncation.

VI. LOCAL SIMPLICITY AND GLOBAL DISTINGUISHABILITY

The smooth two-qubit controller in Eq. (11) provides a fixed reference profile through Eq. (12). We compare the autonomous field to the field-only reference in Eq. (13) without fitting the profile to the output.

For the complete right-terminal three-site algebra, the trace distance decreases as

$$4.84 \times 10^{-3}, \quad 7.09 \times 10^{-4}, \quad 1.37 \times 10^{-4}, \quad 7.95 \times 10^{-5} \quad (35)$$

for $N = 12, 16, 20, 24$. A permanently open reference gives 0.0473, 0.0108, 0.00593, and 0.00453, respectively. At $N = 24$, matching the boundary history improves the local comparison by a factor 57.

At $N = 16$, the largest three-site mismatch, 0.09768, is concentrated around the shutter. The terminal distance is 7.09×10^{-4} , while the global packet-reference fidelity 0.852066 implies, by the Fuchs–van de Graaf bounds [15],

$$D_{\text{global}} \geq 0.07693. \quad (36)$$

The global lower bound is more than 108 times the terminal-block distance. The discrepancy is stored mainly in the switching region, controller correlations, and non-local field correlations.

To examine outgoing degrees of freedom, we construct left and right positive-frequency packet modes. In one analysis the modes are fixed independently by propagating the input packet through the prescribed smooth symplectic history before evaluating the autonomous state. At

