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Quantum-Information Probes of Collapse Dynamics in Causal Fermion Systems

Kernel-resolved decoherence, correlation scaling, and fault tolerance

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Abstract

Finster, Kleiner, and Paganini derived from causal fermion systems (CFS) a nonrelativistic effective collapse model whose statistical operator obeys, to leading order, a Kossakowski–Lindblad equation. In the factorized regime of that derivation, the dissipator is built from purely spatial operators $N_{k,a}$ with positive rates fixed by microscopic time-nonlocal kernels. We translate this CFS-specific generator into quantities accessible to coherent quantum processors. The resulting CFS Dirichlet form obeys an exact purity-loss identity under arbitrary Hamiltonian control, and its null algebra is the commutant of the $*$-algebra generated by the projected $N_{k,a}$. For joint eigenbranches we obtain an exact distinguishability functional controlling exponential coherence loss and a lower bound on channel distance from ideal unitary evolution. In an n-qubit dephasing sector, the projected kernels generate a positive-semidefinite covariance matrix C, with GHZ visibility $\exp[-2T\mathbf{1}^{\mathsf{T}}C\mathbf{1}]$. We prove that uniformly summable finite-range correlations yield at most $O(n)$ GHZ decay, whereas $O(n^2)$ scaling requires a collective non-summable component. Single- and two-qubit cat coherences reconstruct C, giving a covariance-spectroscopy protocol, while Knill–Laflamme conditions characterize first-order correctability of the projected CFS error algebra. We also derive an explicit compatibility inequality between macroscopic state reduction and long coherent computation and show why absence of probe heating does not imply absence of coherence signatures. The results do not exclude CFS: its covariance strength, $\ell_{\min}$, and hardware projection remain underdetermined. They instead identify the microscopic input required for a genuine no-go theorem and establish scalable quantum processors as direct probes of the CFS collapse sector.

1 Introduction

Objective-collapse models replace exact linear quantum evolution by a dynamics that suppresses sufficiently macroscopic superpositions while leaving ordinary microscopic quantum mechanics nearly intact. The best-developed phenomenological example is continuous spontaneous localization (CSL) [1–4]. Because such modifications are dynamical rather than merely interpretive, they are in principle testable. Existing searches have therefore focused on interferometric loss of coherence, anomalous diffusion, spontaneous radiation, and other noninterferometric signatures [5–7].

Causal fermion systems (CFS) provide a different route to collapse. Rather than adding a phenomenological stochastic term to the Schrödinger equation, Finster, Kleiner, and Paganini derive an effective stochastic and nonlinear dynamics from the causal action principle in a nonrelativistic limit [8]. Their statistical operator obeys, to leading order, a deterministic Kossakowski–Lindblad equation. The same construction is nonlocal in time on a microscopic scale $\ell_{\min}$, and the strength of

the stochastic background remains an effective parameter. A later analysis by Finster and Paganini argues that this collapse mechanism does not produce probe heating in the statistical mean, up to higher-order corrections [9]. If that conclusion persists in the relevant experimental regimes, the strongest heating- and radiation-based strategies developed for CSL cannot simply be transferred to CFS.

Quantum-information processors are therefore natural probes of the remaining effect: loss of coherent phase information. This point has precedent. Vischi, Ferialdi, Trombettoni, and Bassi analyzed CSL on superconducting transmon processors and found that direct collapse-induced suppression of computational superpositions is small for standard CSL parameters, while an indirect quasiparticle-production channel can become more important for large devices [10]. CFS changes that hierarchy if its no-heating mechanism removes the indirect channel. Then direct coherence, many-body correlation scaling, and logical error rates become comparatively more informative observables.

The relevant issue is *not* whether a microscopic CFS spacetime is discrete and therefore unable to support quantum algorithms. The general CFS formalism begins with a complex Hilbert space, allows general entangled states in an appropriate limiting construction [11], and now has limiting regimes that recover unitary Fock-space QED and perturbative non-Abelian gauge theory [12, 13]. The sharper question is whether the *collapse correction* can remain sufficiently weak, local, degenerate on encoded subspaces, or otherwise correctable to coexist with scalable coherent computation.

This paper develops that question directly from the CFS collapse kernels. The main point of departure from a generic Lindblad analysis is that the operators and rates are inherited from the factorization used in the CFS derivation. This produces a concrete positive quadratic form on processor observables. We use it to establish five results.

First, the leading CFS dissipator defines a Dirichlet form whose value is exactly the rate at which coherent purity is lost. The identity survives arbitrary Hamiltonian control. Second, the null algebra of this form is the commutant of the algebra generated by the CFS collapse operators, so passive decoherence-free quantum information is characterized algebraically. Third, when the CFS operators are projected onto a dephasing-dominated qubit register, their coefficients generate a positive-semidefinite correlation matrix C. This matrix determines the scaling of cat-state decoherence and can be reconstructed experimentally with one- and two-qubit coherences. Fourth, standard Knill–Laflamme theory gives exact first-order conditions for active correction of the same CFS error algebra. Fifth, the simultaneous requirements of macroscopic collapse and long coherent computation define an explicit compatibility window. Closing that window would constitute a genuine CFS no-go theorem; the presently published CFS analysis does not yet fix the microscopic quantities tightly enough to do so.

