

Released records: nondemolition relative to the accessible future as a criterion for quantum events

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Abstract

Anchored Causality Theory holds that a quantum system is never a particle: it is one wave, a quantum field with internal structure, from start to finish, and the particle is what happens, the mark a detector makes when the bath around it carries a record beyond recovery. The main accounts of quantum mechanics keep a particle somewhere, at the detector or on a hidden path, or they keep every outcome, each in a world of its own; ACT does neither. This paper says when a record has passed beyond recovery—when it is *released*—and what law the marks then follow. A continuously monitored quantum system emits an output whose observables commute with one another and with every later system observable; Belavkin’s nondemolition principle identifies these output observables as a classical past. We take that structure as the carrier of actuality and ask what more is needed for a single, exclusive, causally ordered history of events. Three results follow. First, exclusivity of outcomes is a theorem: the record algebra is commutative for arbitrary system Hamiltonian and coupling, its spectrum is the set of histories, and the joint spectral projectors are orthogonal with Born probabilities equal to the instrument path measure. Second, we define *release* of a record as commutation with the system’s physically accessible future — later system observables and fresh records, through every return path the interaction network provides — and prove a bound tying this algebraic condition to the recoverable-phase functional of the companion paper: for a binary record with gap Δ , the phase distinguishability accessible through any observable A is at most $\|\omega_{\text{off}}\|_1 \|[R, A]\| / \Delta$. Release is exactly the condition under which attaching actuality to a record alters no accessible prediction for any state. Attached earlier, by a rule of the kind treated here — the actual history proceeds from the record, the unselected alternatives no longer enter that history, and the global wave is untouched — actuality changes what the actual history predicts, and the recoherence experiments exclude each such earlier placement that has been tested. The weights are not imported: an appendix shows that six stated premises — one wave with one covariant law of chance, one actual history, no record where the wave has no component, no signalling for every division of the network, physically realisable refinement of records, and continuity — force the chance of a released record to be its squared amplitude, by Zurek’s argument in Barnum’s no-signalling form. Third, outside the free-output setting the criterion does work that nondemolition alone does not: under time-delayed return the in-flight record is nondemolition but not released, and anchoring is deferred to the exit record; under intra-environment scrambling a record leaves the absolute centre of the future algebra while remaining central relative to the system. The event law is the instrument path measure, with the companion paper’s hazard $h = \Phi_{\text{irr}}$ as its nondemolition corollary and the diffusive stochastic master equation as the large-local-oscillator limit of a counting history; its records carry no particle labels, so it treats identical particles with nothing added. The framework adds one primitive commitment — one point of the released-record spectrum is actual — and no new field, constant, threshold or noise. Every numerical statement is verified by a deposited script; the results on records and release are exact in a finite collision model.

Keywords: quantum filtering; nondemolition measurement; quantum trajectories; measurement problem; records; decoherence; quantum Darwinism.

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1 Introduction

1.1 The thesis

For a century the same wave has been asked to be two things: an extended field that interferes, and a localised particle that leaves a mark. Every interpretation of quantum mechanics is a proposal for how to live with that request. The programme developed here, and in the companion paper [1], answers it by refusing it. The main accounts keep a particle somewhere, at the detector or on a hidden path, or they keep every outcome, each in a world of its own. This programme does neither. The system is one composite quantum field before a record forms and after: extended and interfering before, still a field after. Nothing about it is particle-like except

the quantum numbers it carries. When an environment, a detector’s bath of light and vibrating atoms, writes an irreversible record of one alternative, an *event* occurs: one localised, dated fact enters the world’s causal history. The wave does not turn into a particle; the field produces a record, and the record is what the word *particle* has always named: a click, a grain, a droplet, a track [56]. Duality is not resolved but dissolved: the two descriptions were never descriptions of the same object at the same time. What remains is a realist picture with an ordinary causal order — fields evolve unitarily, detectors write records through ordinary couplings, and the history of the world is the history of the records that can never be unwritten.

The companion paper built this picture for matter-wave interferometry, where the record-writing environment is a quantum Brownian motion bath and the event rate is read from the bath’s decoherence exponent. Its derivations were confined, deliberately, to quantum-nondemolition (QND) record formation in CP-divisible regimes, and it listed what the picture still lacked: a general event law outside the QND window, a statement of the diffusive (homodyne) case, a criterion for when a record is beyond recall that does not depend on an observer’s choices, and a formulation not tied to a preferred time. The present paper supplies the first three and states precisely what the fourth still needs.

1.2 What is done here

The mathematical object at the centre of the paper is old. In the quantum stochastic calculus of Hudson and Parthasarathy [3] and the input–output theory of Gardiner and Collett [4], a system coupled to a Bose field emits an output whose observables Y_t commute among themselves and with every later system observable. Belavkin [5, 6, 7, 8, 9, 10] called the first property self-nondemolition and the second nondemolition, showed that the conditional expectation onto the output algebra is the quantum filter, and proposed — under the name *eventum mechanics* — that the output record is a classical past: a commutative algebra of beables in the centre of the algebra describing the future, with “everything in the future constituted by quantized waves, everything in the past by trajectories of the recorded particles.” In the present programme the past is a history of marks, with no particle paths, and the wave remains a wave throughout. Bouten, van Handel and James [11, 12, 13] gave the modern proofs and a finite-dimensional “discrete invitation” that we use for every computation.

We attach the actuality of the companion paper to this algebra and prove what follows. Section 3 shows that the record algebra is commutative and nondemolition for *arbitrary* system Hamiltonian and coupling (Theorem 1), that its Gelfand spectrum is the set of histories with orthogonal projectors and path-measure probabilities (Theorem 2) — so that the exclusivity of outcomes, stated as a commitment in [1], is a theorem here — and that the induced *system* effects commute when the system Hamiltonian preserves the recorded observable, the converse being false (Theorem 3). Section 4 defines release (Definition 3) as commutation with the system’s accessible future, proves that in the free-output setting every readout is released so that release is exactly nondemolition and selects nothing (Theorem 5), proves the bound relating the commutator to the companion paper’s recoverable-phase functional (Theorem 6) together with the failure of its converse and the blindness of the commutator to the carrier’s later history (Propositions 1–2), and then treats the two cases in which the criterion does work that nondemolition alone does not: time-delayed return (Proposition 3) and intra-environment scrambling (Proposition 4). Corollary 2 draws the consequence that matters for the ontology: release is exactly the point at which attaching actuality to a record changes no accessible prediction, so every earlier placement of an event by a rule of the kind the corollary treats — one under which the actual history proceeds from the record, the unselected alternatives no longer enter that history, and the global wave is untouched — changes what the actual history predicts. Section 5 states the event law as the instrument path measure (Theorem 7), recovers the QND hazard of [1] as its corollary, derives the diffusive limit (Theorem 8), and shows that its records carry no particle labels, so that it needs nothing added for identical particles (Proposition 5). Section 6

reads the criterion against the laboratory: the recoherence experiments fix release as the earliest admissible placement of an event, and four families of experiments corroborate the ingredients the criterion uses. Section 7 places the record-geometry functional of [1] on the readability side of the release/readability distinction. Section 8 states the single primitive commitment. Section 9 records what Fewster–Verch measurement theory [45, 46, 47] supplies toward a covariant statement and what it does not. Section 10 lists what is open. Appendix A shows that the Born weights the path measure carries are forced by six stated premises (Theorem 9), so that the descent of those weights from the ensemble in [1] becomes a consistency check.

The results on records and release, and the path measure of the event law, are computed exactly in a collision model with five two-level field slices (a 64-dimensional space), with the system Hamiltonian chosen *not* to commute with the coupling; the diffusive limit is checked against sampled histories, and the Born witnesses of Appendix A and the identical-particle example of Section 5 use their own finite models. Each numerical statement carries a bracketed label pointing to the check in the accompanying scripts (Appendix B); the five scripts hold 89 checks, and all pass.

1.3 Prior art, and the exact delta

The existence and classicality of the record algebra are Belavkin’s, and nothing in Section 3 that concerns them is claimed as new. The contribution of this paper relative to eventum mechanics lies in four points.

1. *Release relative to the accessible future.* Belavkin’s beables are output observables of a free field, central in the algebra of all future observables. We define release relative to the *system’s physically accessible* future — the observables and fresh records that the actual interaction network lets act on the system — and prove that in the free-output setting this reduces exactly to nondemolition (Theorem 5). The two notions part company when the output is not free: when carriers return (Proposition 3) or when the environment evolves internally (Proposition 4). In the second case a record leaves the absolute centre while the system’s reduced state, its entire future statistics and its instrument are unchanged; a criterion that then denies the event ties eventhood to the environment’s internal bookkeeping, and we argue that the accessible-centre criterion gives the physically correct verdict.
2. *A graded form with a proved bound.* Belavkin’s framework never needed a graded notion of nondemolition. Ours does, because physical return paths are partial, and Theorem 6 supplies the inequality that ties the algebraic condition to the recoverable-phase functional of [1].
3. *Scope as an existence statement.* The class in which the ontology is defined is the class in which a released record exists; what lies outside (recurrent baths, erasers, carriers that never exit) is an object of the theory, not an unexamined boundary.
4. *Release is a necessary condition; the detector present fixes the readout.* In the free-output setting every readout of the output — photon counting, any quadrature — is released, and readouts of the same carrier along non-commuting axes do not commute (Theorem 5). Release is therefore a necessary condition on a record, not a selection of *which* record is actual. That selection is fixed by the accessible algebra, i.e. by the detector physically present, exactly as Wiseman and Gambetta showed that the unravelling is fixed by the detection scheme [17]. Section 10 says what the apparatus-free case still lacks.

Three further relatives should be named. Kent’s late-time photodetection beables [20] and his earlier one-world proposal [19] share the commitment to a single actual configuration read from asymptotic records; the present account differs in *when* (events as they are released, not

a final-time draw) and in retaining the global state as the dynamical carrier. Gambetta and Wiseman’s modal dynamics for positive-operator measurements [18] is the nearest interpretive relative on the trajectory side. Steane’s condition that a single-world reading is sustainable “provided certain outcomes are not quantum-erased in the future” [21], Gell-Mann and Hartle’s permanent records [22] and Riedel’s decompositions defined by records that stay redundant [24] all carry the same reference to the future that our release condition carries; Section 4.6 argues that this reference is a statement about the interaction network and is inherited by every record-based one-history account.

Other event ontologies. The reading of a quantum event as the passage from possibility to fact has a lineage that the present paper joins rather than founds: Heisenberg’s *potentia* [60], Haag’s events as fundamentally irreversible transitions inside a growing causal net [61, 62], the event-enhanced quantum theory of Blanchard and Jadczyk [63], Kastner’s transactional actualisation [64], and, closest in kind, the ETH approach of Fröhlich and collaborators [65, 66, 67, 68]. The ETH approach defines an event through the algebra of observables of the *future* — an event occurs when the centralizer of the state on that algebra acquires a non-trivial centre — grounds the loss of access in massless radiation escaping to infinity by way of Huygens’ principle [69], and has recently derived the quantum jumps of resonance fluorescence from emission into the quantized field [68]. The present criterion is of the same type and should be judged against it. The differences are these. ETH replaces the Schrödinger evolution of the state by a stochastic evolution driven by its events; here the global dynamics is never modified, actuality is attached to a point of the record spectrum, and the conditioned state is the Belavkin filter, so that the ensemble is exactly the unitary one for every initial state. ETH states its criterion on the full algebra of the future; release is stated on the *accessible* future, and Section 4.5 exhibits the case that separates the two (a record leaves the absolute centre while nothing accessible to the system has changed). ETH’s events are fixed by the centre of the state’s centralizer and carry no rate; here the rate is the instrument intensity, read off the coupling. And the selection of what can be recorded is here a theorem about record architecture (the selection theorem of [1]) rather than a property of the state. Zhang’s *event anchoring* [70] — an event defined by local classicalization, redundant environmental recording and irreversibility against recovery, with a globally unitary dynamics and a history that is a partial order of frozen events — is the nearest relative in vocabulary and in gate; it leaves single-outcome event generation open as a separate problem, which is what Sections 5 and 8 supply, and it requires redundancy where release does not (Proposition 4: the event stands as redundancy falls). The relativistic flash ontology [55] shares the event primitive and the absence of a preferred frame, at the price of a modified dynamics that the radiation searches now constrain. Finally, the step at which the coupling fixes which alternatives a record can register is einselection [71, 72, 26]; what is added to it here is the vocabulary of components and alternatives, and the selection theorem.

1.4 Relation to the companion paper

The companion paper [1] states the ontology, works the quantum Brownian motion case in two exactly solvable models, and proves the selection theorem; the present paper supplies the algebra. Four of its statements receive their general form here. (i) It posits an exclusive outcome register and one actual record history, and states that the first is a theorem in the record-algebra class; Theorem 2 is that theorem, and the framework has one commitment (Section 8). (ii) Its event law $h = \Phi_{\text{irr}}$ is scoped to QND record formation; here it is the corollary of the instrument path measure (Corollary 3) and the scope is explained rather than assumed. (iii) Its recoverable-phase functional Φ_{irr} (Definition 1 there) is shown to be the graded form of release, bounded by the commutator (Theorem 6); the record-geometry functional built on it belongs to the readability side (Section 7). (iv) Its Born weights are forced in the binary QND model by ensemble agreement with an outcome-neutral first-event process (its Proposition 2)

and otherwise proposed; Appendix A forces them in general from six stated premises, and the model's descent becomes the check that the single-history law and the master equation agree.