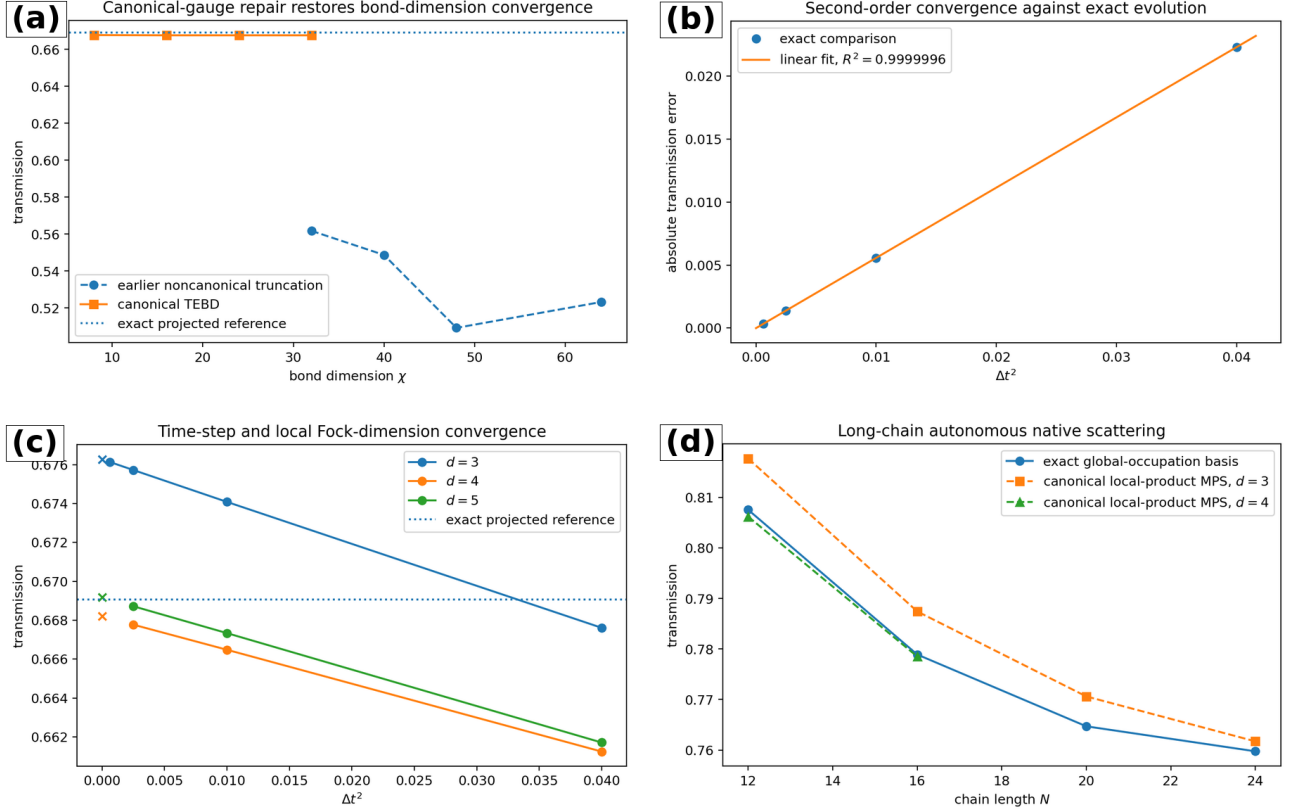


FIG. 4. Numerical validation. (a) Mixed-canonical truncation repairs the earlier nonmonotonic bond-dimension scan. (b) The symmetric TEBD transmission error is proportional to Δt^2 . (c) Joint time-step and local-dimension convergence approaches the independent global-occupation result. (d) A separate sparse-versus-TEBD calculation shows the same second-order scaling.

$N = 16$, $n_{\max} = 3$,

$$\begin{aligned} D_R &= 0.009474, & D_L &= 0.011174, \\ D_{RL} &= 0.030743. \end{aligned} \quad (37)$$

while the terminal three-site distance is 7.01×10^{-4} . A complementary regional-mode reconstruction, used to resolve number statistics, gives $D_R = 0.004023$, $D_L = 0.007483$, but $D_{RL} = 0.101910$. Tracing the left mode yields nonzero weight in every represented right-mode number sector, with $P_R(n \geq 2) = 1.53 \times 10^{-4}$ at $N = 16$.

A cutoff-free Gaussian long-lead calculation supplies a separate asymptotic check. With detector modes fixed independently of the autonomous output, the root-mean-square mismatch of local Weyl characteristic functions decreases to

$$\epsilon_R = 6.23 \times 10^{-6}, \quad \epsilon_L = 4.78 \times 10^{-6} \quad (38)$$

by $N = 224$. These are restricted-algebra statements. They do not assert equality of the full global states, and they are not yet the exact continuum one-photon/vacuum theorem of Ref. [11].

VII. QUANTUM-TO-CLASSICAL CONTROLLER LIMIT

A finite quantum controller is not equivalent to its mean switching profile. To quantify the classical limit, replace the single controller by the symmetric spin- $J = M/2$ sector of M two-level systems,

$$H_C = \Omega J_x, \quad P_C = \frac{J + J_z}{2J}, \quad (39)$$

initialized in $|J, J\rangle_z$. Its mean remains Eq. (12), whereas the connected fluctuation is

$$\langle \delta P_C(t) \delta P_C(t') \rangle = \frac{\sin(\Omega t) \sin(\Omega t')}{4M}. \quad (40)$$

Let $N[x]$ be the field particle number generated by the perturbed classical profile

$$p_x(t) = p_{\text{cl}}^{\text{free}}(t) + x \sin(\Omega t). \quad (41)$$

A spin-coherent expansion gives the leading profile-fluctuation contribution

$$\frac{\delta N}{N_0} = \frac{1}{8N_0} \left. \frac{d^2 N[x]}{dx^2} \right|_{x=0} \frac{1}{M} + O(M^{-3/2}). \quad (42)$$