Throughout we analyze the leading nonrelativistic statistical equation of Ref. [8]. We do not claim that the full time-nonlocal CFS dynamics is itself a Markov semigroup, nor do we identify ordinary laboratory noise with CFS noise. The results should instead be read as conditional consequences of the effective collapse generator already derived within CFS.

2 From the CFS kernels to a positive decoherence generator

The 2024 CFS collapse construction starts from a nonlocal stochastic potential. After diagonalizing the covariance and imposing the factorization used in Sec. 4.4 of Ref. [8], the spatial operators take the form

$$M_{k,a}(t,t') = \Delta_{k,a}(t'-t)\, N_{k,a}, \tag{1}$$

where $N_{k,a}$ is a purely spatial symmetric operator and $\Delta_{k,a}$ is a compactly supported temporal kernel. Under the sign assumption adopted there, the symmetric temporal combinations entering the collapse coefficient are non-negative. In the notation of Eqs. (4.19)–(4.20) of Ref. [8], these combinations may be written

$$a_{k,a}(\zeta) := \Delta_{k,a}(2\zeta) + \Delta_{k,a}(-2\zeta), \tag{2}$$

$$b_{k,a}(\zeta) := \int_{-\infty}^{-\zeta} \left[\Delta_{k,a}(-2\nu) + \Delta_{k,a}(2\nu)\right] \mathrm{d}\nu, \tag{3}$$

with $a_{k,a}(\zeta), b_{k,a}(\zeta) \geq 0$ in the simplified regime considered there.

Theorem 4.1 of Ref. [8] then gives, to leading order, the Kossakowski–Lindblad equation

$$\dot{\rho} = -i[H, \rho] - \frac{1}{2} \sum_{k,a} \int_{-\infty}^{\infty} \mathrm{d}\zeta \, [K_{k,a,\zeta}, [K_{k,a,\zeta}, \rho]], \tag{4}$$

where Eq. (4.18) of that reference has the factorized form

$$K_{k,a,\zeta} = \kappa_{k,a}(\zeta) \, N_{k,a}. \tag{5}$$

The non-negative coefficient $|\kappa_{k,a}(\zeta)|^2$ is fixed there by the temporal factors in Eqs. (2)–(3). We denote it by

$$q_{k,a}(\zeta) := |\kappa_{k,a}(\zeta)|^2 \geq 0. \tag{6}$$

This notation deliberately takes the normalized $K_{k,a,\zeta}$ of the published Lindblad equation as primary, so that none of the processor bounds below depends on a convention-dependent redistribution of numerical factors between $q_{k,a}$ and $N_{k,a}$.

Theorem 4.1 quotes a relative correction of order $O(\ell_{\min}\|B\|)$ to the leading equation; Sec. 4.5.4 separately shows that the subset of non-adjacent Gaussian pairings is suppressed by $O(\ell_{\min}^2\|B\|^2)$. All exact identities below refer to the leading generator in Eq. (4). They remain useful in the full theory only to the extent that these neglected terms are controlled in the relevant processor regime.

Because $N_{k,a}$ does not depend on ζ in the factorized ansatz, the temporal kernel can be compressed into the positive rate

$$\lambda_{k,a} := \frac{1}{2} \int_{-\infty}^{\infty} q_{k,a}(\zeta) \, \mathrm{d}\zeta \geq 0. \tag{7}$$

Assuming the resulting sum converges on the finite-energy sector of interest,

$$\dot{\rho} = -i[H, \rho] + \mathcal{D}_{\mathrm{CFS}}(\rho), \qquad \mathcal{D}_{\mathrm{CFS}}(\rho) := -\sum_{k,a} \lambda_{k,a} [N_{k,a}, [N_{k,a}, \rho]]. \tag{8}$$

Equivalently, the Lindblad operators in the compressed representation are $L_{k,a} = \sqrt{2\lambda_{k,a}}\, N_{k,a}$.

Equation (8) is the object a quantum processor sees at the statistical level. The microscopic non-locality in time has not disappeared conceptually: it determines the positive coefficients $\lambda_{k,a}$ through the CFS temporal kernels. But in the leading factorized ensemble dynamics its experimentally relevant action is encoded by a completely positive quadratic decoherence form.

2.1 Restriction to a processor subspace

Let P project onto a finite-dimensional computational or code subspace $\mathcal{H}_Q$. If

$$(\mathbb{1} - P) N_{k,a} P = 0 \tag{9}$$

for the operators relevant on the experimental time scale, the leading collapse dynamics closes on $\mathcal{H}_Q$ and one may replace $N_{k,a}$ by $PN_{k,a}P$. If Eq. (9) fails, CFS collapse also produces leakage out of the chosen processor subspace. Such leakage is itself experimentally constrained, but it is not described by the pure-dephasing formulas of Secs. 4–5. We therefore make Eq. (9), exact or to the accuracy being tested, an explicit assumption whenever a closed qubit reduction is used.

Table 1 summarizes the resulting translation.

Table 1: CFS collapse data and their processor-level images in the leading factorized regime.

CFS quantity	Processor image	Direct quantum-information probe
$N_{k,a}$	Hermitian error generator on $\mathcal{H}_Q$	basis dependence, leakage, logical action
$\lambda_{k,a}$	positive collapse rate from the temporal kernels	Ramsey/cat coherence rate
$\mathrm{alg}^*\{N_{k,a}\}$	CFS noise algebra	decoherence-free and correctable subspaces
projected coefficients of $N_{k,a}$	positive covariance matrix C	one-/two-qubit covariance spectroscopy
$\ell_{\min}$	microscopic nonlocality scale entering the kernels	spatial/temporal correlation structure, if inherited by the processor projection

3 A CFS Dirichlet form and purity loss under arbitrary control

The dissipator in Eq. (8) defines a natural positive quadratic form on operators.

