

A Diósi–Penrose Decoherence Rate Breaks Lorentz Invariance

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Abstract

Diósi and Penrose proposed that a mass superposition loses coherence at the rate $\Gamma = CE_G/\hbar$, first order in Newton’s constant. We ask what an observation of this rate would imply about spacetime symmetry. A companion paper bounded, through the orbital decay of the double pulsar, the fluctuation–dissipation temperature of a reciprocal bath that supplies the rate at kilometre scales, $T_{\text{eff}}/C > 3.3 \times 10^{16}$ K. Here we show that the law of the noise cannot be Lorentz invariant if matter is local relativistic quantum field theory with a stable vacuum, the noise excites that vacuum at a finite rate per unit volume, the noise is a classical field or a Gaussian record fed back through translation-invariant kernels with an exactly Newtonian infrared limit, its law does not depend on the state, and no further degrees of freedom exist. A noise without feedback whose spectrum reaches timelike momenta excites the vacuum without bound at second order unless its Lorentz stabilizer is compact, so it has a rest frame; a Lorentz-invariant noise with purely spacelike support does so at fourth order, for time-ordered dynamics, where two spacelike transfers sum to a timelike one. For a record with feedback the kernels cannot be Lorentz covariant. Independently, a digitized-figure refit of XENONnT data limits the spacelike noise to $C < 2 \times 10^{-8}$. A law that reads its time direction from the state is a dynamical aether, nonlinear over mixtures; the branch-averaged cosmic energy flow is one possible carrier. The conclusion is not established for infrared-modified Newtonian responses, non-Gaussian or jump records, or non-local couplings, which bring their own ordering. A frame is observable as a Gaussian cutoff of the spontaneous X-ray spectrum near $\hbar c/R_0$, which no covariant noise produces; for a minimal frame theory, left to separate work, the same refit gives $1.9 \text{ \AA} < R_0 < 4 \times 10^{-5} \text{ m}$ at $C = 1$.

1 Introduction

Diósi [1, 2] and Penrose [3] proposed that a mass in a superposition of two locations loses coherence at the rate

$$\Gamma = \frac{CE_G}{\hbar}, \quad E_G = \frac{G}{2} \iint \frac{\Delta\rho(\mathbf{x})\Delta\rho(\mathbf{y})}{|\mathbf{x}-\mathbf{y}|} d^3x d^3y, \quad \Delta\rho = \rho_L - \rho_R, \quad (1)$$

with the mass density smeared over a length R_0 and C a number of order one ($C = 1$ in Diósi’s normalization). The rate is first order in G and needs no environment. It is within

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reach of a free-fall experiment: a sphere of radius 150 nm and density 5000 kg m^{-3} held in a superposition of positions $1 \text{ }\mu\text{m}$ apart for 100 s accumulates $\Gamma t = 2.2 C$, independently of R_0 for $R_0 \ll 150 \text{ nm}$ [4, 5, 6]. Suppose the fringes fade at this rate once every other channel is subtracted. What has been learned about the symmetry of spacetime?

A companion paper [6] answered part of the question. A bath that supplies the rate through a Hamiltonian coupling drags moving bodies, and the orbital decay of the double pulsar PSR J0737–3039A/B bounds its fluctuation–dissipation temperature at kilometre scales by $T_{\text{eff}}/C > 3.3 \times 10^{16} \text{ K}$; noise without back-action is the member of that class without a temperature parameter. This paper asks how the noise can act relativistically. Under stated hypotheses on matter and on the noise, among them a finite rate of vacuum excitation, the law by which the noise acts on matter cannot be Lorentz invariant; for a noise without feedback whose spectrum reaches timelike momenta its Lorentz stabilizer is compact, so it has a rest frame. A noise with timelike spectral support excites the vacuum without bound at second order; a noise with purely spacelike support does so at fourth order for time-ordered dynamics, and X- and γ -ray data exclude it independently at $C > 2 \times 10^{-8}$. A law that reads its time direction from the state instead of carrying it is a dynamical aether (Theorem 2.2).