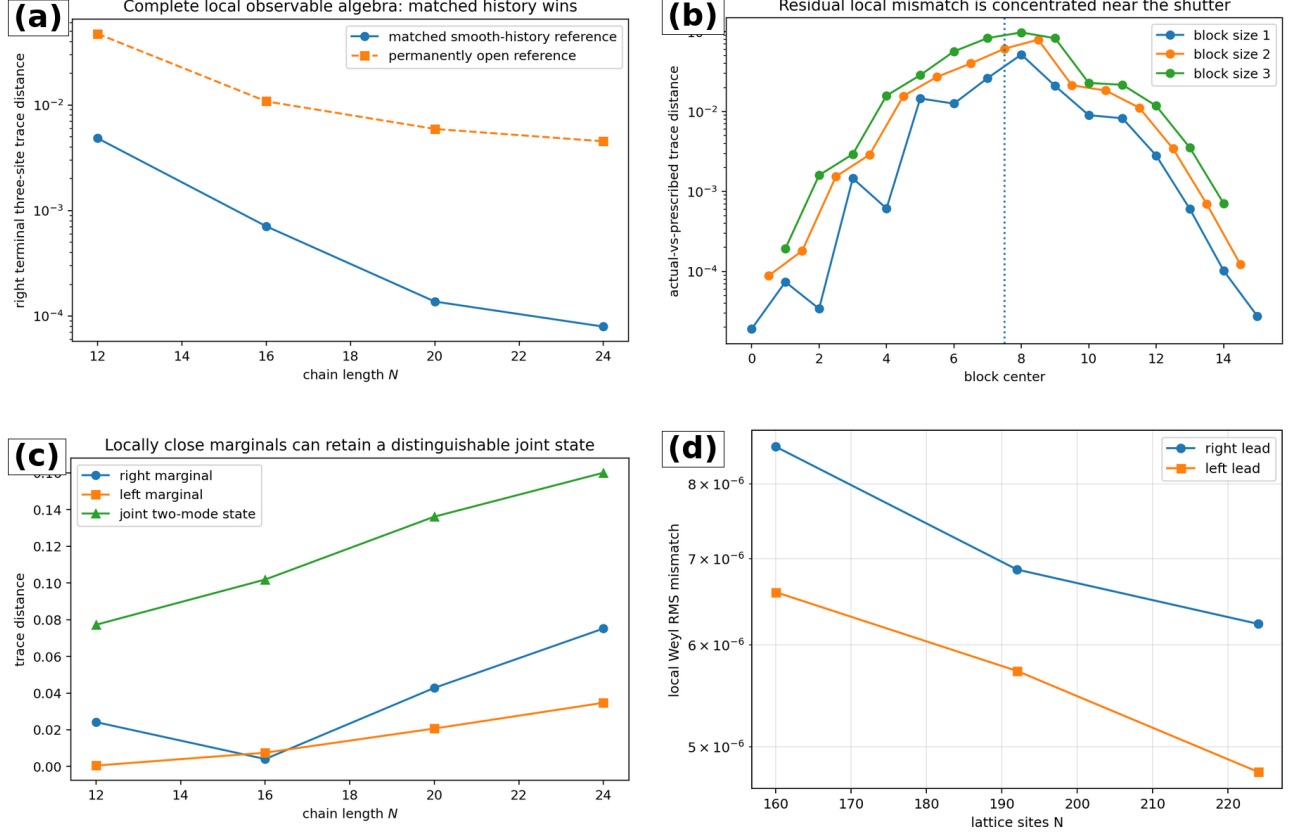


FIG. 5. Local versus global structure. (a) Complete terminal-block algebras approach the independently prescribed smooth-history reference much more closely than a permanently open reference. (b) The residual local mismatch is concentrated near the shutter. (c) Left and right outgoing marginals can be close while the joint two-mode state remains distinguishable. (d) In a separate unrestricted Gaussian long-lead calculation, fixed positive-frequency detector packets give local Weyl mismatches of order 10^{-6} .

The continuum-regulated curvature calculation yields coefficient $A_{\text{cont}} = 0.315339$. Exact native evolution at $N = 7$, $n_{\text{max}} = 10$ gives

$$\frac{N_{\text{auto}} - N_{\text{ref}}}{N_{\text{ref}}} = \frac{0.303678}{M}, \quad R^2 = 0.9999959 \quad (43)$$

for the asymptotic fit. The analytic curvature term accounts for 96.2% of the finite-regulator coefficient; the remainder measures coherent response and backreaction omitted by a classical profile mixture.

In the long-chain scattering calculation,

$$D_{RL}(M) = 0.10348M^{-0.9403}, \quad (44)$$

$$D_{\text{block3}}(M) = 7.42 \times 10^{-4}M^{-0.9685}. \quad (45)$$

with $R^2 > 0.9994$ for $M \geq 2$. A two-packet connected response falls from 0.1102 above the prescribed-switch floor at $M = 1$ to 3.25×10^{-5} at $M = 16$. The mean classical switch is therefore recovered in a controlled collective limit, whereas a single quantum shutter retains branch fluctuations and memory.