Definition 1 (CFS decoherence form)**.** For a Hilbert–Schmidt operator X on $\mathcal{H}_Q$, define

$$\mathcal{E}_{\mathrm{CFS}}(X) := \sum_{k,a} \lambda_{k,a} \, \|[N_{k,a}, X]\|_{\mathrm{HS}}^2 \, . \tag{10}$$

This form is entirely determined by the CFS spatial kernels and the positive rates inherited from their temporal nonlocality.

Theorem 2 (Exact purity-loss identity)**.** *Let $\rho(t)$ obey Eq. (8), allowing an arbitrary time-dependent Hermitian control Hamiltonian $H(t)$. Then*

$$\frac{\mathrm{d}}{\mathrm{d}t} \operatorname{Tr} \rho(t)^2 = -2 \, \mathcal{E}_{\mathrm{CFS}}(\rho(t)) = -2 \sum_{k,a} \lambda_{k,a} \, \|[N_{k,a}, \rho(t)]\|_{\mathrm{HS}}^2 \leq 0. \tag{11}$$

Proof. Using $\frac{\mathrm{d}}{\mathrm{d}t} \operatorname{Tr} \rho^2 = 2 \operatorname{Tr}(\rho \dot{\rho})$, the Hamiltonian contribution vanishes by cyclicity of the trace. For Hermitian N,

$$\operatorname{Tr}\big(\rho[N, [N, \rho]]\big) = \|[N, \rho]\|_{\mathrm{HS}}^2 \, . \tag{12}$$

Substitution into Eq. (8) gives Eq. (11). □

For a pure state $\rho = |\psi\rangle\langle\psi|$,

$$\|[N, \rho]\|_{\mathrm{HS}}^2 = 2\,\mathrm{Var}_\psi(N), \tag{13}$$

so Theorem 2 yields

$$\frac{\mathrm{d}}{\mathrm{d}t}\,\mathrm{Tr}\,\rho^2\Big|_{\rho=|\psi\rangle\langle\psi|} = -4\sum_{k,a}\lambda_{k,a}\,\mathrm{Var}_\psi(N_{k,a}). \tag{14}$$

Thus a pure processor state begins losing purity precisely to the extent that it is a superposition across eigenspaces resolved by the projected CFS operators.

The role of coherent control is worth emphasizing. In an interaction picture generated by the intended unitary $U(t)$, the collapse operators rotate as

$$\widetilde{N}_{k,a}(t) = U(t)^\dagger N_{k,a} U(t), \tag{15}$$

while Eq. (11) remains a sum of non-negative terms. Hamiltonian control can move the state into more or less sensitive sectors and can average structured noise, but it cannot make already incurred statistical purity loss cancel with an opposite Hamiltonian contribution. This is distinct from active quantum error correction, which uses ancillary degrees of freedom, syndrome extraction, and recovery operations.

3.1 Noise algebra and passive decoherence-free information

Let

$$\mathcal{A}_{\mathrm{CFS}} := \mathrm{alg}^*\{N_{k,a} : \lambda_{k,a} > 0\} \tag{16}$$

be the unital $*$-algebra generated by the active projected CFS operators.

Theorem 3 (Null algebra of the CFS dissipator)**.** *On a finite-dimensional processor subspace,*

$$\mathcal{E}_{\mathrm{CFS}}(X) = 0 \quad\Longleftrightarrow\quad X \in \mathcal{A}'_{\mathrm{CFS}} := \{X : [X, A] = 0 \text{ for all } A \in \mathcal{A}_{\mathrm{CFS}}\}. \tag{17}$$

Consequently the operators immune to the leading CFS dissipator form the commutant of the CFS noise algebra.

Proof. Every term in Eq. (10) is non-negative. Hence $\mathcal{E}_{\mathrm{CFS}}(X) = 0$ if and only if $[N_{k,a}, X] = 0$ for every active generator. Commuting with all generators is equivalent to commuting with the unital $*$-algebra they generate. □

Theorem 3 gives a sharper question than "does CFS noise exist?" If $\mathcal{A}_{\mathrm{CFS}}$ acts irreducibly on a proposed register, its commutant is only the scalars and there is no nontrivial passive decoherence-free operator algebra. If the representation has degeneracies or multiplicity factors, nontrivial protected subsystems can exist. The structure of the actual $N_{k,a}$ on a laboratory register is therefore phenomenologically decisive.

4 Kernel-resolved branch distinguishability

For an idling or interferometric interval, suppose $|u\rangle$ and $|v\rangle$ are simultaneous eigenstates of the relevant spatial CFS operators,

$$N_{k,a}|u\rangle = n_{k,a}^{(u)}|u\rangle, \qquad N_{k,a}|v\rangle = n_{k,a}^{(v)}|v\rangle. \tag{18}$$

Define the CFS branch-distinguishability functional

$$\mathfrak{D}_{\mathrm{CFS}}(u,v) := \sum_{k,a} \lambda_{k,a} \left(n_{k,a}^{(u)} - n_{k,a}^{(v)} \right)^2 . \tag{19}$$

Equivalently, substituting Eq. (7),

$$\mathfrak{D}_{\mathrm{CFS}}(u,v) = \sum_{k,a} \int \mathrm{d}\zeta \, w_{k,a}(\zeta) \left(n_{k,a}^{(u)} - n_{k,a}^{(v)} \right)^2 . \tag{20}$$

This formula retains the microscopic CFS origin of the rate: a branch is fragile exactly when the nonlocal-kernel operators resolve it strongly.