The obstruction to covariant collapse dynamics is old. Ghirardi, Grassi and Pearle found that relativistic dynamical reduction excites the vacuum [7]; Pearle avoided the excitation at lowest order with a noise of spacelike, tachyonic support [8] and later with a quasi-relativistic model whose collapse operator annihilates the vacuum and whose time ordering is preferred [9]. Myrvold showed that a Lorentz-invariant Markovian collapse dynamics with a stable vacuum needs non-standard degrees of freedom [10]. Jones, Gasbarri and Bassi showed that a mass-coupled covariant collapse model cannot avoid both a divergent energy rate and superluminal signalling [11]. Diósi formulated continuous measurement of quantum fields relativistically [12] and for retarded observables [13], examined relativistic Gorini–Kossakowski–Lindblad–Sudarshan equations [14], and found that the local, Lorentz-invariant kernels relativity demands of a hybrid classical–quantum dynamics heat at an infinite rate [15]. On the other side, Bedingham built a covariant, finite collapse model with a pointer field [16], and Tumulka a frame-free relativistic flash model for non-interacting particles [17]. For gravity, Kafri, Taylor and Milburn modelled Newtonian interaction as a classical channel [18]; Tilloy and Diósi sourced a classical Newtonian field by a continuous record of the mass density and recovered the Diósi–Penrose rate [19], and their principle of least decoherence fixes $C = 1$ [20]. Oppenheim and collaborators built a classical–quantum theory of gravity [21], derived its decoherence–diffusion trade-off [22], obtained the Diósi–Penrose kernel in its Newtonian limit [23], and wrote covariant path integrals [24, 25]. Gisin showed that dynamics nonlinear over mixtures can signal [26].

What is new here is limited to three items. (1) Theorem 4.4 links a causal, exactly Newtonian classical response and complete positivity to an infinite vacuum excitation in every covariant record-and-feedback theory. It uses neither the Markov property nor a non-radiation hypothesis; the boost-orbit divergence it ends in is the standard one [7, 10, 11]. (2) The spacelike-support door is closed twice: by theory at fourth order for time-ordered dynamics, applying Pearle’s higher-order re-excitation [8] to the Diósi–Penrose normalization, and by data, because covariance ties the noise near the light cone to the noise that fixes the rate, so its spontaneous emission cannot be cut off (Section 5). (3) A refit of the XENONnT electronic-recoil spectrum [27, 28] with the cut emission spectrum of a frame theory (Appendix A). Everything else is assembled from the cited work.

Section 2 fixes the setting and states the theorem. Section 3 recalls arm (i). Sections 4–6 prove arm (ii): state-independent laws, the spacelike door, and the state-dependent alternative. Section 7 gives the observable signature of a frame, the X-ray cutoff.

2 Setting and statement

2.1 What a rate measures

A rigid body is prepared in two branches of equal mass at static positions; the source is isolated and inertial. A rate is a constant Γ with fringe visibility $|z(t)| \propto e^{-\Gamma t}$ at late times. Let X denote whatever decoheres the body: a classical stochastic potential, a quantum bath, or a classical field driven by a measurement record. X acts on matter through $H_{\text{int}} = \int \rho \Phi d^3x$, and

$$S(\mathbf{k}, \omega) = \frac{1}{2} \int dt e^{i\omega t} \langle \{\delta\Phi_{\mathbf{k}}(t), \delta\Phi_{-\mathbf{k}}(0)\} \rangle \quad (2)$$

is its symmetrized spectrum per unit volume. Signature is $(-+++)$, $q = (\omega, \mathbf{k})$, $q^2 = k^2 - \omega^2/c^2$; q is spacelike for $q^2 > 0$. The mass density is smeared by a normalized Gaussian of width R_0 , so that

$$E_G = \frac{G}{2} \int \frac{d^3k}{(2\pi)^3} \frac{4\pi}{k^2} e^{-k^2 R_0^2} |\Delta\rho_{\mathbf{k}}|^2. \quad (3)$$