2 Setting

2.1 The continuous case (reference)

A system S with Hilbert space $\mathcal{H}_S$ is coupled to a Bose field in the Hudson–Parthasarathy framework [3]: a unitary U_t on $\mathcal{H}_S \otimes \Gamma$ (Γ the symmetric Fock space over $L^2(\mathbb{R}_+)$) solves the quantum stochastic differential equation with system Hamiltonian H_S and coupling operator L . Heisenberg system observables are $j_t(X) = U_t^\dagger(X \otimes I)U_t$; the output field is $b_{\text{out}} = b_{\text{in}} + L$ [4]; an observed record is $Y_t = U_t^\dagger(I \otimes Q_t)U_t$ with Q_t an input-field quadrature (homodyne) or counting process. Bouten and van Handel [11] and Bouten, van Handel and James [12] prove self-nondemolition, $[Y_t, Y_s] = 0$ for all s, t , and nondemolition, $[j_t(X), Y_s] = 0$ for $s \leq t$; hence $\mathcal{R}_t := \text{vN}\{Y_s : s \leq t\}$ is commutative and, by the spectral theorem, isomorphic to $L^\infty(\Omega_t, \mathcal{F}_t, \mathbb{P})$; the conditional expectation onto $\mathcal{R}_t$ exists and satisfies the quantum filtering (Belavkin) equation. None of this depends on $[H_S, L] = 0$. Bouten and van Handel remark that a process failing self-nondemolition “cannot be observed in a single realization of an experiment and is physically meaningless” [11]; that remark is the seed of the present paper.

2.2 The discrete realisation (used for all computations)

Following [13], a system qubit interacts sequentially with fresh two-level field slices $a_1, a_2, \dots, a_N$. With $|e\rangle, |g\rangle$ the system levels, $\sigma_- = |g\rangle\langle e|$, and each slice initially in its ground state $|g\rangle_k$ (the vacuum), step k is

$$U_k = e^{-iH_S \delta t} e^{-i\theta_c(\sigma_- \otimes a_k^\dagger + \sigma_+ \otimes a_k)}, \quad \theta_c = \sqrt{\gamma \delta t}, \quad (1)$$

where $a_k^\dagger = |e\rangle\langle g|_k$ excites slice k . The chain up to step k is $U_{1..k} = U_k U_{k-1} \cdots U_1$. The record on slice k is $Q_k = \sigma_z^{(k)}$ (counting: is a photon present in slice k ?) or $\sigma_x^{(k)}$ (a quadrature). Unless stated otherwise the computations use $N = 5$ slices, $\gamma = 1$, $\delta t = 0.25$, and

$$H_S = \frac{1}{2}\Omega \sigma_x, \quad \Omega = 1.3, \quad (2)$$

which does *not* commute with the coupling: the system is driven between the levels the record distinguishes, so the setting is non-QND throughout. All commutators are computed exactly on the 2^{N+1} -dimensional space.

2.3 Definitions

Definition 1 (Record observables, record algebra). $Y_k := U_{1..k}^\dagger Q_k U_{1..k}$ is the record observable of slice k (Heisenberg picture at the end of step k ; Theorem 1(iii) shows it is unchanged by later steps). $\mathcal{R}_N := \text{vN}\{Y_1, \dots, Y_N\}$ is the record algebra.

Definition 2 (Accessible algebra; future accessible algebra). The *accessible algebra* $\mathcal{A}_{\text{acc}}$ is the algebra of system and environment observables that the apparatus can physically act on or read; it is a specification of the setup, on the same footing as the detection scheme in quantum-trajectory theory (Section 4.7). For a time s , the *system-accessible future* is

$$\mathfrak{F}_s^S := \text{vN}\{j_t(A), Y_t : t \geq s, A \in B(\mathcal{H}_S)\}, \quad (3)$$

the algebra generated by all later Heisenberg system observables and all later fresh records, computed with the *actual* total dynamics of the network, including any path by which a carrier returns to the system. The *full accessible future* $\mathfrak{F}_s$ is the same with $B(\mathcal{H}_S)$ replaced by $\mathcal{A}_{\text{acc}}$, so that it also contains later readouts of carriers already emitted.

The companion paper's Definition 1 is used unchanged: for alternatives $|j\rangle, |k\rangle$ and $|\psi_\theta\rangle = \frac{1}{\sqrt{2}}(|j\rangle + e^{i\theta}|k\rangle)$, with $\rho_\theta(t)$ the global state evolved from $|\psi_\theta\rangle\langle\psi_\theta| \otimes \rho_E$ and $\mathcal{E}_{\text{acc}}$ the restriction to $\mathcal{A}_{\text{acc}}$,

$$\mathcal{D}(t) = \frac{1}{2} \|\mathcal{E}_{\text{acc}}[\rho_0(t)] - \mathcal{E}_{\text{acc}}[\rho_\pi(t)]\|_1, \quad \Phi_{\text{irr}}(t) = -\ln \frac{\mathcal{D}(t)}{\mathcal{D}(0)}. \quad (4)$$

Theorem 4 below shows that $\mathcal{D}$ is the restriction norm $\frac{1}{2} \|(\omega_0 - \omega_\pi)|_{\mathcal{A}_{\text{acc}}}\|$, defined for any von Neumann subalgebra whether or not it is a tensor factor.

3 The record algebra

Theorem 1 (Commutativity and nondemolition for arbitrary dynamics). *In the collision model (1), for arbitrary H_S and arbitrary coupling operator in place of σ_- : (i) $[Y_k, Y_l] = 0$ for all k, l ; (ii) $[j_N(X), Y_k] = 0$ for every system observable X and every $k \leq N$; (iii) $U_{1..N}^\dagger Q_k U_{1..N} = Y_k$ for all $N \geq k$.*

Proof. Slice k enters the dynamics only at step k . Write $U_{1..N} = U_{k+1..N} U_{1..k}$ with $U_{k+1..N} = U_N \cdots U_{k+1}$. Each factor of $U_{k+1..N}$ acts on $\mathcal{H}_S \otimes \mathcal{H}_{a_{k+1}} \otimes \cdots$ and as the identity on $\mathcal{H}_{a_k}$, so $[U_{k+1..N}, Q_k] = 0$; this gives (iii): $U_{1..N}^\dagger Q_k U_{1..N} = U_{1..k}^\dagger U_{k+1..N}^\dagger Q_k U_{k+1..N} U_{1..k} = U_{1..k}^\dagger Q_k U_{1..k}$. For (i) with $l > k$, $Y_l = U_{1..k}^\dagger (U_{k+1..l}^\dagger Q_l U_{k+1..l}) U_{1..k}$ and the bracketed operator acts on $\mathcal{H}_S \otimes \mathcal{H}_{a_{k+1}} \otimes \cdots \otimes \mathcal{H}_{a_l}$, hence commutes with Q_k ; conjugation by $U_{1..k}$ preserves commutation. (ii) is the same argument with X in place of Q_l : $j_N(X) = U_{1..k}^\dagger (U_{k+1..N}^\dagger X U_{k+1..N}) U_{1..k}$ and the bracketed operator does not act on $\mathcal{H}_{a_k}$. No property of H_S or of the coupling was used. $\square$

Numerically, with (2): $\max_{k,l} \|[Y_k, Y_l]\| = 1.2 \times 10^{-15}$, $\max \|[j_N(X), Y_k]\| = 1.2 \times 10^{-15}$ over $X \in \{\sigma_x, \sigma_y, \sigma_z\}$, and the stability deviation in (iii) is 1.2×10^{-15} .^[R1, R2] In the continuous case the same three statements are the self-nondemolition and nondemolition theorems of [11, 12].

Theorem 2 (Histories; exclusivity as a theorem). $\mathcal{R}_N \cong L^\infty(\Omega_N, \mathbb{P})$ with $\Omega_N = \{0, 1\}^N$ the set of records. The joint spectral projectors

$$\Pi_r = U_{1..N}^\dagger (P_{r_1} \otimes \cdots \otimes P_{r_N}) U_{1..N}, \quad r \in \Omega_N, \quad (5)$$

are mutually orthogonal and sum to the identity, and for a product initial state $\rho_0 = \rho_S \otimes |g \cdots g\rangle\langle g \cdots g|$,

$$\mathbb{P}(r) = \text{Tr}[\rho_0 \Pi_r] = \text{Tr}[\mathcal{I}_{r_N} \circ \cdots \circ \mathcal{I}_{r_1}(\rho_S)], \quad (6)$$

where $\mathcal{I}_r$ is the single-step instrument on the system obtained by coupling to a fresh slice, projecting the slice onto P_r and tracing it out.

Proof. By Theorem 1(i), $\mathcal{R}_N$ is a commutative von Neumann algebra generated by N commuting self-adjoint operators with spectra $\{\pm 1\}$, so its Gelfand spectrum is $\{0, 1\}^N$ and its minimal projectors are the joint spectral projectors. Since $Y_k = U_{1..N}^\dagger Q_k U_{1..N}$ for all k (Theorem 1(iii)) and the Q_k act on distinct tensor factors, the joint projector for the record r is Π_r as displayed; orthogonality and completeness are inherited from those of the P_{r_k} under the unitary conjugation. For (6), insert $U_{1..N} = U_N \cdots U_1$ and use that P_{r_k} commutes with U_l for $l > k$ to move each P_{r_k} next to U_k : $\text{Tr}[\rho_0 \Pi_r] = \text{Tr}[P_{r_N} U_N \cdots P_{r_1} U_1 \rho_0 U_1^\dagger P_{r_1} \cdots U_N^\dagger P_{r_N}]$. Tracing out slice k immediately after P_{r_k} is applied (it is no longer acted upon) gives the composition of single-step instruments. $\square$

Numerically the $2^5 = 32$ history probabilities computed from the spectral projectors and from the sequential instrument agree to 10^{-16} and sum to one; the largest product of two distinct projectors has norm 1.7×10^{-16} .^[R4] Here $\mathbb{P}$ is the spectral measure of ρ_0 on the record algebra, standard quantum probability; that it is the law of chance of the actual record is what Appendix A derives.

Remark 1 (What this does to the postulate ledger). The companion paper posits, as a commitment separate from the actuality of one history, that “the outcome register is exclusive — each anchoring event contributes exactly one outcome from the record-supported alternatives.” Once actuality is attached to a point of Ω_N , that is not a further commitment: the points of the spectrum of a commutative algebra are exclusive by construction. Section 8 draws the consequence.

Theorem 3 (Induced system effects; where classicality was sought in the wrong algebra). *Let the slice coupling at every step be controlled by a system observable O_{rec} (the slice is rotated conditionally on the eigenvalue of O_{rec}), and let $E_k(r) = \text{Tr}_E[(I \otimes \rho_E) U_{1..k}^\dagger (I \otimes P_r^{(k)}) U_{1..k}]$ be the effect induced on the system by the record r at step k . If $[H_S, O_{\text{rec}}] = 0$ then all $E_k(r)$ are functions of O_{rec} and commute. The converse is false: commutation of the effects is a property of the pair (H_S, O_{rec}) , not of H_S alone.*

Proof. If $[H_S, O_{\text{rec}}] = 0$, every U_k is block-diagonal in the eigenbasis of O_{rec} , so the pullback of any slice projector to the system is diagonal in that basis; diagonal operators commute. For the converse take $H_S = \frac{1}{2}\Omega\sigma_x$. With $O_{\text{rec}} = \frac{1}{2}(I + \sigma_x)$ the hypothesis holds and the effects at steps 1, 2, 3 commute to 4×10^{-17} ; with the same H_S and $O_{\text{rec}} = \frac{1}{2}(I + \sigma_z)$, $\max \| [E_k(r), E_l(r')] \| = 1.3 \times 10^{-2}$. [R3, R8] $\square$

The content of Theorem 3 is negative and clarifying: the algebra whose classicality carries a single history is the record algebra $\mathcal{R}_N$, which is commutative for every (H_S, O_{rec}) , and not the algebra of system effects, which is commutative only under the QND condition. Attempts to obtain a Gelfand spectrum of histories from system effects fail outside QND for this reason, and the failure says nothing against the record-algebra ontology.

Theorem 4 (The accessible restriction norm). *Let ω_0, ω_π be the relative-phase pair of (4) on $\mathcal{H}_A \otimes \mathcal{H}_B$ and $\mathcal{A}_{\text{acc}} = B(\mathcal{H}_A) \otimes I$. Then*

$$\sup\{\frac{1}{2}|\omega_0(A) - \omega_\pi(A)| : A \in \mathcal{A}_{\text{acc}}, \|A\| \leq 1\} = \frac{1}{2} \|\text{Tr}_B \omega_0 - \text{Tr}_B \omega_\pi\|_1. \quad (7)$$

Hence $\mathcal{D}$ of (4) is the restriction norm $\frac{1}{2} \|(\omega_0 - \omega_\pi)|_{\mathcal{A}_{\text{acc}}}\|$, which is defined for any von Neumann subalgebra $\mathcal{A}_{\text{acc}}$, tensor factor or not, and is invariant under $*$ -isomorphisms of $\mathcal{A}_{\text{acc}}$.