Proposition 4 (Exact branch-coherence decay)**.** *Assume Eq. (18) and that the coherent Hamiltonian is diagonal on the two-dimensional span of $|u\rangle, |v\rangle$ during the interrogation interval. Then*

$$\rho_{uv}(T) = e^{-i\phi_{uv}(T)} e^{-T\mathfrak{D}_{\mathrm{CFS}}(u,v)} \rho_{uv}(0), \tag{21}$$

where ϕ_{uv} is the usual Hamiltonian relative phase.

Proof. For every $N_{k,a}$,

$$\langle u|[N_{k,a},[N_{k,a},\rho]]|v\rangle = \left(n_{k,a}^{(u)} - n_{k,a}^{(v)} \right)^2 \rho_{uv}. \tag{22}$$

The corresponding scalar first-order differential equation integrates to Eq. (21). □

For the balanced cat state $|+_{uv}\rangle = (|u\rangle + |v\rangle)/\sqrt{2}$, the ideal output after removing the known Hamiltonian phase has unit visibility, whereas the CFS channel has visibility

$$V_{uv}(T) = e^{-T\mathfrak{D}_{\mathrm{CFS}}(u,v)}. \tag{23}$$

This immediately gives a channel-level witness.

Corollary 5 (Distance from ideal unitary evolution)**.** *Let Φ_T^{CFS} denote the channel generated on the two-branch subspace by Eq. (8) and $\mathcal{U}_T$ the corresponding ideal unitary channel. With the normalized diamond distance*

$$d_\diamond(\Phi,\Psi) := \frac{1}{2} \left\| \Phi - \Psi \right\|_\diamond , \tag{24}$$

we have

$$d_\diamond(\Phi_T^{\mathrm{CFS}}, \mathcal{U}_T) \geq \frac{1 - e^{-T\mathfrak{D}_{\mathrm{CFS}}(u,v)}}{2}. \tag{25}$$

Proof. Evaluate the channel difference on $|+_{uv}\rangle\langle+_{uv}|$. The ideal and dephased outputs differ only in the two off-diagonal entries by $(1 - e^{-T\mathfrak{D}})/2$. The trace distance of these outputs is $(1 - e^{-T\mathfrak{D}})/2$, which lower-bounds the normalized diamond distance. □

The statistical result is logically distinct from the state-vector reduction argument in Ref. [8]. In Sec. 4.6 of that work, the collapse proof assumes that the measured observable commutes with the operators describing the dynamics and that the measurement scale is much larger than $\ell_{\min}$. Equation (11), by contrast, is a statement about the leading ensemble master equation and does not require a measurement observable commuting with an externally applied control Hamiltonian. This makes controlled processors useful even in regimes where a complete CFS trajectory-level collapse theorem has not yet been established.

5 Projected qubit covariance and many-body scaling

We now specialize to a dephasing-dominated qubit projection. On n physical qubits, suppose the branch-resolving part of each CFS operator can be written

$$PN_{k,a}P = c_{k,a}^{(0)}\mathbb{1} + \sum_{i=1}^{n} z_{k,a;i} Z_i, \tag{26}$$

where Z_i is a Pauli Z operator. Identity terms drop out of commutators. Define

$$C_{ij} := 2\sum_{k,a} \lambda_{k,a}\, z_{k,a;i} z_{k,a;j}. \tag{27}$$

Then the projected dissipator is

$$\mathcal{D}_Z(\rho) = -\frac{1}{2}\sum_{i,j=1}^{n} C_{ij}[Z_i,[Z_j,\rho]]. \tag{28}$$

Proposition 6 (CFS covariance is positive semidefinite)**.** *The matrix C in Eq. (27) is positive semidefinite.*

Proof. For any real vector $x \in \mathbb{R}^n$,

$$x^{\mathsf{T}} C x = 2\sum_{k,a} \lambda_{k,a} \left(\sum_i x_i z_{k,a;i}\right)^2 \geq 0. \tag{29}$$

Thus C is a Gram matrix of the projected CFS kernel coefficients. □

Let $|x\rangle, |y\rangle$ be computational-basis states and let $s_i(x) \in \{+1,-1\}$ denote the eigenvalue of Z_i on $|x\rangle$. Set $d_i = s_i(x) - s_i(y)$. Equation (28) gives the exact branch rate

$$\Gamma(x,y) = \frac{1}{2}\,\boldsymbol{d}^{\mathsf{T}} C \boldsymbol{d}. \tag{30}$$

For the GHZ branches $|0^n\rangle$ and $|1^n\rangle$, $\boldsymbol{d} = 2\mathbf{1}$, so

$$V_{\mathrm{GHZ}}(T) = \exp\left[-2T\,\mathbf{1}^{\mathsf{T}} C \mathbf{1}\right]. \tag{31}$$

Two familiar scaling regimes follow immediately:

$$C = \gamma_{\mathrm{loc}} I \quad\Longrightarrow\quad V_{\mathrm{GHZ}} = e^{-2\gamma_{\mathrm{loc}} nT}, \tag{32}$$

$$C = \gamma_{\mathrm{col}} \mathbf{1}\mathbf{1}^{\mathsf{T}} \quad\Longrightarrow\quad V_{\mathrm{GHZ}} = e^{-2\gamma_{\mathrm{col}} n^2 T}. \tag{33}$$

The first is independent local dephasing; the second is a fully collective rank-one component. Equation (27) identifies this n versus n^2 distinction directly through the Gram matrix of the projected CFS collapse kernels.