Lemma 2.1 (A rate is a zero-frequency spectrum). *Let X be stationary with integrable correlations and correlation time τ_c . For $t \gg \tau_c$ the term of $-\ln|z(t)|$ quadratic in $\Delta\rho$ is*

$$\frac{t}{2\hbar^2} \int \frac{d^3k}{(2\pi)^3} |\Delta\rho_{\mathbf{k}}|^2 S(\mathbf{k}, 0), \quad (4)$$

and the remaining terms are of fourth and higher order in $\Delta\rho$. Hence $\Gamma = CE_G/\hbar$ at this order, for every $\Delta\rho$ supported in a band of wavenumbers, if and only if

$$S(\mathbf{k}, 0) = \frac{4\pi C\hbar G}{k^2} e^{-k^2 R_0^2} \quad \text{on that band.} \quad (5)$$

Proof. The relative phase of the branches is $\varphi(t) = \hbar^{-1} \int_0^t dt' \int \Delta\rho \Phi d^3x$ and $|z| = |\langle e^{i\varphi} \rangle|$. The cumulant expansion gives $-\ln|z| = \frac{1}{2}\kappa_2 - \frac{1}{24}\kappa_4 + \dots$, with $\kappa_n = \mathcal{O}(\Delta\rho^n)$. The second cumulant is

$$\kappa_2 = \frac{1}{\hbar^2} \int \frac{d^3k}{(2\pi)^3} |\Delta\rho_{\mathbf{k}}|^2 \int \frac{d\omega}{2\pi} S(\mathbf{k}, \omega) \frac{4\sin^2(\omega t/2)}{\omega^2}, \quad (6)$$

and $4\sin^2(\omega t/2)/\omega^2 \rightarrow 2\pi t \delta(\omega)$ on functions that vary on the scale $1/\tau_c \gg 1/t$. This gives Eq. (4). Inserting Eq. (5) gives $\Gamma = (C/\hbar) (G/2) \int (4\pi/k^2) e^{-k^2 R_0^2} |\Delta\rho_{\mathbf{k}}|^2 = CE_G/\hbar$ by Eq. (3); conversely, equality for every $\Delta\rho$ in the band fixes $S(\mathbf{k}, 0)$ there. $\square$

Gaussianity is not used. A compound-Poisson noise of phase kicks has the same quadratic coefficient and leaves the Diósi–Penrose class only for large kicks, where its rate saturates at the jump rate and stops scaling with E_G . The Markov property is not used either. A coloured noise gives a rate for $t \gg \tau_c$; for $t \ll \tau_c$ the exponent is quadratic in t and there is no rate. What is used is stationarity, integrable correlations, and $\Gamma\tau_c \ll 1$.

2.2 Hypotheses

- (R) *Rate*. For an isolated, inertial, equal-mass two-branch superposition, the late-time decoherence exponent at second order in $\Delta\rho$ is $CE_G/\hbar$ with $C > 0$, over at least one decade of laboratory wavenumbers, with $\Gamma\tau_c \ll 1$.
- (Q) *Standard matter*. Matter is a local relativistic quantum field theory with a unique Poincaré-invariant vacuum Ω obeying the spectrum condition. X couples to a local density $O(x)$, a component or contraction of $T^{\mu\nu}$, whose vacuum Källén–Lehmann weight $\sigma_O(s)$ is positive for almost every $s > 0$, and whose time-ordered two-point amplitude $\langle n|T\tilde{O}(q_1)\tilde{O}(q_2)|\Omega\rangle$ to two-particle states is nonzero on a set of spacelike (q_1, q_2) of positive measure.
- (S) *Stable vacuum under the noise*. At each order in G at which it is computed, the probability per unit volume and time that the noise excites the matter vacuum is finite.
- (N) *No further degrees of freedom*. Either (a) X is a classical stochastic field ϕ , or a bath in a tracial state, acting through $\int O\phi d^4x$ with time-ordered evolution; or (b) X is a classical field h that responds linearly to a continuous record of O and acts back through $\int hO d^4x$. Nothing else carries energy, momentum or excitations.
- (L) *State-independence*. The law of X , meaning its spectrum, its kernels and its record, does not depend on the quantum state of matter.
- (G) *Gaussian record, Newtonian response* (case (b)). The record noise and the induced influence functional are Gaussian; the decoherence kernel $\Lambda(q)$ and the response kernel $\tilde{D}(q)$ are translation-invariant; $\tilde{D}$ is retarded and *exactly Newtonian in the infrared*, $q^2\tilde{D}(q) \rightarrow 2\pi G/\hbar$ as $q^2 \rightarrow 0$ in the normalization of Section 4.3.