VIII. EXACT EMBEDDING IN ORDINARY MATTER FRUSTRATION

The controller need not be an abstract register. Introduce one reference matter qubit and R clock qubits. In the unary sector with the reference unexcited and exactly one clock excitation, define $W|s\rangle = |s\rangle_{\text{unary}}$. The local XY Hamiltonian

$$H_m = -\frac{J}{2} \sum_{s=0}^{R-2} (X_{s+1}X_{s+2} + Y_{s+1}Y_{s+2}) \quad (46)$$

obeys

$$W^\dagger H_m W = -J \sum_s (|s+1\rangle\langle s| + |s\rangle\langle s+1|). \quad (47)$$

With frustration projectors

$$F_{0,s+1} = \frac{\mathbb{I} - Z_0 Z_{s+1}}{2}, \quad (48)$$

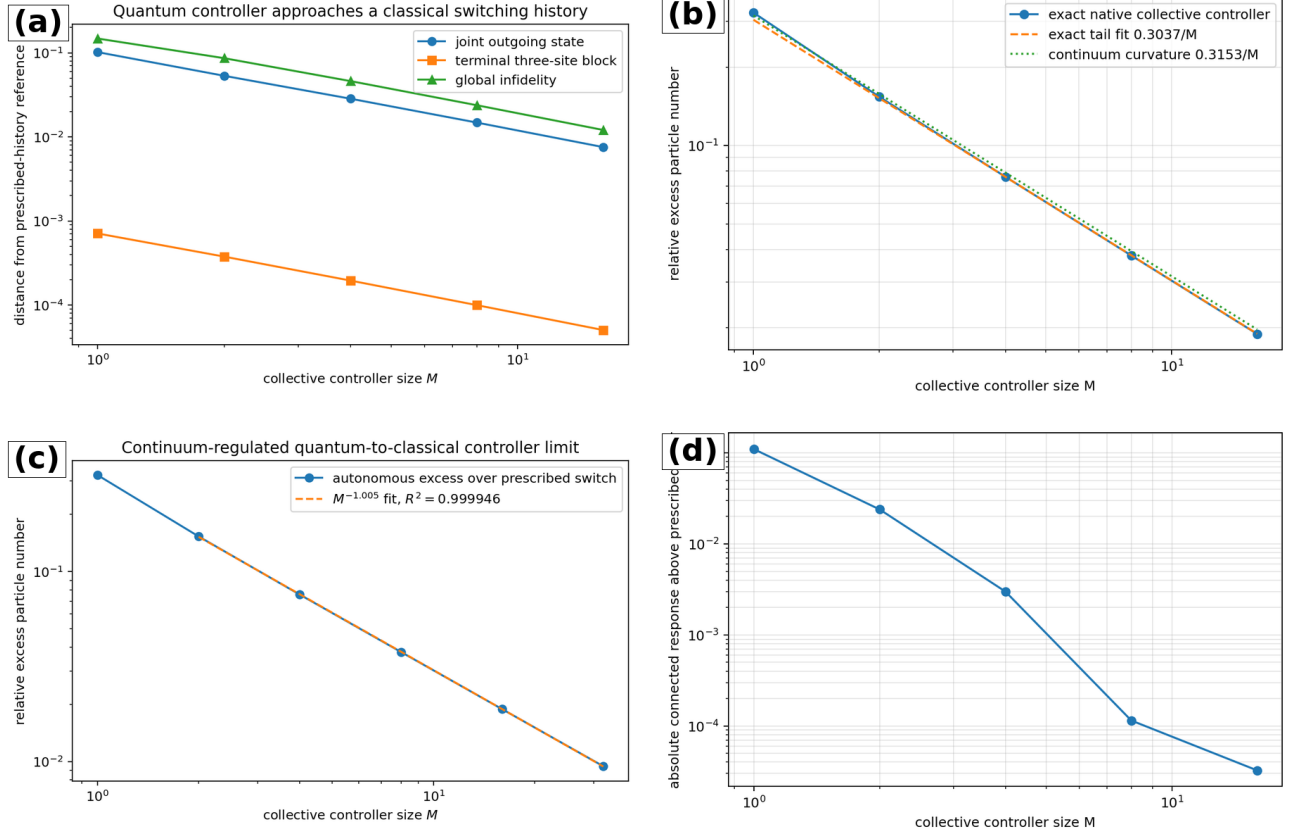


FIG. 6. Quantum-to-classical controller limit. (a) Complete local and joint outgoing distances decrease with collective controller size. (b) Excess vacuum excitation above the prescribed switch scales approximately as $1/M$. (c) The exact native coefficient is close to the independently calculated spin-coherent curvature coefficient. (d) A two-packet connected response, generated by controller memory and backaction, disappears toward the prescribed-switch floor as M increases.

we have

$$W^\dagger \left(\sum_{s < s_*} F_{0,s+1} \right) W = P_C. \quad (49)$$

Thus the unary matter sector is exactly the controller in Eq. (10). Appendix B gives the proof.