Proof. For $A = A' \otimes I$, $\text{Tr}[(\omega_0 - \omega_\pi)A] = \text{Tr}[(\text{Tr}_B \omega_0 - \text{Tr}_B \omega_\pi)A']$, and the supremum of $|\text{Tr}[\delta A']|$ over $\|A'\| \leq 1$ is the trace norm of δ , attained at $A' = \text{sgn } \delta$. Invariance under $*$ -isomorphisms is immediate from the definition as a norm on the dual of $\mathcal{A}_{\text{acc}}$. (200 random pairs: maximiser attains the trace norm to 10^{-16} , random contractions never exceed it. [R7]) $\square$

4 Release

4.1 Definition

Definition 3 (Release; permanence). Let R be a record observable formed at time s .

- R is **released** (system-release, condition (S)) if $[R, F] = 0$ for every $F \in \mathfrak{F}_s^S$: nothing in the system’s accessible future — no later system observable and no later fresh record, through any return path the network provides — can act on R .
- R is **permanent** (record-permanence, condition (E)) if $[R, F] = 0$ for every $F \in \mathfrak{F}_s$, i.e. additionally for later readouts of R ’s own carrier in any basis.

Release is *anchoring*; permanence is *readability* or objectivity in Zurek’s sense [26, 27]. Released records commute among themselves (Theorem 1); the actual history is one point of their joint spectrum.

The two conditions coincide when the output is free and both fail for an eraser; they separate under intra-environment scrambling (Proposition 4). A first version of this work used a single condition and conflated them.

Remark 2 (The nucleation analogy). Definition 3 has a classical counterpart that we state once, for orientation. In a supersaturated medium a nucleus below the critical size can dissolve back; past it, the nucleus grows irreversibly. An unreleased record is subcritical in exactly this sense — a return path can still recover its phase — release is the critical size, and permanence (Section 7) is the subsequent growth to a record that many fragments can read. The analogy is exact at the level of the survival law $S = e^{-\Phi_{\text{irr}}}$ of the companion paper, where the hazard plays the nucleation rate, and it is used nowhere in the proofs.

4.2 Release in the free-output setting is nondemolition, and selects nothing

Theorem 5 (Free output). *In the collision model (1) with every slice meeting the system once: (i) for every unit vector $n \in \mathbb{R}^3$ and every k , the readout $Y_k^{(n)} = U_{1..k}^\dagger (n \cdot \vec{\sigma}^{(k)}) U_{1..k}$ is released, and commutes with every readout $Y_l^{(m)}$, $l \neq k$, along any axis m ; (ii) for non-parallel n, m , $[Y_k^{(n)}, Y_k^{(m)}] \neq 0$.*

Proof. (i) is Theorem 1 with $Q_k = n \cdot \vec{\sigma}^{(k)}$, which is any observable of slice k ; the proof used only that slice k is not acted on after step k . (ii) $[Y_k^{(n)}, Y_k^{(m)}] = U_{1..k}^\dagger [n \cdot \vec{\sigma}, m \cdot \vec{\sigma}]^{(k)} U_{1..k} = 2i U_{1..k}^\dagger ((n \times m) \cdot \vec{\sigma}^{(k)}) U_{1..k}$, of norm $2|n \times m|$. $\square$

Numerically, over eight axes (the three Pauli axes and five random ones) and five slices, the largest commutator with a future system observable and the largest commutator between readouts of different slices are both 1.2×10^{-15} , while $\|[Y_k^z, Y_k^x]\| = 2.000$.^[S1a, S1b]

Remark 3 (The delta from Belavkin, stated exactly). Theorem 5 says two things. First, in the setting in which Belavkin’s nondemolition principle is formulated — a free output — release *is* nondemolition: every output observable satisfies it. The present paper adds nothing there and claims nothing there. Second, release does not select the record. The counting-type and quadrature-type readouts of the same slice are both released and generate non-commuting algebras; their union has no joint spectrum and hence no joint history space. Which of them carries the actual history is fixed by $\mathcal{A}_{\text{acc}}$, i.e. by the detector physically present — which is what Wiseman and Gambetta established for unravellings [17]. The claims of this paper concern what release adds *outside* the free-output setting (Sections 4.4–4.5) and the graded form that a physical, partial return path requires (Section 4.3).

4.3 The graded form: a bound, its converse, and the carrier

A physical return path is rarely all or nothing. The companion paper’s functional Φ_{irr} is graded; release as defined is binary. The following theorem is the bridge.

Theorem 6 (Graded release bound). *Let $R = r_+ P_+ + r_- P_-$ be a binary record, $P_+ + P_- = I$, $\Delta = |r_+ - r_-|$. For an operator A write $A_{\text{off}} = P_+ A P_- + P_- A P_+$, and for a state ω write $\omega_{\text{off}} = P_+ \omega P_- + P_- \omega P_+$. Let $S = P_+ - P_-$ and $\omega_\pi = S \omega S$ (the state with the relative phase between the two record alternatives flipped — the state that (4) compares ω against). Then:*

- (a) $\|[R, A]\| = \Delta \|A_{\text{off}}\|$ in every unitarily invariant norm;
- (b) for every A with $\|A\| \leq 1$, $|\text{Tr}(\omega_{\text{off}} A)| \leq \|\omega_{\text{off}}\|_1 \|[R, A]\| / \Delta$;
- (c) consequently

$$\mathcal{D}_{\text{acc}}(\omega, \omega_\pi) := \sup_{A \in \mathcal{A}_{\text{acc}}, \|A\| \leq 1} \frac{1}{2} |\text{Tr}[(\omega - \omega_\pi) A]| \leq \frac{\|\omega_{\text{off}}\|_1}{\Delta} \sup_{A \in \mathcal{A}_{\text{acc}}, \|A\| \leq 1} \|[R, A]\|, \quad (8)$$

and in particular, if R is released then $\mathcal{D}_{\text{acc}}(\omega, \omega_\pi) = 0$ for every state ω .

Proof. (a) Insert $I = P_+ + P_-$ on both sides of A in $[R, A] = r_+ P_+ A + r_- P_- A - r_+ A P_+ - r_- A P_-$. The diagonal blocks $P_\pm A P_\pm$ cancel and $[R, A] = (r_+ - r_-)(P_+ A P_- - P_- A P_+)$. Since $S A_{\text{off}} = (P_+ - P_-)(P_+ A P_- + P_- A P_+) = P_+ A P_- - P_- A P_+$, we have $[R, A] = (r_+ - r_-) S A_{\text{off}}$ with S unitary, whence $\|[R, A]\| = \Delta \|A_{\text{off}}\|$ in any unitarily invariant norm. (b) ω_{off} has only off-diagonal blocks, so $\text{Tr}(\omega_{\text{off}} A) = \text{Tr}(\omega_{\text{off}} A_{\text{off}})$: the diagonal blocks of A pair with the vanishing diagonal blocks of ω_{off} . Hölder's inequality gives $|\text{Tr}(\omega_{\text{off}} A_{\text{off}})| \leq \|\omega_{\text{off}}\|_1 \|A_{\text{off}}\|$, and (a) converts $\|A_{\text{off}}\|$. (c) S leaves the diagonal blocks of ω and negates the off-diagonal ones, so $\omega - \omega_\pi = 2\omega_{\text{off}}$ and $\frac{1}{2} |\text{Tr}[(\omega - \omega_\pi) A]| = |\text{Tr}(\omega_{\text{off}} A)|$; take the supremum over the accessible unit ball and apply (b). If $[R, A] = 0$ for all $A \in \mathcal{A}_{\text{acc}}$ every term vanishes. $\square$

Two checks: the identity (a) holds for 60 random A to 1.6×10^{-15} , and the bound (b) holds for 60 random state-observable pairs (largest ratio of left to right side 0.05) and, with the accessible maximiser computed exactly by Theorem 4, on all three return families of Table 1 (largest ratio 1.0000).^[S2a]

Remark 4 (Tightness). Over the full unit ball, $\sup_A |\text{Tr}(\omega_{\text{off}} A)| = \|\omega_{\text{off}}\|_1$, attained by the polar part of $\omega_{\text{off}}^\dagger$, which is itself off-diagonal; for that A , $\|[R, A]\| = \Delta \|A_{\text{off}}\| = \Delta$. So (8) is an equality whenever $\mathcal{A}_{\text{acc}}$ contains the polar unitary of ω_{off} , and Table 1 shows equality along the coherent-return family. For a record with $m > 2$ values and smallest gap $\Delta_{\min}$, $\|[R, A]\|_2^2 = \sum_{r \neq r'} (r - r')^2 \|P_r A P_{r'}\|_2^2 \geq \Delta_{\min}^2 \|A_{\text{off}}\|_2^2$ gives $|\text{Tr}(\omega_{\text{off}} A)| \leq \|\omega_{\text{off}}\|_2 \|[R, A]\|_2 / \Delta_{\min}$.

Proposition 1 (The converse is false). *There are states ω with $\omega_{\text{off}} = 0$; for them $\mathcal{D}_{\text{acc}}(\omega, \omega_\pi) = 0$ under every return coupling, while $\|[R, j(X)]\|$ can be maximal.*

Proof. Take $\omega_{\text{dec}} = P_+ \omega P_+ + P_- \omega P_-$ for any ω : then $\omega_{\text{dec}} - S \omega_{\text{dec}} S = 0$ identically. In the model, under a full return of slice 1 to the system, $\|[R, j(X)]\| = 1.683$ while $\mathcal{D}_{\text{acc}} = 0$ and $\|\omega_{\text{off}}\|_1 = 1.1 \times 10^{-16}$.^[S2b] (The system pointer state $|e\rangle$ is *not* such a state: partial exchange makes the click random, $\|\omega_{\text{off}}\|_1 = 0.841$, and the phase is recoverable, $\mathcal{D}_{\text{acc}} = 0.671$.) $\square$

The commutator says a return path *could* recover phase; whether there is phase to recover is a property of the state. Release is the state-independent, sufficient condition; Φ_{irr} is the state-level quantity.

Proposition 2 (The commutator is blind to the carrier's later history). *Let V be any unitary inserted after the record forms that commutes with the carrier observable Q_k (for instance a dephasing or scrambling of the carrier in its record basis). Then for every later system observable X , $\|[Y_k, U^\dagger X U]\|$ is unchanged by the insertion of V , while $\mathcal{D}_{\text{acc}}$ is not.*

Proof. With $U = U_{\text{ret}} V U_{1..k}$,

$$[Y_k, U^\dagger X U] = U_{1..k}^\dagger [Q_k, V^\dagger U_{\text{ret}}^\dagger X U_{\text{ret}} V] U_{1..k} = U_{1..k}^\dagger V^\dagger [Q_k, U_{\text{ret}}^\dagger X U_{\text{ret}}] V U_{1..k}$$

because $[V, Q_k] = 0$; the commutator is unitarily conjugated and its norm is unchanged (max change 4.4×10^{-16}).^[S2d] That $\mathcal{D}_{\text{acc}}$ changes is Table 1, family (ii). $\square$

Table 1 is the empirical content of this subsection. Along a coherent partial return (family (i)) the commutator and the recoverable phase rise together and the bound is tight. Along family (ii) the commutator does not move while the recoverable phase falls to zero: the commutator measures the size of A 's off-diagonal blocks, which the return coupling fixes, but the phase actually recovered is the overlap $\text{Tr}(\omega_{\text{off}} A_{\text{off}})$, and dephasing the carrier rotates ω_{off} within the off-diagonal sector without changing its norm. Across the three families there are 15 order inversions among 113 ordered pairs: the commutator does not rank recoverability.^[S2c]

Table 1: Three partial-return families for the record on slice 1 (formed under $H_S = \frac{1}{2}\Omega\sigma_z$ so that the return is exact), initial state $|+\rangle \otimes |g \cdots g\rangle$, $\mathcal{A}_{\text{acc}} =$ later system observables. $\mathcal{D}_{\text{acc}}$ is computed exactly with the maximiser of Theorem 4; the commutator is maximised over $X \in \{\sigma_x, \sigma_y, \sigma_z\}$. Family (i): $e^{+i\varepsilon H_{\text{int},1}}$ after step 1. Family (ii): slice 1 dephased in its record basis by coupling to slice 2 with strength p , then returned fully. Family (iii): return through $\sigma_x \otimes \sigma_x^{(1)}$ with strength ε . [S2a, S2c]

family	parameter	$\mathcal{D}_{\text{acc}}$	$\max_X \ [R, j(X)]\ $
(i) coherent partial return	ε 0, .1, .25, .5, .75, 1	= 0, .047, .115, .225, .324, .407	0, .200, .495, .959, 1.363, 1.683
(ii) dephased carrier, full return	p 0, .25, .5, .75, 1	= .407, .376, .288, .156, 0	1.683 at every p
(iii) return via $\sigma_x \otimes \sigma_x$	ε 0, .1, .25, .5, .75, 1	= 0, .079, .191, .336, .398, .363	0, .397, .959, 1.683, 1.995, 1.819

Corollary 1 (What “graded release” is). *Release (Definition 3) is the exact, state-independent, coupling-only condition. Its graded form is Φ_{irr} of (4): the recoverable phase that the accessible future actually leaves. Theorem 6 bounds the second by the first; the bound is tight for coherent return and loose whenever the carrier’s own history has already destroyed the alignment. The companion paper’s choice of Φ_{irr} as the operative quantity is thereby confirmed, and the release/permanence distinction of Definition 3 is sharpened: the commutator with $\mathfrak{F}_s^S$ is purely on the release side; the carrier’s history is on the permanence side; the phase actually recoverable depends on both.*

Corollary 2 (Actuality is empirically idle exactly at release). *Let $R = r_+P_+ + r_-P_-$ be a binary record formed at time s , let ω be any state, and let $\omega_{\text{dec}} = P_+\omega P_+ + P_-\omega P_-$ be the mixture obtained by attaching a definite value to R with the Born weights and conditioning on it without reading it. Then for every $A \in \mathfrak{F}_s^S$ with $\|A\| \leq 1$,*

$$|\text{Tr}(\omega A) - \text{Tr}(\omega_{\text{dec}} A)| \leq \frac{\|\omega_{\text{off}}\|_1}{\Delta} \|[R, A]\|. \quad (9)$$

Consequently: (i) if R is released, attaching actuality to it changes no accessible prediction, for every state; (ii) if R is not released, there exist an accessible observable $A \in \mathfrak{F}_s^S$ and a state ω for which the two predictions differ.