5.1 A finite-range scaling theorem

The localization scale $\ell_{\min}$ in CFS motivates a precise question: if the processor projection inherits only finite-range or absolutely summable correlations, how fast can GHZ collapse grow with system size?

Theorem 7 (Uniformly summable correlations are at most extensive)**.** *Consider a sequence of n-qubit covariance matrices $C^{(n)} \succeq 0$. If there is a constant c_* independent of n such that*

$$\max_i \sum_{j=1}^{n} |C_{ij}^{(n)}| \le c_*, \tag{34}$$

then

$$0 \le \mathbf{1}^{\mathsf{T}} C^{(n)} \mathbf{1} \le n c_*, \qquad \Gamma_{\mathrm{GHZ}}^{(n)} \le 2 n c_*. \tag{35}$$

Consequently superextensive GHZ decay requires failure of a uniform absolute-summability bound such as Eq. (34).

Proof. Positive semidefiniteness gives $\mathbf{1}^{\mathsf{T}} C \mathbf{1} \ge 0$. For the upper bound,

$$\mathbf{1}^{\mathsf{T}} C \mathbf{1} = \sum_{i,j} C_{ij} \le \sum_{i,j} |C_{ij}| \le n \max_i \sum_j |C_{ij}| \le n c_*. \tag{36}$$

Substituting into Eq. (31) yields Eq. (35). □

A useful corollary is obtained when qubits occupy a fixed-density D-dimensional architecture and

$$|C_{ij}| \le c_0 f\left(\frac{r_{ij}}{\xi}\right), \qquad \sup_i \sum_j f\left(\frac{r_{ij}}{\xi}\right) < \infty. \tag{37}$$

Then the GHZ exponent grows at most linearly in n. If the CFS projection inherits a finite correlation range controlled by $\ell_{\min}$, Eq. (37) is the natural processor-level hypothesis and quadratic scaling is not expected. This last implication is conditional: Ref. [8] states that the microscopic fields are localized and the effective dynamics is nonlocal on the scale $\ell_{\min}$, but it does not derive the qubit covariance matrix C for any present hardware platform.

The theorem also explains why scalable quantum computation does not automatically conflict with the CFS amplification argument. Ref. [8] states that a many-particle Hartree–Fock state collapses once one constituent one-particle wave function collapses, giving the usual amplification toward mesoscopic and macroscopic systems. If the relevant processor noise is local, both macroscopic collapse and GHZ decoherence can scale extensively while the enormous difference between the number of microscopic constituents in an apparatus and the number of controlled qubits leaves a broad possible hierarchy of time scales.

6 A quantum processor as a covariance spectrometer

The projected CFS matrix C is not merely a theoretical bookkeeping device. In the dephasing sector it can be reconstructed from simple coherent witnesses.

For a single qubit i, the coherence between $|0\rangle_i$ and $|1\rangle_i$ decays at rate

$$\Gamma_i = 2 C_{ii}. \tag{38}$$

For a pair i, j, consider the two cat coherences

$$|00\rangle\langle 11| : \qquad \Gamma_{ij}^{(+)} = 2(C_{ii} + C_{jj} + 2C_{ij}), \tag{39}$$

$$|01\rangle\langle 10| : \qquad \Gamma_{ij}^{(-)} = 2(C_{ii} + C_{jj} - 2C_{ij}). \tag{40}$$

Therefore

$$C_{ij} = \frac{\Gamma_{ij}^{(+)} - \Gamma_{ij}^{(-)}}{8}, \qquad C_{ii} = \frac{\Gamma_i}{2}. \tag{41}$$

Proposition 8 (Pair-cat covariance reconstruction)**.** *Within the closed dephasing reduction Eq. (28), the complete symmetric matrix C can be reconstructed from n single-qubit coherence rates and the two pair-cat rates in Eqs. (39)–(40) for each pair.*

This is a simple but useful experimental translation of the CFS kernel data. Varying the physical separation and geometry of the qubits turns Eq. (41) into a correlation-length probe. Repeating the construction under basis rotations probes components of the projected $N_{k,a}$ beyond the pure Z sector. If multi-qubit Pauli terms are appreciable in $PN_{k,a}P$, higher-order cat coherences or process tomography are required; the pair protocol should then be regarded as a diagnostic of the one-body dephasing block, not the complete CFS channel.

There is an unavoidable identification problem: conventional environmental dephasing generates the same mathematical covariance matrix. Thus a measured C is not evidence for CFS. The CFS contribution would have to be isolated by universality across devices, a distinctive spatial or temporal scaling, a lower floor that survives improved materials and control, or an independent theoretical prediction tying C to the same microscopic parameters that enter other CFS limits. Nevertheless, a processor places a direct upper bound on any additional positive CFS contribution compatible with the assumed additive dephasing model.