For the traceless components of $T^{\mu\nu}$ the weight σ_O near $s \rightarrow 0$ comes from photon (and graviton) pairs. For $T^\mu{}_\mu$ it comes, above the hadronic and e^+e^- thresholds, from the trace anomaly; below $2m_e c^2$ there are no light charged particles and gluons are confined, and the two-photon weight comes from the Euler–Heisenberg operators, suppressed by powers of $s/m_e^2 c^2$ but positive. The amplitude condition in (Q) holds at tree level whenever O is bilinear in a massive field, as for $:\psi^2:$ or $m\bar{\psi}\psi$ (Section 5).

2.3 The theorem

Theorem 2.2 (Chain). *Assume (R).*

- (i) *Back-action at kilometre scales. If X acts through a Hamiltonian and obeys the conditions of the drag theorem of Ref. [6] at the wavenumbers a neutron star resolves, then $|T_{\text{eff}}|/C > 3.3 \times 10^{16}$ K. The member of the class free of a temperature parameter is $T_{\text{eff}} = \infty$: noise without back-action.*
- (ii) *Lorentz invariance is broken. Assume (Q), (S), (N) and (L). In case (a), with time-ordered evolution, the law of X is not Lorentz invariant. If its spectrum is nonzero on a set of timelike q of positive measure, its Lorentz stabilizer is compact; it then fixes a unit timelike vector u , the rest frame of the noise, unique if the law is isotropic about it. In case (b) with (G), the kernels by which X acts on matter are not Lorentz covariant.*

(iii) An independent experimental closure. *Within the semiclassical model of spontaneous emission used by the X-ray searches [29, 28], a Poincaré-invariant noise with spacelike spectral support is consistent with the XENONnT spectrum only if $C < 1.8 \times 10^{-8}$.*

Arms (i) and (ii) are independent: neither uses the other, and neither uses the Newtonian mean field except through (G). Arm (i) is Section 3. In (ii), case (a) is Lemma 4.1 for a noise with timelike spectral support and Lemma 5.2 for a Poincaré-invariant one with purely spacelike support; whether a spacelike-only noise invariant under a smaller non-compact group can satisfy (R) is open (Section 5); case (b) is Theorem 4.4. (iii) is Proposition 5.4; it does not use time ordering, so it also covers dynamics that drop it. Section 6 treats laws that violate (L). Case (b) is weaker than case (a): Theorem 4.4 excludes Lorentz-covariant kernels, not kernels invariant under a non-compact subgroup such as $SO(2, 1)$ (Remark 4.2).

Remark 2.1 (Where the chain is not established). *An infrared-modified Newtonian law:* a response whose Källén–Lehmann weight is spread over $(0, \mu^2)$ gives $k^2 \tilde{D}(0, k) \rightarrow 0$ as $k \rightarrow 0$, Newtonian only at $r \ll 1/\mu$, and the step of Theorem 4.4 that uses the exact infrared limit fails. *Non-Gaussian or jump records:* the trade-off and the excitation rate of Section 4.3 need a Gaussian record, and a Poissonian hybrid with a stable vacuum exists [30]; flash-sourced Newtonian gravity is of this kind [31]. For Gaussian kick events at rate ν the floor on the rate is $\nu(1 - e^{-E_G/\hbar\nu})$, which falls below $E_G/\hbar$ only where the rate saturates at ν and stops scaling with E_G ; whether such a record reproduces Eq. (1) with $C \sim 1$ over the laboratory band and stays covariant is open. *Non-local couplings:* an O that does not commute with itself at spacelike separation violates (Q) and brings its own ordering surface (Section 4.4).