Proof. $\omega - \omega_{\text{dec}} = \omega_{\text{off}}$, so the left side of (9) is $|\text{Tr}(\omega_{\text{off}} A)|$ and Theorem 6(b) gives the bound; (i) follows because every commutator on the right vanishes. For (ii), not released means $[R, A] \neq 0$ for some $A \in \mathfrak{F}_s^S$, which we may take self-adjoint (if $[R, A] \neq 0$ then R fails to commute with the Hermitian or the anti-Hermitian part of A), and which by Theorem 6(a) means $P_+AP_- \neq 0$; choose unit vectors $u \in \text{ran } P_+$, $v \in \text{ran } P_-$ with $a = \langle u|A|v \rangle \neq 0$ and set $\psi = (u + e^{i\varphi}v)/\sqrt{2}$, $\omega = |\psi\rangle\langle\psi|$. Then $\text{Tr}(\omega_{\text{off}} A) = \text{Re}(e^{i\varphi}a)$, which equals $|a| > 0$ for $\varphi = -\arg a$. $\square$

Corollary 2 is the statement that fixes where a single-history ontology over unmodified quantum mechanics may place its events. A rule that makes a record definite before release, in the sense of the corollary — the actual history proceeds from the record, the unselected alternatives no longer enter that history, and the global wave is untouched — changes what the actual history predicts: by (ii) that history predicts, for some preparation and some accessible observable, a value that unitary evolution does not, and the observable in question is precisely one through which the recorded phase can still return. A rule that assigns values at all times and lets the unselected alternatives keep entering the actual history, as the whole wave keeps guiding a Bohmian configuration, is not of this kind, and nothing here bears on it. A rule that places the event at release, or later, is empirically idle by (i). Section 6 shows that the laboratory has tested the earlier placements and found the unitary prediction every time; the

present criterion is therefore the earliest placement the data leave standing, and nothing in the data chooses among the later ones.

4.4 Deferred anchoring under time-delayed return

The free-output assumption fails whenever an emitted carrier re-interacts with the system. The simplest such network is an atom in front of a mirror — the half-cavity of Dorner and Zoller [42] and the time-delayed coherent feedback of Grimsmo [43] and Pichler and Zoller [44] — which is non-Markovian for the system with memory equal to the round-trip time. In the collision model: slice k interacts at step k , is reflected, interacts again at step $k + \tau$ with an additional phase ϕ in the coupling, and then leaves for good. We use $\tau = 2$, $\phi = 0.9$, H_S as in (2). Write $Y_k^{(1)} = U^{(k)\dagger} Q_k U^{(k)}$ for the record pulled back through the dynamics up to step k (the in-flight record) and $Y_k^{\text{exit}} = U^{(k+\tau)\dagger} Q_k U^{(k+\tau)}$ for the record pulled back through step $k + \tau$ (the exit record).

Proposition 3 (Deferred anchoring). *In the delayed-return network:*

- (a) *the in-flight record is nondemolition at intermediate times, $[Y_k^{(1)}, j_t(X)] = 0$ for $k < t < k + \tau$, but is not released: $[Y_k^{(1)}, j_t(X)] \neq 0$ for $t \geq k + \tau$;*
- (b) *the exit record is released, is a stable Heisenberg observable after exit, and the exit-record algebra $\text{vN}\{Y_1^{\text{exit}}, \dots, Y_N^{\text{exit}}\}$ is commutative; its joint spectral projectors are orthogonal and complete, and their probabilities equal the path measure of sequential readout of each slice at its exit, i.e. of the instrument on the enlarged system (system plus in-flight slices);*
- (c) *the in-flight outcome and the exit outcome differ with finite probability.*

Proof. (a) For $k < t < k + \tau$ slice k is not acted on between steps k and t , so the argument of Theorem 1(ii) applies. For $t \geq k + \tau$ the factor $U^{(t)}$ contains the second interaction of slice k with the system, which does not commute with Q_k , and the commutator is generically nonzero: in the model, $\max_X \|[Y_2^{(1)}, j_3(X)]\| = 5.2 \times 10^{-16}$ while $\max_X \|[Y_2^{(1)}, j_T(X)]\| = 1.614$ for $T = N + \tau$.^[S3a] (b) After step $k + \tau$ slice k is never acted on again, so Theorem 1 applies verbatim to the exit records with $U^{(k+\tau)}$ in place of $U_{1..k}$: $\max_X \|[Y_2^{\text{exit}}, j_T(X)]\| = 1.2 \times 10^{-15}$, stability 1.2×10^{-15} , $\max_{k,l} \|[Y_k^{\text{exit}}, Y_l^{\text{exit}}]\| = 1.5 \times 10^{-15}$, orthogonality 1.7×10^{-16} , normalisation $1 - 10^{-12}$; the 32 history probabilities agree with sequential exit readout to 5.6×10^{-17} .^[S3b] The path-measure statement is Theorem 2 applied to the enlarged system. (c) With the initial state $\sqrt{0.7}|e\rangle + \sqrt{0.3}|g\rangle$, $\mathbb{P}(\text{first-pass click on slice 2}) = 0.123$ and $\mathbb{P}(\text{no photon at exit} \mid \text{first-pass click}) = 0.195$.^[S3c] Without the return the first-pass record is already released (6.1×10^{-16}).^[S3d] $\square$

Proposition 3 is the sharpest positive content of the release condition. A theory that anchored on the in-flight record would, one time in five here, assert a fact that the network later overturns. Release refuses to anchor until the carrier’s future interaction cone no longer meets the system’s; anchoring is *deferred to exit*. This is a statement about the network, not about an observer. Two further consequences: the actual history is a point of the spectrum of the exit-record algebra, well defined with no Markov approximation anywhere; and in the class of carriers that never exit — the atom–photon bound states of the half-cavity [42], perfectly reflecting resonators — no record is ever released and no event occurs, which is the physically correct verdict: the excitation oscillates and records nothing. The continuum limit of this construction (a stochastic differential equation with delay is not of Hudson–Parthasarathy form, and the matrix-product-state treatments [44] define the output algebra at the final time only) is stated as open in Section 10.

4.5 Absolute centre versus accessible centre: scrambling

The second setting in which the output is not free is an environment that continues to evolve internally after the record forms. Riedel, Zurek and Zwolak [25] showed that in quantum Brownian motion the redundancy of the position record rises on the decoherence timescale and then *falls* on the environment-mixing timescale, as intra-environment interactions delocalise the record, while the system’s reduced dynamics is unaffected. In the collision model: after steps 1 and 2 a random unitary acts on slices 1, 2, 4, 5 (all present slices except the one about to interact), then step 3 proceeds.

Proposition 4 (Scrambling). *After full-environment scrambling the record Y_1 (i) is released — it commutes with all later system observables and fresh records; (ii) is not permanent — it fails to commute with readouts of the scrambled carriers; and (iii) the system’s reduced state, and hence its entire future statistics and the step-1 instrument on it, are identical with and without the scrambling. An eraser (return path) fails both conditions. A coherent operation on “the friend’s laboratory” (system plus record carrier) that is part of the network makes the record fail release.*

Proof. (i) is Theorem 1: the scrambler does not touch the system or slice 3, so it commutes with Q_3 and with X , and slice 1 is never re-coupled to the system; numerically 1.3×10^{-15} . (ii) The scrambler entangles slice 1 with slices 2, 4, 5, so Q_1 pulled back through it does not commute with $\sigma_{x,y,z}$ on those slices: $\max \| [Y_1, \cdot] \| = 1.995$. (iii) The scrambler acts on the environment only and is traced out: $\| \rho_S^{\text{scr}} - \rho_S \| = 4.1 \times 10^{-16}$. Eraser: $\| [Y_1, j(X)] \| = 1.683$; laboratory operation: 1.999.^[S4, R9] $\square$

Here an absolute-centre criterion and the accessible-centre criterion give different verdicts, and this is the case that decides between them. A criterion that requires the record to commute with *everything* in the future algebra — the natural reading of “beable” when the output is not free — says that after scrambling there is no beable and hence no event. But nothing that touches the system has changed: its reduced state, its future statistics and the instrument that acted on it at step 1 are the same to machine precision. Such a criterion makes the occurrence of an event depend on the environment’s internal bookkeeping after the carrier has left, which is non-locality of the kind that matters and is contrary to the reading that Riedel, Zurek and Zwolak themselves give of their result. The accessible-centre criterion says the event occurred and its readability degraded (the conditional-state trace distance of the carrier falls from 0.230 to 0.172^[R5c]), which is what quantum Darwinism says happens. On the other two rows of the table the frameworks agree: wherever a return path to the system exists, whether an eraser or a coherent operation on the whole laboratory, the record is not released and no event occurred for it. The extended Wigner’s-friend scenarios [21] are treated the same way: whether the friend’s record is released is a fact about whether the reversing operation is physically in the friend’s future cone, not a fact relative to an observer.

4.6 The future-reference objection

Release refers to the *future* accessible algebra. Read naively, whether a record is released now depends on what the apparatus will do later. Three things settle this.

First, release is a property of the interaction network, not of a realised future. The commutant condition is evaluated in the Heisenberg picture with the actual total dynamics, apparatus included; in the collision model “the slice never returns” is a statement about the coupling graph, knowable at s if the graph is known. A record is released if and only if it lies outside the system’s future interaction cone in that graph. An experimenter who “later adds a return path” was always part of a network with that path.

Second, the residual reference to the future is inherited, and shared by every record-based one-history account: Steane’s non-erasure condition [21], Gell-Mann and Hartle’s records that

are never decohered away [22], Riedel’s decompositions defined by records that stay redundant [24]. Release adds no future-reference; it states the one that every irreversibility notion carries as a commutator rather than as prose.

Third, this is the one place where the atemporal reading of the companion paper’s founding picture does work. What is actual is a point of the spectrum of the released-record algebra of the whole network, and “the whole network” is not a time-slice. Belavkin’s “quantum future / classical past” superselection is the same structure. We say this once and lean on it nowhere else; every derivation in this paper is in laboratory time.

4.7 The accessible algebra is not a free parameter

Three points answer the objection that $\mathcal{A}_{\text{acc}}$ replaces the free parameters of collapse models. (i) Standard quantum-trajectory theory needs the same input: “different detection schemes result in different types of jumps (or diffusion)” [17], and no theory of continuous measurement writes a conditioned state without specifying the detection scheme. $\mathcal{A}_{\text{acc}}$ is that specification; nothing is added relative to the theory inherited. (ii) What this paper adds is a *property*, not a choice: given the setup, release is a statement about commutators in the actual dynamics, and Theorem 4 makes Φ_{irr} a coordinate-free restriction norm. (iii) What remains open is the autonomous case — a dust grain in the microwave background, a molecule in a gas with no designated apparatus — where “accessible” must be specified by the interaction network itself: which fragments can, by the actual dynamics, ever act back on the system or be read jointly. Riedel meets the same question in defining decompositions by spatially local redundant records [24]; Theorem 5 makes it sharper here, since release does not select and the free-output setting therefore has no selection principle at all until $\mathcal{A}_{\text{acc}}$ is given. Riedel’s demonstration that Markovian quantum Brownian motion is an iterated entanglement-breaking measurement with coherent-state pointers on a preferred timescale [23] is the candidate answer for the matter-wave case and is a conjecture.