7 No heating does not imply no coherent signature

Finster and Paganini argue that the CFS collapse mechanism does not heat the probe in the statistical mean, up to higher-order corrections [9]. That result is important because many of the strongest collapse-model bounds exploit diffusion, spontaneous radiation, or energy deposition [5–7]. But energy conservation and coherence conservation are logically different.

For the leading generator Eq. (8), let H be time independent. Then

$$\frac{\mathrm{d}}{\mathrm{d}t}\,\mathrm{Tr}(H\rho) = -\sum_{k,a} \lambda_{k,a}\,\mathrm{Tr}(\rho[N_{k,a}, [N_{k,a}, H]])\,. \tag{42}$$

In particular, if $[N_{k,a}, H] = 0$ for all active kernels,

$$\frac{\mathrm{d}}{\mathrm{d}t}\,\mathrm{Tr}(H\rho) = 0, \tag{43}$$

while Theorem 2 still gives strict purity loss whenever $[N_{k,a}, \rho] \neq 0$.

Proposition 9 (Silent dephasing)**.** *A double-commutator CFS channel can preserve the mean energy of a probe while irreversibly suppressing coherence. Therefore a no-heating statement does not remove direct interferometric or processor-level signatures of the collapse channel.*

The proposition is *not* a derivation of the more subtle no-heating mechanism of Ref. [9]; that work uses the field-dependent scalar product and the full structure of the CFS interaction. Its role here is operational. Even if the full CFS cancellations eliminate net heating, a nonzero statistical decoherence form Eq. (10) remains directly testable through coherent evolution.

This is the principal reason quantum processors are complementary to noninterferometric collapse searches. They need not infer collapse from a secondary energy channel: they interrogate the phase information that the effective collapse dynamics is designed to erase.

8 Fault tolerance and the projected CFS error algebra

The existence of nonzero CFS decoherence does not by itself rule out scalable quantum computation. Quantum error correction is specifically designed to protect coherence against weak physical noise, including dephasing.

Over a short interval δt, Eq. (8) has Lindblad operators

$$L_{k,a} = \sqrt{2\lambda_{k,a}}\, N_{k,a}. \tag{44}$$

To first order in δt, the Kraus error span is generated by $\mathbb{1}$ and the active $N_{k,a}$. Let P_L project onto a code space.

Proposition 10 (First-order exact correction of CFS collapse errors)**.** *For a finite active set of CFS generators, the leading short-time collapse errors are exactly correctable on $P_L\mathcal{H}_Q$ if and only if the Knill–Laflamme relations for the error span $\{\mathbb{1}, N_{k,a}\}$ hold. Equivalently, there exist scalars $\alpha_{k,a}$ and $\beta_{(k,a),(l,b)}$ such that*

$$P_L N_{k,a} P_L = \alpha_{k,a} P_L, \tag{45}$$

$$P_L N_{k,a} N_{l,b} P_L = \beta_{(k,a),(l,b)} P_L \tag{46}$$

for all active pairs. For a continuous or countable family, the same condition applies to the closed error span on the code subspace.

Proof. The short-time channel has jump Kraus operators $\sqrt{\delta t}\, L_{k,a}$ and a no-jump Kraus operator equal to the identity plus $O(\delta t)$. Applying the Knill–Laflamme theorem to the resulting error span gives Eqs. (45)–(46); the positive factors $\sqrt{2\lambda_{k,a}}$ can be absorbed into the scalar coefficient matrix [14]. □

The passive criterion of Theorem 3 and the active criterion of Proposition 10 are complementary. A code contained in a joint degenerate eigenspace of all active $N_{k,a}$ is passively immune to the leading dissipator. More general codes may actively correct a finite error algebra even when the individual physical qubits decohere.

This is also why an asymptotic "BQP no-go" does not follow from the existence of CFS collapse. Threshold theorems already cover sufficiently weak local non-Markovian noise [15] and some algebraically decaying long-range correlated noise [16]. Whether CFS lies inside or outside such correctable classes depends on the norm, spatial decay, temporal correlations beyond the leading Markovian reduction, and logical action of the actual $N_{k,a}$. These are theory-specific quantities still to be fixed.

A particularly sharp failure mode would be a processor-visible component of $\mathcal{A}_{\mathrm{CFS}}$ that acts as an unavoidable logical operator at every code distance or a collective covariance whose strength

does not decrease under encoding. Establishing such a statement from the causal action principle would be far closer to a true quantum-computational no-go theorem than merely observing that the effective theory contains a Lindblad term.

9 Collapse–computation compatibility and a no-go criterion

The central phenomenological tension can be stated without committing to a specific hardware platform. A viable collapse model must make sufficiently distinct macroscopic branches decohere rapidly, while allowing experimentally realized microscopic and mesoscopic coherent branches to remain coherent long enough for controlled processing.

Suppose, along a one-parameter slice of CFS parameter space, that the positive rates factor as

$$\lambda_{k,a} = g\,\bar{\lambda}_{k,a}, \qquad g > 0, \tag{47}$$

where g is an overall stochastic-covariance strength and the barred quantities fix the normalized kernel shape for a chosen $\ell_{\min}$ and background. Define the normalized branch functionals $\bar{\mathfrak{D}}_{\mathrm{M}}$ for a macroscopic measurement branch pair and $\bar{\mathfrak{D}}_{\mathrm{Q}}$ for a quantum-computational witness.