3 Arm (i): back-action at kilometre scales

Ref. [6] proves arm (i); we restate it and add one reading. If the noise of Eq. (5) is carried by a bath operator coupled through a Hamiltonian, the same correlator carries the dissipative response $\chi''(\mathbf{k}, \omega)$, and a body moving through the bath is dragged. Under stationarity, whiteness over the drift band, passivity and linear response at wavenumbers $k \leq k_* = 10/R_{\text{NS}}$, the friction rate is $\gamma \geq (2\pi/3)(C\hbar G/k_B T_{\text{eff}})\rho_{k_*}$, and the orbital decay of PSR J0737–3039A/B gives

$$|T_{\text{eff}}|/C > T_* = 3.3 \times 10^{16} \text{ K} \quad (7)$$

(Theorem 3.1 and Corollary 3.3 of Ref. [6]). The escapes are a kernel cut below kilometre wavenumbers, a bath slower than 1.4 ms at those scales, a temperature that rises toward long wavelengths, and active modes; each changes the bath between the laboratory and the neutron star.

Corollary 3.1 (Response lag). *Under the conditions of Eq. (7), measured against Newton’s static response $4\pi G/k^2$, the response of the bath to the mass density lags by*

$$\tau_{\text{lag}} < \frac{\hbar}{2k_B T_*} = 1.2 \times 10^{-28} \text{ s}, \quad (8)$$

for every C .

Proof. At $\hbar\omega \ll k_B T_{\text{eff}}$ the fluctuation–dissipation relation reads $S(\mathbf{k}, 0) = 2k_B T_{\text{eff}} \partial_\omega \chi''(\mathbf{k}, 0)$.

With Eq. (5) at $k \ll 1/R_0$, $k^2 \partial_\omega \chi''(\mathbf{k}, 0) = 2\pi C \hbar G / (k_B T_{\text{eff}}) < 2\pi \hbar G / (k_B T_*)$, and C cancels. Writing $\chi(\mathbf{k}, \omega) = (4\pi G/k^2)(1 + i\omega \tau_{\text{lag}}) + \mathcal{O}(\omega^2)$ gives $\tau_{\text{lag}} = k^2 \partial_\omega \chi'' / (4\pi G)$, hence Eq. (8). $\square$

The bound is a statement about the dissipative slope alone. The member of the class without a temperature parameter is $T_{\text{eff}} = \infty$, $\chi'' \equiv 0$: a classical noise, a bath in a tracial state, or a classical field sourced instantaneously by a record. These are the cases (a) and (b) of hypothesis (N).

4 Arm (ii): Lorentz invariance is broken

4.1 Noise without feedback

Lemma 4.1 (Covariant noise excites the vacuum). *Assume (Q) and case (a) of (N). At second order in ϕ , the probability per unit volume and time that the vacuum is excited is*

$$\mathcal{R} = \frac{1}{\hbar^2} \int \frac{d^4 q}{(2\pi)^4} W(q) \rho_O(q), \quad \rho_O(q) = \theta(q^0) \sigma_O(-q^2), \quad (9)$$

where W is the spectrum of ϕ (symmetrized, for a tracial bath) and σ_O the Källén–Lehmann weight of O . If W is nonzero on a set of timelike q of positive measure and the law of ϕ is invariant under a closed non-compact subgroup H of the Lorentz group (a one-parameter group of boosts or of null rotations, or the whole group), then $\mathcal{R} = \infty$.