Full-Hilbert-space evolution confirms the isometry. With no leakage, the matter-clock and abstract-register transmissions differ by 1.35×10^{-14} and their excess one-body matrices by 7.05×10^{-14} . Generic transverse leakage of strength 0.05 reduces the physical-sector weight to 0.4639 without protection. Adding a penalty $U(N_{\text{clk}} - 1)^2$ with $U = 4$ raises it to 0.997795, with transmission error 0.003278 and one-body error 0.007243 relative to the ideal register.

This is the direct matter-clock interpretation of the benchmark: ordinary matter frustration controls a bilinear between square-root edge modes. The broader claim that such an architecture is the microscopic origin of an electromagnetic field is neither assumed nor established here.

IX. DISCUSSION

The calculations isolate what is and is not new. The mode-mixing physics in Eqs. (19)–(22) is standard dynamical-Casimir physics. The new element is that the boundary condition is generated by a finite quantum subsystem within the same time-independent Hamiltonian as the field. This gives three effects absent from a prescribed classical switch.

First, the field and controller exchange energy and become entangled. The final controller state depends on whether an incident packet was present, and conditioning on it separates reflected and transmitted branches. Second, matching the mean switching profile is sufficient for very accurate far-local observables but not for the joint field state: controller branch information survives in nonlocal correlations. Third, controller fluctuations generate additional radiation and multipacket memory that vanish in a controlled collective limit.

The model also clarifies several scope boundaries. The edge ladders are scalar bosonic modes. No Gauss law, transverse polarization structure, or emergent electromag-

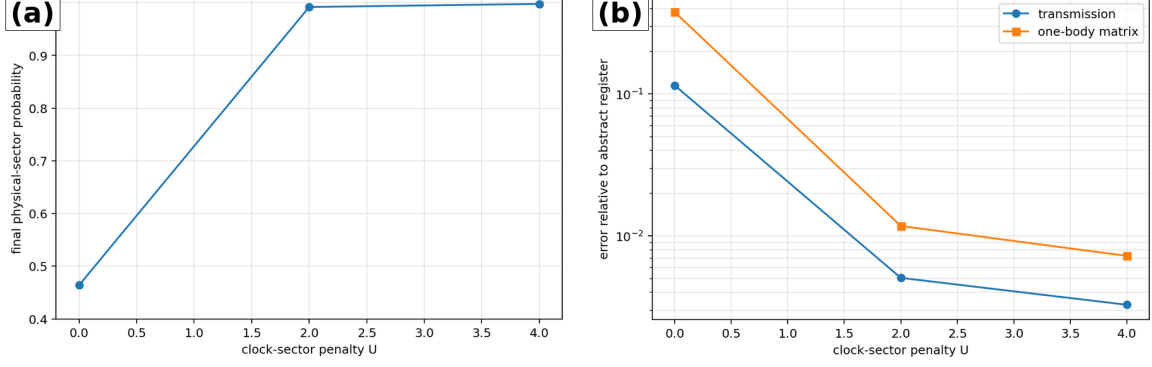


FIG. 7. Matter-clock embedding and leakage protection. (a) The unary physical-sector probability is restored by a number-sector penalty. (b) Transmission and one-body errors relative to the abstract register decrease correspondingly.

netic $U(1)$ is used. Every finite ladder or global occupation sector has a finite number cutoff, so the positive reduced probabilities reported here establish support on every represented sector only. The ideal finite-box divergence in Eq. (25) belongs to an abrupt perfect-cut regulator and is not a normalized outcome with literally infinitely many quanta. Finally, the long-lead Weyl tests and fixed outgoing packets support a finite-regulator local/global separation, but they do not replace a full continuum scattering theorem.