If a measurement requires residual coherence at most ε_{M} after time T_{M},

$$g\,\bar{\mathfrak{D}}_{\mathrm{M}} T_{\mathrm{M}} \geq \ln\frac{1}{\varepsilon_{\mathrm{M}}}. \tag{48}$$

If a processor requires visibility at least η_{Q} after coherent time T_{Q},

$$g\,\bar{\mathfrak{D}}_{\mathrm{Q}} T_{\mathrm{Q}} \leq \ln\frac{1}{\eta_{\mathrm{Q}}}. \tag{49}$$

The two requirements are compatible only if

$$\boxed{\frac{\bar{\mathfrak{D}}_{\mathrm{M}}}{\bar{\mathfrak{D}}_{\mathrm{Q}}} \geq \frac{T_{\mathrm{Q}}}{T_{\mathrm{M}}}\frac{\ln(1/\varepsilon_{\mathrm{M}})}{\ln(1/\eta_{\mathrm{Q}})}} \tag{50}$$

for the chosen kernel shape and $\ell_{\min}$.

Equation (50) is the cleanest route to a genuine no-go theorem. If a microscopic CFS calculation were to fix the normalized kernels strongly enough that the inequality fails for *every* parameter choice capable of producing the required macroscopic state reduction, the corresponding collapse model would be incompatible with the observed coherent processor. Conversely, if Eq. (50) holds with a large margin, quantum computation does not threaten the collapse mechanism.

At present the published CFS collapse model explicitly retains $\ell_{\min}$ and the stochastic covariance strength as free parameters [8]. The recent QFT-limit program uses related stochastic data and therefore offers a plausible route to reducing this freedom [12,13], but no published relation currently supplies the lower bound on processor-visible $\bar{\mathfrak{D}}_{\mathrm{Q}}$ needed to close Eq. (50). The correct present conclusion is therefore a constraint program, not an exclusion theorem.

10 Experimental program

The preceding results suggest a processor-based CFS phenomenology program with increasing levels of model dependence.

Level 1: direct coherence bounds. For any branch pair to which Proposition 4 applies, an observed visibility $V(T)$ implies

$$\mathfrak{D}_{\mathrm{CFS}}(u,v) \leq -\frac{1}{T}\ln V(T) \tag{51}$$

if all observed loss is conservatively attributed to CFS. Ordinary noise makes this bound weak but safe: known environmental channels can only tighten it once independently modeled.

Level 2: spatial covariance spectroscopy. Use Eqs. (38)–(41) on qubit pairs at varying physical separations. This reconstructs the dephasing block of C and tests whether the residual covariance is local, finite range, algebraically decaying, or collective. If CFS predicts a processor projection tied to $\ell_{\min}$, the measured spatial decay can be compared directly with that prediction.

Level 3: many-body scaling. Measure GHZ or cat-state visibility as n increases while controlling geometry. Linear scaling is compatible with a uniformly summable finite-range CFS covariance; superlinear scaling would point to growing row sums, and an n^2 term to a collective component. This is more informative than a single coherence-time measurement because it interrogates the correlation structure of the collapse noise.

Level 4: logical scaling. Repeat the experiment on encoded logical qubits as code distance increases. If the physical CFS operators satisfy Eqs. (45)–(46) approximately and the residual noise is sufficiently local, logical errors should fall under appropriate fault-tolerant protocols. A code-distance-independent dephasing floor would instead indicate a logical or collective component not removed by the encoding.

Level 5: cross-platform universality. The strongest evidence for a fundamental collapse channel would not be an anomalous rate on one device but a common residual law across physically different platforms after conventional noise is reduced. Superconducting circuits are particularly relevant because objective-collapse constraints have already been studied there [10]; trapped ions, neutral atoms, spin qubits, and matter-wave interferometers would supply complementary spatial and material couplings.

The present paper does not attach numerical exclusions to existing hardware because CFS has not yet provided the map $N_{k,a} \mapsto$ hardware operators or a fixed stochastic covariance. Inventing such a map would convert a rigorous constraint into a model-dependent estimate. The purpose here is instead to specify exactly what must be calculated before processor data can become a quantitative CFS test.

11 Discussion

The effective-collapse program in CFS occupies an unusual experimental position. On one hand, the leading statistical dynamics is rigid enough to produce a positive Kossakowski–Lindblad decoherence form. On the other, the microscopic scale $\ell_{\min}$, covariance strength, and processor matrix elements of the spatial kernels remain underdetermined. The theory can therefore be constrained structurally before it can be excluded numerically.

The strongest CFS-specific object isolated here is not a generic "decoherence rate" but the operator family

$$\{\lambda_{k,a}, N_{k,a}\}, \tag{52}$$

with $\lambda_{k,a}$ inherited from the temporal kernels of the nonlocal stochastic potential. From this family follow three experimentally distinct structures:

(i) the Dirichlet form $\mathcal{E}_{\mathrm{CFS}}$, which measures instantaneous purity loss;

(ii) the noise algebra $\mathcal{A}_{\mathrm{CFS}}$, which determines passive protected information and enters active error correction;

(iii) the projected covariance matrix C, which controls the spatial scaling of dephasing on a qubit register.

These structures are more informative than a single phenomenological collapse constant.

The analysis also makes clear why the simple "discrete substrate versus BQP" objection is misplaced. CFS is not a finite-state classical cellular automaton, and its QFT-limit work explicitly constructs quantum states, entanglement, and unitary Fock-space dynamics [11, 12]. The potentially testable tension lies instead between the unitary QFT limit and the nonunitary collapse correction. Quantum computation is useful because it deliberately operates near that interface for long times and with controllable many-body entanglement.