Proof. To first order in the coupling the amplitude to reach $|n\rangle$ is $-(i/\hbar) \int d^4 x \phi(x) \langle n | O(x) | \Omega \rangle$. Averaging its square over ϕ and summing over n gives $\hbar^{-2} \iint d^4 x d^4 y \langle \phi(x) \phi(y) \rangle \langle \Omega | O(y) O(x) | \Omega \rangle$; translation invariance turns this into Eq. (9) per unit four-volume, and the spectrum condition restricts ρ_O to the closed forward cone. A Poincaré-invariant law has $W(q) = w(q^2)$. Writing $s = -q^2$ on the forward cone,

$$\mathcal{R} = \frac{1}{\hbar^2 (2\pi)^4} \int ds w(-s) \sigma_O(s) \int \frac{d^3 q}{2\sqrt{q^2 + s}}. \quad (10)$$

The inner integral is the volume of the mass hyperboloid; below $|\mathbf{q}| < P$ it equals $\pi[P\sqrt{P^2 + s} - s \operatorname{arsinh}(P/\sqrt{s})]$ and grows as πP^2 . By the support assumption and (Q), $w\sigma_O > 0$ on a set of s of positive measure, so $\mathcal{R} = \infty$. For a general closed non-compact H , the hyperboloid is $SO^+(1, 3)/SO(3)$ and its point stabilizers are compact, so H acts properly: for compact A the set $\{h \in H : hA \cap A \neq \emptyset\}$ is compact, and there are $h_1, h_2, \dots \in H$ with the $h_i A$ pairwise disjoint. The set B where $W\rho_O > 0$ is H -invariant; choosing A with $\mu(B \cap A) > 0$, B contains the disjoint sets $h_i(B \cap A)$ of equal invariant measure, so the shell integral is infinite on a set of shells of positive measure. $\square$

Corollary 4.2 (Compact stabilizer). *Under (Q), (S), (N)(a) and (L), the Lorentz stabilizer of a noise law with timelike spectral support is compact, and every compact subgroup K of $SO^+(1, 3)$ fixes a unit timelike vector u .*

Proof. The stabilizer is closed; by Lemma 4.1 and (S) it is not non-compact. For a compact K , average any future timelike unit vector over the Haar measure of K : the average lies in the open future cone, which is convex, and is K -invariant; normalize it. $\square$

The vector u is unique when the law is isotropic about it; an anisotropic law with a smaller compact stabilizer still distinguishes frames but may fix several timelike directions. The mechanism of Lemma 4.1, a positive rate density integrated over a non-compact orbit, is the standard obstruction to relativistic collapse [7, 8, 10, 11]; no novelty is claimed for it. No resummation makes a positive, invariant excitation density of nonzero-momentum states finite. The divergence is the price of the symmetry, not of the floor (5): a noise with that zero-frequency spectrum and a frequency profile that falls faster than any power in a frame, for instance $e^{-\omega^2/\Omega^2}$, excites the vacuum at a finite rate that does not depend on a momentum cutoff, even for components of $T^{\mu\nu}$, whose weight grows as $\sigma_O \propto s^2$. A Lorentzian profile does not suffice there: against $\sigma_O \propto s^2$ its frequency integral diverges. Spectral support on timelike q is used; a noise without it is Section 5.

4.2 Covariant positivity

Lemma 4.3 (Positive covariant kernels see the trace). *Let $\Lambda(q)$ be a positive semidefinite Hermitian form on symmetric tensors that is Lorentz covariant, $\Lambda(Lq)(Lt, Lt) = \Lambda(q)(t, t)$. At spacelike q , $\Lambda(q)(t, t)$ is a positive form in $(t^\mu_\mu, q_\mu q_\nu t^{\mu\nu})$; on conserved tensors, $q_\mu t^{\mu\nu} = 0$, it equals $a(q^2) |t^\mu_\mu|^2$ with $a \geq 0$.*