These limitations are productive rather than cosmetic: they separate a closed finite-model claim from the stronger gauge-emergence and continuum claims that require independent papers. The present result is therefore best viewed as an autonomous quantum-boundary benchmark for a matter–edge lattice architecture.

X. CONCLUSION

We have constructed and tested a quantum shutter whose actuator is part of the modeled quantum system. One local matter-controlled bilinear changes the field’s vacuum structure, produces a nonzero Bogoliubov β block, and converts a one-quantum state into a one-quantum-added squeezed state. Exact and tensor-network evolutions show autonomous scattering, backreaction, controller–field entanglement, reduced many-number support, and conserved total energy. Complete far-local algebras can follow the mean boundary history even when the joint outgoing state remains distinguishable, and the discrepancy from a prescribed switch vanishes approximately as $1/M$ for a collective controller. The controller itself has an exact realization in ordinary matter qubits and frustration projectors.

The defensible result is accordingly precise: a finite matter–edge Hamiltonian can autonomously realize the standard boundary-change Bogoliubov mechanism with backreaction and nontrivial local/global state structure. Electromagnetic gauge emergence and an unregulated

continuum limit lie beyond that statement.

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DATA AND CODE AVAILABILITY

All source code, raw numerical data, figure-generation scripts, parameter files, and SHA-256 checksums required to reproduce the reported benchmarks are included in the accompanying reproducibility archive. The submission bundle also includes the exact TeX source and figures used to build this manuscript. A public archival DOI can be added before journal submission.

Appendix A: Finite-box coefficients and ultraviolet asymptotics

Consider a continuum scalar field on $0 < x < L$. After opening, the normalized modes and frequencies are

$$v_m(x) = \sqrt{\frac{2}{L}} \sin \frac{m\pi x}{L}, \quad \Omega_m = \frac{cm\pi}{L}. \quad (\text{A1})$$

Before opening, a normalized left-half mode is

$$u_n^L(x) = \sqrt{\frac{4}{L}} \sin \frac{2n\pi x}{L}, \quad 0 < x < L/2, \quad \omega_n = \frac{2cn\pi}{L}. \quad (\text{A2})$$

The overlap is

$$S_{mn}^L = \begin{cases} 1/\sqrt{2}, & m = 2n, \\ 0, & m \text{ even}, m \neq 2n, \\ (-1)^{n+(m-1)/2} \frac{4\sqrt{2}n}{\pi(m^2 - 4n^2)}, & m \text{ odd}. \end{cases} \quad (\text{A3})$$

Substitution into Eq. (20) yields Eq. (24); the factor $m - 2n$ cancels for odd m . In the degenerate channel $m = 2n$, the frequency ratio is one and $\beta = 0$ despite nonzero overlap.

The two half-cavities contribute equally, giving

$$N_{\text{vac}}(K) = \frac{8}{\pi^2} \sum_{\substack{m \leq K \text{ odd} \\ n \leq K}} \frac{n}{m(m+2n)^2}. \quad (\text{A4})$$

For $1 \ll m, n \lesssim K$, the inner sum has a logarithmic contribution and summation over odd m produces a second logarithm. Euler–Maclaurin replacement of the bulk domain gives

$$N_{\text{vac}}(K) = \frac{1}{2\pi^2} (\log K)^2 + O(\log K). \quad (\text{A5})$$

A numerical fit over $K \in [50, 6400]$ gives coefficient 0.050947, compared with $1/(2\pi^2) = 0.050661$.

Appendix B: Unary matter-clock isometry

Let the reference qubit be site 0 and the clock qubits sites $1, \dots, R$. Define number operators $n_j = (\mathbb{I} - Z_j)/2$ and the unary subspace

$$\mathcal{H}_1 = \text{span}\{|s\rangle_{\text{unary}} : s = 0, \dots, R-1\}, \quad (\text{B1})$$

where the reference is in $Z = +1$ and exactly clock qubit $s+1$ is excited. The XY identity

$$\frac{1}{2}(X_j X_{j+1} + Y_j Y_{j+1}) = \sigma_j^+ \sigma_{j+1}^- + \sigma_j^- \sigma_{j+1}^+ \quad (\text{B2})$$

shows that Eq. (46) preserves $\mathcal{H}_1$ and acts there as nearest-neighbor hopping. Moreover, on $\mathcal{H}_1$, $F_{0,s+1} = n_{s+1}$ because the reference is unexcited. Therefore Eqs. (10) and

(49) follow exactly under the isometry W . A penalty

$$H_U = U \left(\sum_{s=1}^R n_s - \mathbb{I} \right)^2 + h_{\text{ref}} n_0 \quad (\text{B3})$$

protects the unary sector against generic leakage terms.