5 The event law

Theorem 7 (The WHEN law, general form). *For a released record the law of events is the instrument path measure (6), equivalently the law of the output process under the quantum filter. In the continuous case: for a released counting record with jump operator L_r the event intensity is*

$$\lambda_r(t) = \text{Tr} (L_r^\dagger L_r \rho_c(t)), \quad (10)$$

and for a released quadrature record $dY_t = \langle X_\theta \rangle_t dt + dW_t$ with $X_\theta = e^{-i\theta} L + e^{i\theta} L^\dagger$; the conditioned state ρ_c is the Belavkin filter (the conditional expectation onto $\mathcal{R}_t$), and its ensemble average reproduces the unconditioned Lindblad dynamics for every initial state.

Proof. This is the content of the filtering theorems [5, 11, 12] together with Theorem 2: the probability of a history is $\text{Tr}[\rho_0 \Pi_r]$, which is the path measure; the counting intensity is the infinitesimal version of the single-step instrument’s outcome probability, $\text{Tr}[\mathcal{I}_r(\rho_c)] = \text{Tr}(L_r^\dagger L_r \rho_c) dt + O(dt^2)$; exactness in ensemble is completeness of the instrument, $\sum_r \mathcal{I}_r = e^{\mathcal{L} dt}$. $\square$

Corollary 3 (The QND hazard of the companion paper). *If the record channels are $L_\pm = \sqrt{\Lambda} P_\pm$ with $\sum_\pm L_\pm^\dagger L_\pm = \Lambda I$, then the total intensity is Λ , independent of the conditioned state, the no-event evolution is outcome-neutral, and $\Lambda = \Phi_{\text{irr}}$: the event law $h = g \Phi_{\text{irr}}$ of [1], with its Propositions 1 and 2 (uniqueness of the hazard; Born weights forced by outcome neutrality), is the QND case of Theorem 7.*

Proof. $\sum_\pm \text{Tr}(L_\pm^\dagger L_\pm \rho_c) = \Lambda \text{Tr} \rho_c = \Lambda$. The no-event generator $-\frac{i}{\hbar} H_{\text{eff}} = -\frac{i}{\hbar} H_S - \frac{1}{2} \Lambda I$ is proportional to the identity in its non-Hermitian part, so the no-event evolution does not reweight

the alternatives. For a QND record the accessible distinguishability of the phase pair decays as $\mathcal{D}(t) = \mathcal{D}(0)e^{-\int_0^t \Lambda}$, so $\dot{\Phi}_{\text{irr}} = \Lambda$. $\square$

Outside the QND window the hazard $h = \dot{\Phi}_{\text{irr}}$ is not merely unestablished but wrong: for a driven two-level atom with a photodetected output, the extrapolated law gives the wrong lifetime, the wrong click probability and no antibunching, while (10) with $L = \sqrt{\gamma}\sigma_-$ gives $\rho_{ee}(t) = \rho_{ee}(0)e^{-\gamma t}$, $\mathbb{P}(\text{click}) = \rho_{ee}(0)$ and the waiting-time density $w(\tau) = \gamma|\langle e|e^{-iH_{\text{eff}}\tau}|g\rangle|^2$ with $w(0^+) = 0$ — Carmichael’s delay function [29, 30]. The companion paper scoped its law to QND record formation and so is untouched; Theorem 7 is what stands in the general case. The scope condition of [1] (CP-divisible record-forming regimes) receives its explanation here: it is the class in which a released record exists at all (Theorem 1 needs fresh, non-returning carriers), and Hall *et al.* [28] show that within it a time-local jump unravelling with non-negative rates exists exactly when the reduced dynamics is CP-divisible.

Theorem 8 (Diffusive limit). *Balanced homodyne detection with local-oscillator amplitude α realises the counting channels $L_{\pm} = (\alpha \pm e^{-i\theta}L)/\sqrt{2}$, with intensities $\lambda_{\pm} = \frac{1}{2}[\alpha^2 + \langle L^{\dagger}L \rangle \pm \alpha\langle X_{\theta} \rangle]$. The scaled record $dY = (dN_+ - dN_-)/\alpha$ satisfies $\mathbb{E}[dY] = \langle X_{\theta} \rangle dt + O(1/\alpha)$ and $\text{Var}[dY] = dt + O(1/\alpha^2)$, and as $\alpha \rightarrow \infty$ the conditioned state obeys $d\rho_c = \mathcal{L}\rho_c dt + \mathcal{H}[e^{-i\theta}L]\rho_c dW$ [14, 15, 16]. The ensemble is exact at every α .*

Proof. $L_{\pm}^{\dagger}L_{\pm} = \frac{1}{2}(\alpha^2 \pm \alpha X_{\theta} + L^{\dagger}L)$ gives the intensities and, since $\sum_{\pm} L_{\pm}^{\dagger}L_{\pm} = \alpha^2 + L^{\dagger}L$ and $\sum_{\pm} L_{\pm}\rho L_{\pm}^{\dagger} = \alpha^2\rho + L\rho L^{\dagger}$, the unconditioned generator is the original Lindbladian for every α (the α^2 terms cancel between jump and no-jump parts). The moments follow from $\mathbb{E}[dN_{\pm}] = \lambda_{\pm}dt$ and $\text{Var}[dN_{\pm}] = \lambda_{\pm}dt$ to leading order; the limit is the construction of [14] (“ $c \rightarrow c + \gamma$, $|\gamma|^2 \gg \langle c^{\dagger}c \rangle$ ”). Numerically, with a driven atom ($\Omega = 1.3\gamma$) and $\alpha = 3, 8, 20$, the ensemble matches the exact Lindblad solution to $\leq 2.5 \times 10^{-2}$ at every α (the counting step must satisfy $\delta t \ll 1/\alpha^2$), the click rate is α^2 per unit time, $\text{Var}(dY)/dt = 1.000$, the windowed mean of dY regresses on $\langle X_{\theta} \rangle dt$ with slope 1.002, and the spread of the conditioned $\langle \sigma_z \rangle$ at $T = 2$ converges to that of the diffusive stochastic master equation.^[R6] $\square$

Remark 5 (Ontological reading). The actual history in a homodyne experiment [57, 58, 59] is the counting record $N_{\pm}$; every local-oscillator photon click is a released record and an anchoring; the Wiener record Y is a coarse-graining of that history, not a second kind of event. Anchoring is progressive with discrete hits at rate α^2 , each moving the conditioned state by $O(1/\alpha)$. No threshold at which “a diffusive signal becomes an event” is needed or allowed.

Identical particles. Electrons carry no labels, and the records of the event law carry none either. Every coupling that writes a record of electrons is built from the electron field ψ . Photons and phonons couple through $\psi^{\dagger}o\psi$: the charge, current and spin densities. A detector that captures an electron couples through $d^{\dagger}\psi$, the tunnelling of an electron into a reservoir of identical electrons; for a wide-band, empty reservoir in the Markov limit this gives the record channel $L_x = \sqrt{\kappa}c_x$ [31], where c_x removes one electron from the mode of pixel x . The coupling and the record, the reservoir’s occupation, are even in the fermion operators, and even operators on distinct modes commute, so the slice argument of Theorems 1 and 2 holds for fermionic reservoir modes as it does for qubit slices. Effective models sometimes add two-body couplings, such as an exchange modulated by phonons; these too treat all electrons alike.

Proposition 5 (Identical particles). *Write N electrons in labelled coordinates on $\mathcal{H}_1^{\otimes N}$, with T_{π} the permutation of the labels by π and P_A the antisymmetriser, and identify $P_A\mathcal{H}_1^{\otimes N}$ with the N -electron sector of Fock space. Let every record channel be a linear combination of products of: one-body operators $\sum_i o^{(i)}$ (in Fock space $\psi^{\dagger}o\psi$); two-body operators $\sum_{i<j} w^{(ij)}$ with w symmetric under exchange of its two particles; Borel functions of self-adjoint operators so built;*

the field operators c_x and $c_x^\dagger$; and frequency components of any of these under a permutation-symmetric system Hamiltonian with pure point spectrum. Then (i) every record channel built without field operators commutes with every T_π and maps the antisymmetric sector to itself; (ii) every record channel maps antisymmetric states to antisymmetric states, each factor c_x removing one electron and each $c_x^\dagger$ adding one; and (iii) for a one-particle operator o , with $o^{(i)}$ acting on the i -th particle,

$$P_A o^{(1)} P_A = \frac{1}{N} P_A \left(\sum_{i=1}^N o^{(i)} \right) P_A. \quad (11)$$

Proof. One-body and symmetric two-body sums commute with every T_π , and so do linear combinations, products and Borel functions of operators that commute with every T_π . Every spectral projector P_E of a permutation-symmetric Hamiltonian commutes with every T_π , so each frequency component $O(\omega) = \sum_{E'-E=\omega} P_E O P_{E'}$ of such an O does too; an operator that commutes with every T_π commutes with P_A , a linear combination of the T_π , and so maps the antisymmetric sector to itself. This is (i). The field operator c_x acts on Fock space, which carries no labels; in labelled coordinates it sends an antisymmetric $\Psi(x_1, \dots, x_N)$ to $\sqrt{N} \Psi(x, x_2, \dots, x_N)$, which is antisymmetric in the remaining arguments, and its adjoint $c_x^\dagger$ sends antisymmetric $(N-1)$ -electron states to antisymmetric N -electron states. The operators of (i) and the spectral projectors preserve antisymmetry in every sector, and linear combinations and products of maps that preserve antisymmetry preserve it; this is (ii). For (iii), a transposition T_{1j} satisfies $T_{1j} P_A = P_A T_{1j} = -P_A$ and $T_{1j} o^{(1)} T_{1j} = o^{(j)}$, so $P_A o^{(j)} P_A = P_A o^{(1)} P_A$ for every j ; summing over j gives (11). $\square$

Equation (11) says that, inside the antisymmetric sector, the operator for “electron 1 is at x ” acts as the density at x divided by N . An operator $o^{(1)}$ on a labelled particle is a record channel only if o is a multiple of the identity: otherwise $T_{12} o^{(1)} T_{12} = o^{(2)} \neq o^{(1)}$, and $o^{(1)}$ moves some antisymmetric states out of the sector, so no record acts on a labelled particle.^[B1–B7] The proposition is the permutation invariance of observables that the identity of particles requires [32]. The event law inherits it because its records are produced by physical couplings, not by an added dynamics. A collapse dynamics written per labelled particle does not inherit it. The model of Ghirardi, Rimini and Weber was formulated for distinguishable particles, each hit acting on one labelled particle [33, 34]. For two same-spin electrons in the ground state of an eight-site chain with nearest-neighbour repulsion, a Gaussian hit of width one site on “electron 1” leaves, at the worst placement, 51% of the post-hit state in the antisymmetric sector, and if each post-hit state is projected back onto that sector the hit probabilities sum to 0.62 instead of one. The leak shrinks as the hit widens past the electrons’ spacing (96% kept at a width of four sites); the failure is one of principle, not of size. The versions for identical particles couple to the number density instead [35, 36, 34], and the symmetric hit keeps the sector.^[E1–E3]

With absorptive pixels at every site, Theorem 7 gives the joint density of an event at x followed at once by an event at y as

$$p(x, y) = \kappa^2 G_2(x, y), \quad G_2(x, y) = \langle c_x^\dagger c_y^\dagger c_y c_x \rangle, \quad (12)$$

the expectation taken in the state just before the first event; G_2 is the fermionic counterpart of the normally ordered correlation function of Glauber’s theory of photodetection [37], read here as the law of the actual history. When c_x is a single mode of one spin, $p(x, x) = 0$ for every state, interacting or not, because $c_x c_x = 0$: two electrons of the same spin are never registered at one pixel at one moment. After an event, the intensity at the same pixel starts at zero, rises as τ^2 , and recovers only as far as the remaining wave flows back. In 40 000 histories sampled event by event for two same-spin electrons in the ground state of a twelve-site chain (hopping J , nearest-neighbour repulsion $3J$, $\kappa = 0.5J/\hbar$), the second event lands on the first event’s pixel 0.79% of the time; the path measure gives 0.81%, and a remaining electron that ignored the first event would give 9.67%.^[C1–C5]

Place two particles one in each of the chain’s two lowest non-interacting orbitals ϕ_1, ϕ_2 : as same-spin fermions (a Slater determinant), as bosons (the symmetrised product), or as labelled particles (the product $\phi_1(x_1)\phi_2(x_2)$, recorded by pixels that do not read labels, for which $G_2^{\text{labelled}}(x, y) = |\phi_1(x)\phi_2(y)|^2 + |\phi_2(x)\phi_1(y)|^2$). The same law gives $G_2^{\text{bosons}} + G_2^{\text{fermions}} = 2G_2^{\text{labelled}}$ at every pair of points; at coincidence the bosons give twice the labelled value and the fermions zero. The sign of the exchange term belongs to the field, not to the law.^[C6] Fermionic antibunching is measured in conductors, as anticorrelated current fluctuations at an electron beam splitter [38, 39], and in atoms; in one apparatus helium-3 antibunches and helium-4 bunches [40]. The two electrons of helium’s ground state have opposite spins and so no exchange hole: in the Hartree–Fock state, in which each occupied spin-orbital holds exactly one electron, the pair density at coincidence equals its uncorrelated value. In the correlated state no mode holds a whole electron, the best holding 0.992 per spin, and with idealised local absorptive detectors inside the atom, a thought experiment, the part of the state that does not fit one electron per spin-orbital is what makes a second event less likely at the place of the first: the Coulomb hole [41]. Gaussian bases cannot form the electron–electron cusp, so the computed coincidence densities are upper estimates and the hole is, if anything, deeper.^[He1, He3] None of this distinguishes the present account from standard quantum mechanics, and the pixels are idealised: one mode each, unit efficiency, no dead time. It shows that, for record channels built from field couplings, the event law needs nothing added for identical particles.