Proof. The stabilizer of a spacelike q is $SO^+(2, 1)$, connected, simple and non-compact. The null space of $\Lambda(q)$ is invariant under it, and on the quotient $\Lambda(q)$ is an invariant inner product, so the stabilizer acts unitarily on a finite-dimensional space. A finite-dimensional unitary representation of a connected simple non-compact group is trivial. Hence $\Lambda(q)$ factors through the invariant linear functionals on symmetric tensors. Under the stabilizer, symmetric tensors decompose as $\mathbf{1} \oplus \mathbf{1} \oplus \mathbf{3} \oplus \mathbf{5}$, and the two singlets are contraction with $\eta_{\mu\nu}$ and with $q_\mu q_\nu$. On conserved tensors the second vanishes. $\square$

Static Newtonian sources sit at spacelike q . A covariant, completely positive classical sector can therefore respond to them only through the trace: it is Nordström-like, gives no light deflection and no gravitomagnetism. For an ultralocal kernel this is the statement of Oppenheim, Russo and Weller-Davies that a positive semidefinite kernel cannot capture the transverse part of the Einstein equations in Lorentzian signature, and their CQ-Nordström theory is the trace case [25]; the lemma extends it to non-local kernels by the little-group argument. For pure noise without feedback, Lemma 4.3 fixes $O = -T^\mu_\mu$, which for laboratory matter is ρc^2 up to $p/\rho c^2$.

4.3 Record and feedback

In case (b) a classical field responds linearly to a record of O , $h = D*(O + \xi)$ with record noise ξ of spectral density $N(q)$, and acts back through $\int h O$. Integrating out h gives a Gaussian influence functional on matter with decoherence kernel $\Lambda(q)$ and response $\tilde{D}(q)$. We normalize both so that the decoherence exponent of a static superposition is $t \int d^3k (2\pi)^{-3} \Lambda(0, \mathbf{k}) |\Delta \rho_{\mathbf{k}}|^2$ and the Newtonian response is $\tilde{D}(0, \mathbf{k}) = 2\pi G/(\hbar k^2)$; the Diósi–Penrose rate is then $\Lambda(0, \mathbf{k}) = C \tilde{D}(0, \mathbf{k})$. The record costs decoherence $1/(4N)$, and the field noise $|\tilde{D}|^2 N$ acts back on O , so at every q

$$\Lambda(q) \geq \frac{1}{4N(q)} + |\tilde{D}(q)|^2 N(q) \geq |\tilde{D}(q)|, \quad (11)$$

with equality at $N = 1/(2|\tilde{D}|)$. This is the measurement imprecision–back-action inequality resolved in frequency; at $\omega = 0$ it is the floor $C \geq 1$ of Tilloy and Diósi [19, 20] and the decoherence–diffusion trade-off of Oppenheim et al. [22, 23].

Theorem 4.4 (Covariant feedback excites the vacuum). *Assume (Q), case (b) of (N), (L) and (G), and let $\Lambda, \tilde{D}$ be Lorentz covariant. Then the vacuum is excited at an infinite rate per unit four-volume, for every value of G .*

Proof. Step 1: the rate keeps the dissipative term. At first order in G the vacuum transition probability per unit four-volume of the Gaussian hybrid is

$$\frac{dP}{d^4x} = \int \frac{d^4q}{(2\pi)^4} \theta(q^0) [\Lambda(q) \mp \text{Im } \tilde{D}(q)] \rho_O(q), \quad (12)$$

the noise kernel corrected by the absorptive response, as $S(\omega) \mp \hbar\chi''(\omega)$ is for the up-transition rate of a quantum bath; the sign depends on the convention for $\tilde{D}$. In either convention the integrand is at least $\Lambda - |\text{Im } \tilde{D}| \geq |\tilde{D}| - |\text{Im } \tilde{D}| \geq 0$ by Eq. (11), and it is positive wherever $\text{Re } \tilde{D} \neq 0$. Positivity can fail only where $\text{Re } \tilde{D} = 0$, and that set is not empty: $\tilde{D} = c/\sqrt{q^2}$ is retarded, causal and covariant, purely imaginary on the timelike cut, and with a saturated trade-off excites nothing.