There are four theoretical inputs that would substantially sharpen the present bounds.

First, one needs an explicit evaluation of the spatial operators $N_{k,a}$ for realistic low-energy matter and electromagnetic environments, followed by projection onto actual qubit degrees of freedom. Second, the stochastic covariance should be tied quantitatively to the same microscopic parameters entering the CFS QFT limit, reducing the freedom in the rates $\lambda_{k,a}$. Third, one needs control of the $O(\ell_{\min}\|B\|)$ corrections to the leading statistical equation for long coherent protocols. Fourth, the full time-nonlocal and eventually relativistic collapse dynamics should be analyzed under externally driven noncommuting controls rather than only in a measurement setting satisfying the commutation assumptions of the existing state-reduction proof.

If these steps produced a nonzero lower bound on the CFS distinguishability functional for a family of increasingly large fault-tolerant processors, Eq. (50) could become an experimental exclusion criterion. If they instead showed that the collapse operators act almost scalarly on encoded computational sectors while strongly resolving macroscopic apparatus branches, CFS would explain how objective collapse and scalable quantum computation coexist. Either outcome would be physically informative.

12 Conclusion

We have translated the leading effective collapse dynamics derived from causal fermion systems into a quantum-information language that retains its microscopic kernel structure. The positive weights generated by the time-nonlocal CFS kernels define a Dirichlet form, an operator noise algebra, and, after qubit projection, a positive-semidefinite covariance matrix. These objects control purity loss, branch distinguishability, decoherence-free information, many-body scaling, and the correctability of collapse errors.

The resulting statements are rigorous for the leading effective generator. Coherent Hamiltonian control cannot contribute an opposite sign to the purity budget; joint-eigenstate branches decay according to the CFS distinguishability functional; finite-range summable covariance yields at most extensive GHZ decoherence; pair-cat experiments reconstruct the projected covariance; and Knill–Laflamme theory gives the exact first-order criterion for active correction. A no-heating collapse mechanism can still have all of these coherence signatures.

None of these results presently rules out CFS. The missing ingredient is not another generic theorem about Lindblad noise, but a CFS calculation that fixes the processor-visible kernels and their covariance strongly enough to compare the lower requirement for macroscopic collapse with the upper requirement from coherent quantum computation. Equation (50) states that comparison explicitly. Scalable quantum processors therefore provide not a categorical objection to causal fermion systems, but a direct and increasingly stringent experimental arena for its collapse sector.

Acknowledgments

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Data availability

No data were generated or analyzed in this theoretical study.

A Short-time Kraus form

For completeness, consider one Hermitian collapse generator N with rate λ. Its dissipator can be written

$$-\lambda[N,[N,\rho]] = 2\lambda N\rho N - \lambda\{N^2,\rho\}. \tag{53}$$

Over a short interval δt, a Kraus representation through first order is

$$E_0 = \mathbb{1} - iH\delta t - \sum_{k,a} \lambda_{k,a} N_{k,a}^2 \delta t, \tag{54}$$

$$E_{k,a} = \sqrt{2\lambda_{k,a}\delta t}\, N_{k,a}, \tag{55}$$

up to terms $O(\delta t^{3/2})$ in the jump operators and $O(\delta t^2)$ in the channel. The completeness relation satisfies

$$E_0^\dagger E_0 + \sum_{k,a} E_{k,a}^\dagger E_{k,a} = \mathbb{1} + O(\delta t^2). \tag{56}$$

This makes the error span used in Proposition 10 explicit.

B Fixed-point algebra of the dissipator

For an arbitrary Hilbert–Schmidt operator X,

$$-\operatorname{Tr}\Big(X^\dagger \mathcal{D}_{\mathrm{CFS}}(X)\Big) = \sum_{k,a} \lambda_{k,a} \operatorname{Tr}\Big(X^\dagger [N_{k,a},[N_{k,a},X]]\Big) \tag{57}$$

$$= \sum_{k,a} \lambda_{k,a} \,\|[N_{k,a},X]\|_{\mathrm{HS}}^2 . \tag{58}$$

Thus the dissipator is negative semidefinite with respect to the Hilbert–Schmidt inner product. In finite dimensions, if $\mathcal{D}_{\mathrm{CFS}}(X) = 0$, then the left-hand side vanishes and therefore every active

commutator vanishes. This proves the fixed-point statement used in Theorem 3 for the dissipative part alone. A Hamiltonian term can of course move operators out of this fixed algebra unless the Hamiltonian preserves it.

C General branch rates for Pauli dephasing

For the covariance generator Eq. (28), computational-basis states satisfy

$$Z_i|x\rangle = s_i(x)|x\rangle, \qquad s_i(x) \in \{+1, -1\}. \tag{59}$$

Therefore

$$\langle x|[Z_i, [Z_j, \rho]]|y\rangle = \big(s_i(x) - s_i(y)\big)\big(s_j(x) - s_j(y)\big)\rho_{xy}. \tag{60}$$

Setting $d_i = s_i(x) - s_i(y)$ gives Eq. (30). The pair-cat reconstruction Eqs. (39)–(41) follows by choosing difference vectors $2(e_i + e_j)$ and $2(e_i - e_j)$.

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