6 Where the laboratory places the event

Corollary 2 makes the placement of an event a question experiments can answer, and they have. Each candidate rule in Table 2 is read as making the recorded variable definite at the moment named, in the corollary’s sense: the actual history proceeds from the record, the unselected alternatives no longer enter that history, and the global wave is untouched. By Corollary 2(ii), under every such rule earlier than release the actual history predicts, in some recoherence experiment, that the interference does not return, whereas unitary quantum mechanics predicts that it does; in every such experiment performed the unitary prediction was found.

Two things follow. The data exclude every single-history rule of this kind that makes a record definite while the record can still coherently return, and release sits on that boundary at the earliest place the data allow; each eraser and returning-carrier experiment was an occasion to refute it, and none did. In Wheeler’s delayed-choice experiments no environmental record is involved and the result follows the final configuration [86, 87]: no record, no event until detection. Later placements, and interpretations with no events at all, survive as well, because each reproduces the statistics of unmodified quantum mechanics (Corollary 2(i)). The experiments fix the boundary and corroborate the placement; whether anything happens there is the one thing they leave to the postulate of Section 8.

The ingredients the criterion uses are themselves measured physics, and a companion record [2] catalogues the experiments — one hundred at the time of writing, from 1922 to 2026 — with what each corroborates and what it does not establish. They fall into four families. A passive dissipative body with no readout destroys interference in proportion to its loss, and a record is what counts, not an interaction: electron waves passing above a doped-silicon plate lose their fringes near it while beside gold, a million times less lossy, they barely change [89, 91]. The rate of record formation follows every coupling that opens a record channel, by one rule: a cavity-field cat decoheres faster for larger separation, with decoherence and damping tied by one coupling [74, 75], and engineered reservoirs dial the rate in a trapped ion [76, 77]. The coupling selects which alternatives can be recorded and never the outcome, and the selection is unitary: a Stern–Gerlach magnet fixes the axis [92], a change of reservoir changes the pointer basis [77], and the split beams recombine [93]. And every detection is one record, with Born weights and a hazard that restarts after a click, and no event leaves two: electrons sent one

Table 2: Candidate placements of a physical event and the experiments that bear on them. Each rule is read as in Corollary 2: the recorded variable becomes definite at the moment named, the actual history proceeds from that record, the unselected alternatives no longer enter it, and the global wave is untouched.

event placed	what it would imply	contradicted by	status
at first interaction	an atom exchanging a photon with a cavity mode has “emitted or not” definite at once; a magnet that splits a beam has decided the spin; a surface that shifts a phase has recorded a path	vacuum Rabi oscillation [73]; the full-loop Stern–Gerlach interferometer [93]; atoms that pass a surface and keep their fringes [88]; electrons beside gold [91]	excluded
at decoherence of the reduced state	which-path is definite, so no conditioning on the environment’s record could restore the fringes	coherence regained by selecting the scattered photon [78]; conditioned fringes in delayed-choice erasers [79, 80]; reversed partial measurements [81, 82]; a quantum jump caught in flight [83]	excluded
when a record carrier exists	emission is a finished fact, so a returning field could not alter the source’s own emission	the mirror-returned field that modulates an ion’s own emission [84]; returning phonons [85]	excluded
at release (Definition 3)	the record is beyond coherent recovery by the system’s accessible future; everything above, as observed	—	survives
later, never, or always	the same statistics, from interpretations without events or with values at all times	—	survive

at a time through a biprism arrive as single marks, which build fringes when no record of the path is kept [90] and a plain band when a lossy slab close to the paths keeps one [89]; one photon at a beam splitter fires one detector and never both while the same photons interfere [94]; and resonance fluorescence is antibunched, from an atomic beam and from a single stored ion [95, 96], which is the waiting-time law of Section 5 and not the extrapolated QND hazard. On the present account the two patterns are two kinds of run. With no slab the mark on the screen is the run’s first event, and the wave reaches the screen with both components. With the slab the path record is released before the screen, the run already has one event, and the mark is its second, following the component the record selected: two events in one history, each leaving one mark. A slab that records weakly anchors some runs and not others, and the contrast that remains is the fraction not yet anchored. A record that is held but can still coherently return, a scattered photon’s for one, decoheres the reduced state as fully and anchors nothing (Table 2, second row). None of these experiments distinguishes the present account from any other that keeps unitary quantum mechanics; what they do is place almost the whole of the mechanism on measured ground and leave one bit, actuality, to the postulate.

7 Readability geometry

Definition 3 separates two questions that the companion paper’s record functional ran together: whether a record is beyond the system’s recall (release), and how well it can be read from the fragments that carry it (permanence, readability). The companion paper’s Definition 1, and the released-record functional built on it, quantify the second.

Definition 4 (Readability of a carrier). For a record R with alternatives $P_{\pm}$ carried by a

fragment F of the environment, the readability of F is the trace distance between the conditional states of F given the two alternatives, $\beta_F = \frac{1}{2} \left\| \rho_F^{(+)} - \rho_F^{(-)} \right\|_1$.

Readability is a permanence-side quantity: it is changed by operations on the carrier that commute with the system (Proposition 2), and it can fall after release without anything returning to the system (Proposition 4: $0.230 \rightarrow 0.172$). Any functional of readability summed over fragments therefore behaves as follows.

Proposition 6 (Ratchet). *A functional $R_{\text{irr}}(t)$ that counts a record at its release and never afterwards — $R_{\text{irr}}(t) = \sum_{k: s_k \leq t} \phi_k$ with ϕ_k evaluated at the release time s_k — is non-decreasing in t by construction, whatever later happens to the readability of the carriers.*

This restates the monotonicity proposition of the companion record-geometry note in the only form in which it is true: the earlier statement, that a readability functional over accessible fragments is monotone, is contradicted by the fall of redundancy [25], and the fix is to build the functional on released records. The uniqueness result of that note — that composition-additivity and merge-invariance force the summary to be linear in the pairwise record kernel — is a statement about the readability coordinate and is unaffected by the present reorganisation; it is not reproduced here. The readability functional built on the released records, its uniqueness under composition-additivity and merge-invariance, and its relation to the hazard are the subject of a separate note and are not used in this paper; the alignment question of Section 10 is where they meet the WHAT side of the event law.

8 What is postulated

The ontic postulate. One point of the joint spectrum of the released-record algebra of the network is actual.

That is the framework’s one primitive commitment. Everything else in this paper is standard open-system physics, a theorem of it, or — for the Born weights — a theorem under the stated premises of Appendix A, and Table 3 says which.

Actuality itself is not derived. The mathematics supplies exclusive alternatives (Theorem 2) with Born weights (Theorem 9); the postulate says that one of them is the case. Two senses of “real” are in play, and we keep them apart. The components of the wave are real as its structure: they carry amplitude, they interfere while no record is released, and what a record leaves unrecorded can anchor later. An alternative is real as fact only when a released record has made it actual. After an event the unselected alternatives keep the first kind of reality and never acquire the second; they are not destroyed, and they are not things. Attached only at release, the postulate changes no prediction (Corollary 2), and it is justified by the explanatory work it does. It is not a hidden variable of Bohm’s kind: Bohm adds a particle with a definite position at all times; the present account adds no particle, only the fact of which marks were made. No new field, constant, threshold, preferred basis or noise is introduced, and no ensemble prediction is altered.

9 Toward a covariant statement

Fewster and Verch [45] couple a system quantum field to a probe field in a compact spacetime region K , obtain induced system observables and Davies–Lewis instruments from the scattering morphism, and prove causal factorisation: for causally orderable coupling regions the composite instrument is the causally ordered composition, for causally disjoint regions the order is immaterial, and non-selective measurements cannot influence causally disjoint experiments.

Table 3: Postulate ledger: the companion paper’s entries and their status here.

item	companion paper [1]	this paper
one record history actual	primitive commitment	primitive commitment (the sole one)
exclusive outcome register	commitment, with the theorem deferred to this paper	theorem (Theorem 2)
event law	$h = g \Phi_{\text{irr}}$, QND-scoped	theorem (Theorem 7); $h = \Phi_{\text{irr}}$ its QND corollary
Born weights	forced in the binary QND model by ensemble agreement with outcome-neutrality (Proposition 2); otherwise proposed	theorem (Theorem 9), conditional on six stated premises; the model’s descent a consistency check
gate $g_{\mathcal{I}}$	structural indicator	Definition 3 (release) + Theorem 5; readability the permanence-side coordinate
scope	CP-divisible record-forming regimes	existence of a released record (fresh, non-returning carriers) — same class, with its reason
accessible algebra	“fixed by physics” (prose)	detector specification + release as a dynamical property (Section 4.7); autonomous case open
irreversibility measure	Definition 1 (Φ_{irr})	graded release, bounded by the commutator (Theorem 6)
diffusive records	not addressed	Theorem 8
covariance	open	partially ordered record history (Section 9); covariant release open

Bostelmann, Fewster and Ruep [46] show that Sorkin-type impossible measurements are excluded when the coupling is local, and Fewster, Jubb and Ruep [47] that every local observable of the free field admits asymptotic measurement schemes.

Two consequences follow for the present framework. The actual history need not be a globally time-ordered sequence on a preferred foliation: it can be a partially ordered set of record events $\{(K_i, r_i)\}$, ordered by the causal order of the coupling regions, with spacelike records unordered and their instruments commuting. That removes the preferred-frame objection to the ontology; the derivations of this paper remain in laboratory time but are no longer required to be. And release is precisely what would terminate the measurement chain that Fewster and Verch explicitly leave open (the probe’s own readout is a further link): a record is actual when it is central relative to the accessible future of the region that produced it.

What is not done: stating release covariantly. Local algebras in algebraic quantum field theory are factors, so “central” must mean central relative to the accessible subalgebra of the apparatus’s causal future, and specifying that subalgebra covariantly — without a frame entering through $\mathcal{A}_{\text{acc}}$ — is the open problem. We do not solve it here.

9.1 Entangled records and Bell nonlocality

The partial order of record events makes the framework’s position on Bell’s theorem precise. Let records a and b be produced in causally disjoint coupling regions K_A and K_B , with settings x and y . Their instruments commute and their composition is independent of order [45], so the law of the actual point, restricted to the two records, is

$$\Pr(a, b \mid x, y) = \text{Tr}[(E_a^x \otimes E_b^y) \rho], \quad (13)$$

with E the induced effects. Because a non-selective instrument in K_A cannot influence K_B , the marginal $\sum_a \Pr(a, b \mid x, y)$ is independent of x : parameter independence holds, and nothing

can be signalled. The order of justification should be stated. Parameter independence is a requirement of the framework, made in [1] for Bell reasons and placed on the law of chance before the weights are known; in Appendix A it is, for every division of the network and not only for causally disjoint regions, the hypothesis (H4), and the law it forces is the instrument law. What (13) shows is that the instrument honours the commitment. The two statements are consistent and not circular: (H4) is the premise, and the no-signalling of the instrument is the check. Outcome independence fails: conditioned on a , the law of b is that of the state left after the record in K_A . For entangled ρ the statistics (13) violate the Bell–CHSH inequality [48], as loophole-free experiments confirm [51, 52, 53, 54]. Since the framework assumes neither setting–source correlations nor retrocausal influence, the one primitive commitment of Section 8 entails Bell nonlocality, of the outcome-dependent kind only [49, 50]. No foliation is needed: the two records are unordered events of the record history, and conditioning in either order gives (13). A single-history theory with nonlocal correlations and no preferred frame has a precedent in the relativistic flash ontology [55]; what remains open is the covariant statement of release discussed above.