Appendix C: Numerical regulators and diagnostics

The exact global-occupation basis contains all N -tuples satisfying $\sum_j n_j \leq n_{\text{max}}$ and has dimension

$$D_F = \binom{N + n_{\text{max}}}{n_{\text{max}}}. \quad (\text{C1})$$

It preserves the square-root ladder matrix elements exactly within the retained sector but imposes a global ceiling. The independent local-product regulator uses d states on each site and is propagated either exactly for small N or as an MPS for larger N . Convergence requires both low top-rung probability and stability of observables under n_{max} or d .

The canonical TEBD step is a symmetric second-order product formula. The orthogonality center is moved with the sweep so that each SVD truncates the actual Schmidt spectrum. Reported runs monitor the accumulated discarded weight, left/right canonical residuals, norm, total energy, local occupation tails, and comparison with exact sparse evolution. Quadratic Gaussian calculations use the real symplectic equations for the mode functions and independently verify the canonical identities in Eq. (21).

For reduced outgoing modes, finite density matrices are reconstructed from normal-ordered moments. Direct one-mode reconstruction and the partial trace of the joint two-mode reconstruction agree to 4.26×10^{-6} at $N = 16$ and 1.18×10^{-6} at $N = 24$. Weyl-characteristic comparisons use randomly sampled local phase-space displacements fixed independently of the autonomous output.

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- [1] G. T. Moore, “Quantum theory of the electromagnetic field in a variable-length one-dimensional cavity,” *J. Math. Phys.* **11**, 2679 (1970).
 - [2] S. A. Fulling and P. C. W. Davies, “Radiation from a moving mirror in two dimensional space-time: conformal anomaly,” *Proc. R. Soc. Lond. A* **348**, 393 (1976).
 - [3] N. D. Birrell and P. C. W. Davies, *Quantum Fields in Curved Space* (Cambridge University Press, Cambridge, 1982).
 - [4] C. K. Law, “Effective Hamiltonian for the radiation in a cavity with a moving mirror and a time-varying dielectric medium,” *Phys. Rev. A* **49**, 433 (1994).
 - [5] V. V. Dodonov, “Current status of the dynamical Casimir effect,” *Phys. Scr.* **82**, 038105 (2010).
 - [6] P. D. Nation, J. R. Johansson, M. P. Blencowe, and F. Nori, “Colloquium: Stimulating uncertainty: Amplifying the quantum vacuum with superconducting circuits,” *Rev. Mod. Phys.* **84**, 1 (2012).
 - [7] C. M. Wilson *et al.*, “Observation of the dynamical Casimir effect in a superconducting circuit,” *Nature* **479**, 376 (2011).
 - [8] S. Butera and I. Carusotto, “Mechanical backreaction effect of the dynamical Casimir emission,” arXiv:1810.11281.
 - [9] S. Butera and I. Carusotto, “Quantum fluctuations of the friction force induced by the dynamical Casimir emission,” arXiv:1909.09607.
 - [10] Y.-C. Xie, S. Butera, and B.-L. Hu, “Optomechanical backreaction of quantum field processes in dynamical Casimir effect,” arXiv:2308.03129.
 - [11] I. C. O. Rukan, J. Gulla, and J. Skaar, “A truncated photon,” arXiv:2510.21636.

- [12] A. H. Al-Mohy and N. J. Higham, “Computing the action of the matrix exponential, with an application to exponential integrators,” *SIAM J. Sci. Comput.* **33**, 488 (2011).
- [13] G. Vidal, “Efficient simulation of one-dimensional quantum many-body systems,” *Phys. Rev. Lett.* **93**, 040502 (2004).
- [14] U. Schollwöck, “The density-matrix renormalization group in the age of matrix product states,” *Ann. Phys.* **326**, 96 (2011).
- [15] C. A. Fuchs and J. van de Graaf, “Cryptographic distinguishability measures for quantum-mechanical states,” *IEEE Trans. Inf. Theory* **45**, 1216 (1999).