10 Open problems

1. *Covariant release* (Section 9), and whether the Fewster–Verch partial order composes with release without a preferred frame entering through $\mathcal{A}_{\text{acc}}$.
2. *The continuum limit of deferred anchoring*. Proposition 3 is proved in the discrete class; the delayed quantum stochastic differential equation is not of Hudson–Parthasarathy form.
3. *Theorem 6 for continuous-valued records*. The bound is stated for finite spectra with a gap; the diffusive limit needs a version in which $\Delta \rightarrow 0$ and $\|\omega_{\text{off}}\|_1 \rightarrow 0$ together (each local-oscillator click shifts the conditioned state by $O(1/\alpha)$).
4. *Selection in the autonomous case* (Section 4.7). Release is necessary, not selective; for an environment with no designated apparatus the accessible algebra must be supplied by the environment’s own fragment structure. Riedel’s coherent-state pointers for quantum Brownian motion [23] are the candidate answer and are unproved as a selection principle. This is the sharpest open item in the programme.
5. *Beyond fresh carriers*. When does a structured (non-Markovian) environment supply an outgoing field carrying a released record? Input–output theory with memory and the process-tensor formulation are the tools; until the answer is known the scope is the class of Theorem 1 and its delayed-return extension.
6. *Release as a stopping time*. Release at time s is defined through the accessible future computed with the actual dynamics of the network (Definition 3), and Section 4.6 says why that dependence carries no outcome. A formulation in which the event time is a stopping time of a filtration generated by the coupling structure, adapted to the records so far, is not written here.
7. *The alignment conjecture*. Whether the record-selected instrument coincides with the dynamically robust (predictability-sieve or maximally robust unravelling) one in the linear–Gaussian regime remains the open question on the WHAT side; [23] suggests that it does, with coherent-state pointers.

11 Conclusion

The output of a continuously monitored quantum system already has the structure a single history needs: its observables commute, its spectrum is the set of histories, and its conditional

expectation is the filter. Belavkin saw this and called it a classical past. What was missing for a realist account was a criterion, stated in the dynamics and not in an observer’s choices, for when a record is beyond the system’s recall, and a demonstration that the criterion does work where nondemolition alone does not. Release — commutation with the system’s accessible future — is that criterion. It reduces to nondemolition when the output is free; it defers anchoring to exit when carriers return; it affirms the event and degrades only readability when the environment scrambles its own records; and it is tied by a proved bound to the recoverable-phase functional with which the companion paper measured irreversibility. Exclusivity is a theorem, the weights are a theorem under the six stated premises of Appendix A, the event law is the instrument path measure with the QND hazard as its corollary, the diffusive stochastic master equation is a coarse-graining of a counting history, identical particles need nothing added because the records carry no particle labels, and the framework carries one primitive commitment. Attaching actuality at release changes no accessible prediction for any state, and attaching it earlier, by a rule under which the unselected alternatives no longer enter the actual history, would; the laboratory has tested each of the three earlier placements of Table 2 and found the unitary prediction, so release is the earliest placement the data leave standing. Which readout of a free output is the actual one is fixed by the detector physically present, not by an observer, and stating what fixes it when there is no detector is the problem that comes next.

The picture with which the companion paper began survives intact and is now sharper. There is one composite field before release and after; release adds one fact; the field does not become a particle, and no duality is left to interpret.

Statements and Declarations

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Data availability. No datasets were generated or analysed. Every numerical statement in this paper is reproduced by five verification scripts (Python 3; numpy and scipy, and PySCF for the helium script), deposited under the MIT License with the companion paper’s code at Zenodo, [doi:10.5281/zenodo.21844368](https://doi.org/10.5281/zenodo.21844368), which resolves to the latest version:

- ACT_RecordAlgebra_Backbone_v0p2.py: 25 checks; seed 20260902;
- ACT_S_StressTest_v0p1.py: 26 checks; seed 20260903;
- ACT_Born_Route_Checks_Oct09_2026.py: six checks, five of them cited here as W-A, W-B, W-C, W-C’ and W-E; seed 20261009;
- ACT_IdenticalFermions_WorkedExample_v0p1.py: 27 checks, seventeen of them cited here as B1–B7, C1–C7 and E1–E3; seed 20261015;
- ACT_IdenticalFermions_Helium_v0p1.py: five checks, two of them cited here as He1 and He3; no random numbers.

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theoretical claims and conclusions are the author's own, and the author has verified and takes full responsibility for the manuscript in its entirety. These tools do not satisfy authorship criteria and are not listed as authors.

A The weights are Born's

Theorem 2 gives the probability of a history as $\text{Tr}[\rho_0 \Pi_r]$, the standard quantum probability of a joint spectral projector, and Theorem 7 reads the event law off it. This appendix shows that those weights are forced, not imported: six stated premises on the law of chance leave only the squared amplitude; three of them restate the ontology and the postulate, one extends the Bell requirement to every division of the network, one is an assumption about the network, and one is new. The argument is Zurek's [98, 99], in the no-signalling form Barnum gave it [100]; nothing in the mathematics is new, and what is added is the premise accounting, which the ontology paper [97] states in condensed form. The argument has been contested where every outcome occurs [101, 102, 103], because there what the probability of an outcome means, when every outcome occurs, is itself the open question; here one record is actual and its chance is a chance, and each contested step is carried by a named premise.

Setting

One global wave Ψ , a unit vector of a separable Hilbert space $\mathcal{H}$, with unmodified unitary dynamics. The tensor factorisation of $\mathcal{H}$ into parts — system, slices, pointers — is given by the network, as in Section 2, and, as in the collision model, the network holds a supply of fresh pointers, each in a fixed initial state and not yet coupled; a *wave of the network* is a unit vector in that sector, $\Psi = \Psi_0 \otimes |c_0\rangle \otimes |d_0\rangle \otimes \dots$, and the theorem is stated for such waves. Record algebras are atomic, as in the discrete realisation of Section 2: a released record algebra $\mathcal{A}$ is a commutative algebra of released records (Definition 3) whose minimal projectors $\{P_r\}_{r \in \Omega(\mathcal{A})}$, countably many, are mutually orthogonal and sum to the identity (Theorem 2), so records are sharp on the total space; the continuous records of Section 2 are reached from the discrete path measure in the limit $\delta t \rightarrow 0$. Records of two parts compose: if $\mathcal{A} = \{P_a \otimes I\}$ and $\mathcal{B} = \{I \otimes Q_b\}$ are released record algebras of parts A and B , their join $\mathcal{A} \vee \mathcal{B} = \{P_a \otimes Q_b\}$ is a released record algebra (Theorem 1: records on distinct factors commute). A *law of chance* assigns to every released record algebra $\mathcal{A}$ and every wave Ψ a function $\mu(\cdot \mid \Psi; \mathcal{A}): \Omega(\mathcal{A}) \rightarrow [0, 1]$, the chance that record r becomes actual. The theorem says what μ must be. Nothing below uses the path measure, the instrument or the filter; those are what the theorem recovers.

Premises

- (H1) *One law, one wave.* The chance depends on the wave and the record projectors only, and the same function serves in every basis: for every unitary U on $\mathcal{H}$ for which $U\mathcal{A}U^\dagger$, the algebra with projectors $UP_rU^\dagger$, is again a released record algebra of the network,

$$\mu(r \mid U\Psi; U\mathcal{A}U^\dagger) = \mu(r \mid \Psi; \mathcal{A}). \quad (14)$$

The law depends on neither the Hamiltonian, the coupling, the factorisation nor a preferred basis. The proof uses only unitaries that map $\mathcal{A}$ onto itself.

- (H2) *One history.* For each released record algebra exactly one record is actual: $\mu(\cdot \mid \Psi; \mathcal{A})$ is a probability measure on $\Omega(\mathcal{A})$, countably additive when $\Omega(\mathcal{A})$ is infinite.
- (H3) *Support.* No record is made where the wave has no component: $P_r\Psi = 0$ implies $\mu(r \mid \Psi; \mathcal{A}) = 0$, for the joint records of a composite algebra as for single ones.

(H4) *No signalling, for every division.* The chance of a record here is unchanged when a record is made elsewhere, and unchanged whether one is made: for every factorisation $\mathcal{H} = \mathcal{H}_A \otimes \mathcal{H}_B$ given by the network and all released record algebras $\mathcal{A}$ of A and $\mathcal{B}$ of B , the marginals of the joint law are the parts' own laws,

$$\mu(a \mid \Psi; \mathcal{A}) = \sum_{b \in \Omega(\mathcal{B})} \mu(a, b \mid \Psi; \mathcal{A} \vee \mathcal{B}), \quad \mu(b \mid \Psi; \mathcal{B}) = \sum_{a \in \Omega(\mathcal{A})} \mu(a, b \mid \Psi; \mathcal{A} \vee \mathcal{B}). \quad (15)$$

The trivial algebra $\{I\}$, no record made, counts among the record algebras of a part.

(H5) *Refinement.* Any released record can be refined by the unmodified dynamics: for every released record algebra $\mathcal{A} = \{P_r\}$ of a part A and integers $m_r \geq 1$, the network supplies a fresh pointer C in a fixed initial state $|c_0\rangle$ and a unitary $W = \sum_r P_r \otimes W_r$ on $A \otimes C$ with $W_r|c_0\rangle = m_r^{-1/2} \sum_{j=1}^{m_r} |c_{rj}\rangle$, the $|c_{rj}\rangle$ orthonormal, such that the projectors $|c_{rj}\rangle\langle c_{rj}|$, completed by their remainder, form a released record algebra $\mathcal{C}$ of C , so that $\mathcal{A} \vee \mathcal{C}$ is a released record algebra of $A \otimes C$ refining $\mathcal{A}$; on $W\Psi$ its supported records are the $P_r \otimes |c_{rj}\rangle\langle c_{rj}|$. With every $m_r = 1$ this is a copy of the record onto a pointer.

(H6) *Continuity.* For fixed $\mathcal{A}$ and r , $\Psi \mapsto \mu(r \mid \Psi; \mathcal{A})$ is continuous in the norm of $\mathcal{H}$.

Where each premise stands in the framework. (H1) is the ontology of [1]: the wave is what exists, and nothing else is there for a chance to depend on. (H2) is the postulate of Section 8 together with ordinary additivity over exclusive actual events; countable additivity is the one strengthening, needed only for infinite algebras. (H3) is that anchoring selects among the wave's components and creates none. (H4) is the framework's commitment to parameter independence (the Bell position of [1], and Section 9.1 here) stated for every division, and the lift from the causally disjoint case is this: any part can be carried away, with its apparatus, by a unitary acting on that part alone; by (H1) that unitary changes no chance of a record here (Lemma 1); once the part is far away the Bell position applies; and by (H1) again the chances in the far configuration are the chances in the near one. The one further premise the lift uses is that the division is into systems, each of which can be carried away with its apparatus. (H4) is the only noncontextuality the argument uses — a record's chance does not change when a record is made elsewhere — and it is a commitment, not something had for free. (H5) is an assumption about the network, realised in the collision model by Theorem 1 with a controlled coupling: a fresh slice rotated conditionally on a record, and never returning, carries a released record, and a pointer with m_r sectors is m_r such slices or one larger one. (H6) is new to the framework and is stated as a premise.

Three lemmas

Lemma 1 (Unitaries that commute with the records). *If $[U, P_r] = 0$ for every $r \in \Omega(\mathcal{A})$, then $\mu(r \mid U\Psi; \mathcal{A}) = \mu(r \mid \Psi; \mathcal{A})$. In particular a unitary acting on a part on which the records of $\mathcal{A}$ act trivially leaves every chance of $\mathcal{A}$ unchanged.*

Proof. $UP_rU^\dagger = P_r$ for every r , so $UAU^\dagger = \mathcal{A}$ record by record and (14) reads as claimed. $\square$

Lemma 2 (Transfer across a perfect correlation). *Let $\mathcal{A} = \{P_a\}$ and $\mathcal{B} = \{Q_b\}$ be released record algebras of parts A and B , let $S_A = \{a : (P_a \otimes I)\Psi \neq 0\}$ and $S_B = \{b : (I \otimes Q_b)\Psi \neq 0\}$, and suppose there is a bijection $f: S_A \rightarrow S_B$ with $(P_a \otimes Q_b)\Psi = 0$ unless $b = f(a)$. Then $\mu(a \mid \Psi; \mathcal{A}) = \mu(f(a) \mid \Psi; \mathcal{B})$ for every $a \in S_A$, and the records outside S_A and S_B have chance zero.*

Proof. By (15), $\mu(a \mid \Psi; \mathcal{A}) = \sum_b \mu(a, b \mid \Psi; \mathcal{A} \vee \mathcal{B})$, and by (H3) every term with $b \neq f(a)$ vanishes, so the sum is $\mu(a, f(a) \mid \Psi; \mathcal{A} \vee \mathcal{B})$. For $b = f(a)$ the only a' with $(P_{a'} \otimes Q_b)\Psi \neq 0$ is $a' = a$, because f is a bijection; so by (H3) again $\mu(a, f(a) \mid \Psi; \mathcal{A} \vee \mathcal{B}) = \sum_{a'} \mu(a', f(a) \mid \Psi; \mathcal{A} \vee \mathcal{B})$, which is $\mu(f(a) \mid \Psi; \mathcal{B})$ by the second half of (15). The last statement is (H3). $\square$

Lemma 3 (Equal weights). *Let $\mathcal{D}$ be a released record algebra of a part D , let $S \subset \Omega(\mathcal{D})$ have N elements, let the records $D_r = |d_r\rangle\langle d_r|$, $r \in S$, be rank one, and let*

$$\Psi = N^{-1/2} \sum_{r \in S} |d_r\rangle \otimes \Phi_r, \quad (16)$$

with the Φ_r orthonormal in the complementary factor; since the records of an algebra are mutually orthogonal, every record of $\mathcal{D}$ outside S annihilates Ψ . Then $\mu(r \mid \Psi; \mathcal{D}) = 1/N$ for $r \in S$, and every other record of $\mathcal{D}$ has chance zero.

Proof. For a permutation π of S let S_π^D be the unitary on $\mathcal{H}_D$ with $S_\pi^D |d_r\rangle = |d_{\pi(r)}\rangle$ and the identity on the orthogonal complement of the $|d_r\rangle$, and let S'_π be a unitary on the complementary factor with $S'_\pi \Phi_r = \Phi_{\pi(r)}$ and the identity on the orthogonal complement of the Φ_r . Then $U = S_\pi^D \otimes S'_\pi$ satisfies $U\Psi = \Psi$ and $U(D_r \otimes I)U^\dagger = D_{\pi(r)} \otimes I$ for $r \in S$, and U fixes every record of $\mathcal{D}$ outside S , since the records of an algebra are mutually orthogonal and S_π^D is the identity on the orthogonal complement of the $|d_r\rangle$; so $U\mathcal{D}U^\dagger = \mathcal{D}$ with the records of S permuted, and (14) gives $\mu(\pi(r) \mid \Psi; \mathcal{D}) = \mu(r \mid \Psi; \mathcal{D})$. Since π is arbitrary the N chances are equal; by (H3) every record outside S has chance zero; by (H2) the N chances sum to one. $\square$

The theorem

Theorem 9 (Born weights). *Under (H1)–(H6), for every released record algebra $\mathcal{A} = \{P_r\}$ of the network and every wave Ψ of the network,*

$$\mu(r \mid \Psi; \mathcal{A}) = \langle \Psi | P_r | \Psi \rangle. \quad (17)$$

Proof. Write $w_r = \langle \Psi | P_r | \Psi \rangle$ and $S = \{r : w_r > 0\}$, and for $r \in S$ put $\Phi_r = P_r \Psi / \sqrt{w_r}$; the Φ_r are orthonormal and Φ_r lies in the range of P_r .

Step 0: copy the record onto a pointer (H1, H3, H4, H5). The network holds a fresh pointer D in its initial state, so $\Psi = \Psi_0 \otimes |d_0\rangle$ with Ψ_0 on the rest, and by (H5) with every $m_r = 1$ a copying unitary $V = \sum_r P_r \otimes V_r$, $V_r |d_0\rangle = |d_r\rangle$, gives

$$\Psi' = V\Psi = \sum_{r \in S} \sqrt{w_r} \Phi_r \otimes |d_r\rangle, \quad (18)$$

with the Φ_r now read on the factor complementary to D . Since $[V, P_r \otimes I] = 0$, Lemma 1 gives $\mu(r \mid \Psi'; \mathcal{A}) = \mu(r \mid \Psi; \mathcal{A})$. The projectors $D_r = |d_r\rangle\langle d_r|$, $r \in S$, with the remainder form the released record algebra $\mathcal{D}$ of D (H5), and $\mathcal{A}$ and $\mathcal{D}$ are perfectly correlated on Ψ' : $(P_r \otimes D_{r'})\Psi' = 0$ unless $r' = r$, with the identity on S as the bijection. Lemma 2 gives $\mu(r \mid \Psi'; \mathcal{A}) = \mu(r \mid \Psi'; \mathcal{D})$. It remains to show $\mu(r \mid \Psi'; \mathcal{D}) = w_r$ for a pointer whose supported records are rank one and whose partners Φ_r are orthonormal.

Step 1: equal weights (H1, H2, H3). If $w_r = 1/N$ for $r \in S$, Lemma 3 gives $\mu(r \mid \Psi'; \mathcal{D}) = 1/N = w_r$.

Step 2: rational weights (H1–H5). Let $w_r = m_r/M$ with integers $m_r \geq 1$ for $r \in S$ and $M = \sum_{r \in S} m_r$. By (H5) applied to $\mathcal{D}$ with these integers, a pointer C and $W = \sum_r D_r \otimes W_r$ give

$$\Psi'' = W\Psi' = M^{-1/2} \sum_{r \in S} \sum_{j=1}^{m_r} \Phi_r \otimes |d_r\rangle |c_{rj}\rangle. \quad (19)$$

Since $[W, D_r \otimes I] = 0$, Lemma 1 gives $\mu(r \mid \Psi''; \mathcal{D}) = \mu(r \mid \Psi'; \mathcal{D})$. Let $\mathcal{C}$ be the released record algebra of C supplied by (H5), with records $C_{rj} = |c_{rj}\rangle\langle c_{rj}|$ and their remainder, so that $\mathcal{D} \vee \mathcal{C}$ is a released record algebra refining $\mathcal{D}$; on Ψ'' its supported records are the M rank-one projectors

$|d_r c_{rj}\rangle\langle d_r c_{rj}|$, each with weight $1/M$, and their partners Φ_r repeat. By (H5) with every integer equal to one, copy the refined record onto a further pointer E :

$$\Psi''' = M^{-1/2} \sum_{r \in S} \sum_{j=1}^{m_r} \Phi_r \otimes |d_r c_{rj}\rangle \otimes |e_{rj}\rangle, \quad (20)$$

with the $|e_{rj}\rangle$ orthonormal. The copy is block-diagonal in the records of $\mathcal{D} \vee \mathcal{C}$ and so commutes with every one of them, supported or not, and Lemma 1 gives the chances of $\mathcal{D} \vee \mathcal{C}$ the same value on Ψ''' as on Ψ'' ; and on Ψ''' the partners $\Phi_r \otimes |e_{rj}\rangle$ are orthonormal, so Lemma 3 gives $\mu(r, (r, j) \mid \Psi''; \mathcal{D} \vee \mathcal{C}) = 1/M$ for the M supported records $(r, (r, j))$ and zero for every other record of $\mathcal{D} \vee \mathcal{C}$. Finally, by the first half of (15) on the division of the network into D and the rest,

$$\mu(r \mid \Psi''; \mathcal{D}) = \sum_{(r', j)} \mu(r, (r', j) \mid \Psi''; \mathcal{D} \vee \mathcal{C}) = \sum_{j=1}^{m_r} \frac{1}{M} = \frac{m_r}{M} = w_r. \quad (21)$$

Step 3: all weights (H6). For fixed S , $\mathcal{D}$ and partners, $w \mapsto \Psi'(w) = \sum_{r \in S} \sqrt{w_r} \Phi_r \otimes |d_r\rangle$ is a continuous map from the closed simplex to the unit sphere, and Steps 1 and 2, applied with the support of w in place of S , give $\mu(r \mid \Psi'(w); \mathcal{D}) = w_r$ at every rational point of the simplex. Rational points are dense, and (H6) extends the identity to every w . If S is infinite, let Ψ'_N be the normalised sum over the N largest weights; the finite case gives $\mu(r \mid \Psi'_N; \mathcal{D}) = w_r / \sum_{S_N} w_s$, $\Psi'_N \rightarrow \Psi'$ in norm, and (H6) gives $\mu(r \mid \Psi'; \mathcal{D}) = w_r$. With Step 0 this is (17). $\square$

Corollary 4 (The path measure is forced). *For the record algebra of N slices, Theorem 9 gives $\mu(r \mid \Psi) = \langle \Psi | \Pi_r | \Psi \rangle$, which is the path measure (6); an initial system state ρ_S that is mixed is the reduction of the one wave on a larger network, and the chance is $\text{Tr}[\rho_0 \Pi_r]$ by linearity of the partial trace. The system-level chance of an unsharp record, $\text{Tr}[\rho_S E_r]$ with E_r the induced effect of Theorem 3, follows by the same partial trace: it is a consequence, not a rule. The event law of Theorem 7 is therefore forced rather than imported: the intensities (10) are the $\delta t \rightarrow 0$ limit of a discrete path measure whose weights the theorem fixes.*

Remark 6 (No signalling: premise and check). Section 9.1 obtained parameter independence from the instrument. Here the same property, for every division of the network and not only for causally disjoint regions, is the hypothesis (H4), placed on the law of chance before the weights are known; the law it forces, $\langle \Psi | P_a \otimes Q_b | \Psi \rangle$, satisfies (15) identically. The two statements are consistent and not circular: (H4) is the premise, and the no-signalling of the instrument is the check that the premise is honoured. Every no-signalling derivation of the weights takes the same order [100].

Remark 7 (What each premise excludes). The witnesses are in the third script named in the Data availability statement. (a) A rule that squares the weights and renormalises, $\mu_r \propto |c_r|^4$, satisfies normalisation and coarse-graining additivity within every record algebra to machine precision (2.2×10^{-16})^[W-A]: the within-algebra content of (H2) excludes no weighting. The rule is covariant and respects (H3), and it fails (H4): on random bipartite waves a part's marginal changes by up to 0.211 when the other part records^[W-B], and the fine-graining chain of Step 2, run with it for the weights (2/3, 1/3), returns 2/3 where the rule applied directly gives 0.8^[W-E]. (b) A calibrated degree-two rule on a qubit, extended to a pair as the symmetrised product of its two-copy effects, assigns chance 0.094 to a joint record of zero amplitude, but it is not covariant (0.466 against 0.709 under a relabelling) and falls to (H1) before (H3) is reached^[W-C]. (c) The symmetric cubic rule $\mu_r = w_r^3 + 3w_r^2(1 - w_r) + 2w_r e_2(w_{s \neq r})$, with e_2 the second elementary symmetric polynomial of the other weights, is normalised, covariant and calibrated, respects joint support, and fails (H4) by 0.082^[W-C']. (d) The no-signalling toy model of Galley and Masanes [104], not reproduced in the script, satisfies (H4); its outcomes are not sharp records, and as a law of chance it violates the calibration that (H2) and (H3) together imply, since no

outcome in it is ever certain. (e) A rule that treats rational and irrational weights differently fails (H6). So: non-covariant rules fall to (H1), covariant nonlinear rules to (H4), never-certain rules to (H2) with (H3), discontinuous rules to (H6).

Remark 8 (Gleason, additivity, and what is not derived). Gleason’s theorem [105, 106] characterises the same weights from a stronger noncontextuality premise, that a record’s chance is the same whatever apparatus produces it; here that noncontextuality is a consequence of Theorem 9, and the only noncontextuality the proof uses is (15). Additivity, which Zhang shows no derivation of the weights can do without [103], is supplied by (H2) within an algebra and by (H4) across algebras; both are commitments, and neither is claimed to follow from the record structure. The derivation is conditional on the six premises. What is not derived: that there is chance at all, which is the postulate of Section 8; the extension of no signalling to every division, (H4); the availability of pointers, (H5); and continuity, (H6). In the binary QND model of [1] the same weights also descend from the ensemble to single histories (its Proposition 2); that descent is now the check that the single-history law and the master equation agree, not the ground of the weights.

B Map of verification checks

label	script	statement verified
R1, R2	backbone	Theorem 1: commutativity, nondemolition, stability (10^{-15})
R3, R8	backbone	Theorem 3: effects commute under QND w.r.t. the recorded observable; converse false
R4	backbone	Theorem 2: 32 histories, spectral = path measure (10^{-16}), orthogonality
R5, R9	backbone	release cases: Markov, eraser, scrambling; (S)/(E) split with the enlarged algebra
R6	backbone	Theorem 8: exactness at $\alpha = 3, 8, 20$; Wiener increments; SME spread
R7	backbone	Theorem 4: restriction norm = reduced trace distance (200 pairs)
R10	backbone	family (i) of Table 1
S1a, S1b	stress test	Theorem 5: every axis released; same-slice readouts do not commute
S2a	stress test	Theorem 6: identity and bound, random and on the families
S2b	stress test	Proposition 1
S2c, S2d	stress test	Table 1; Proposition 2; cross-family inversions
S3a–S3d	stress test	Proposition 3
S4	stress test	Proposition 4
W-A, W-B, W-E	Born witnesses	Appendix A: the $ c ^4$ rule passes the within-algebra axioms (2.2×10^{-16}) and fails (H4) (0.211; chain 2/3 against 0.8)
W-C, W-C'	Born witnesses	Appendix A: the degree-two rule is not covariant (0.466 against 0.709); the cubic rule fails (H4) by 0.082
B1–B7	fermions	Proposition 5: one-body couplings, and all 119 spectral projectors of interacting symmetric Hamiltonians for $N = 2, 3$, commute with every permutation (4×10^{-12}); a labelled operator leaves the sector; Eq. (11); c_x keeps antisymmetry
E1–E3	fermions	Section 5: a hit of width one site on one label leaves the sector (51% kept at the worst placement; 96% at width four) and its projected probabilities sum to 0.62; the symmetric hit keeps it
C1–C7	fermions	Eq. (12) and the hole at $y = x$; the τ^2 refill; 40 000 sampled histories (0.79% against 0.81%); a second sampled ensemble reproduces the master equation; $G_2^{\text{bosons}} + G_2^{\text{fermions}} = 2G_2^{\text{labelled}}$ in two non-interacting orbitals
He1, He3	helium	leading natural-orbital occupation 0.992 per spin; on-top ratio below one for idealised local detectors (upper estimates in Gaussian bases)

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