Step 2: exact infrared Newton forces $\text{Re } \tilde{D} \neq 0$ near the cone. By covariance and retardation $\tilde{D} = d(q^2)$ is the boundary value of a function analytic in the plane cut along the timelike axis. By (G), $d(q^2) = 2\pi G/(\hbar q^2) + o(|q^2|^{-1})$ as $q^2 \rightarrow 0$, with a real residue. On the timelike side, $q^2 = -s$, this gives $\text{Re } d = -2\pi G/(\hbar s) + o(s^{-1}) \neq 0$ for $0 < s < s_0$, and there the integrand is positive. The exactness matters: a response $\propto (q^2)^{-1/2}$, as in Step 1, is real on the spacelike side and purely imaginary on the timelike side.

Step 3: matter has spectral weight there. By (Q), $\rho_O > 0$ for almost every $s \in (0, s_0)$: photon pairs for the traceless components of $T^{\mu\nu}$, and for $T^\mu{}_\mu$ two-photon states through the Euler–Heisenberg operators, power-suppressed but positive. No local O annihilates the vacuum [32].

Step 4: the boost orbit. The integrand of Eq. (12) is a positive function of q^2 on a set of positive measure, and the integral over the mass hyperboloid at fixed q^2 diverges as in Eq. (10). $\square$

Remark 4.1 (What the theorem uses and what it adds). Theorem 4.4 uses neither the Markov property nor a hypothesis that the classical field does not radiate: a finite record correlation time is a function of q^2 and stays nonzero on timelike q , and radiation enters only through $\text{Im } \tilde{D}$, which Step 1 keeps. The divergence in Step 4 is the standard one. What the theorem adds is the link of Steps 1–2: a causal response that is exactly Newtonian in the infrared, together with complete positivity, cannot keep the excitation density zero on the timelike side of the cone. Non-Markovian records of retarded observables [13] are covered. For a channel with a gap, massive pairs only above $4m^2$, Step 2 would be needed at $s > 4m^2$, which the infrared limit alone does not guarantee; hypothesis (Q) excludes such a gap, and for $T^\mu{}_\mu$ the two-photon weight removes it.

Remark 4.2 (Kernels invariant under a subgroup). If Λ and $\tilde{D}$ are invariant only under a closed non-compact subgroup H , Steps 1, 3 and 4 go through, Step 4 by the argument of Lemma 4.1. Step 2 used covariance to write $\tilde{D}$ as a function of q^2 . It holds without covariance if the

Newtonian limit $q^2 \tilde{D} \rightarrow 2\pi G/\hbar$ holds on approach to the light cone from the timelike side; the stabilizer is then compact as in case (a). A response invariant under, say, $SO(2,1)$ and Newtonian only at $\omega = 0$ is not excluded by Theorem 4.4.

Remark 4.3 (Classical–quantum gravity). The covariant classical–quantum path integrals of Oppenheim and Weller-Davies are manifestly covariant [24]. Their decoherence term is ultralocal, with a free coupling, and decoheres into eigenstates of the trace of $T^{\mu\nu}$: the case Lemma 4.3 allows, and not the Diósi–Penrose kernel. The Diósi–Penrose kernel appears in that programme in the Newtonian limit, obtained by gauge fixing on spatial slices with an instantaneous kernel [23]: on a foliation. A covariant trace channel with a Newtonian kernel $\propto 1/|q^2|$ falls under Theorem 4.4 through $T^\mu{}_\mu$.

4.4 Non-local couplings and white noise

Two structural facts locate the time direction.

Non-local couplings need an ordering. If O does not commute with itself at spacelike separation, as for the spacelike projection T^{SL} of Section 7 or for a number density $a^\dagger a$, the second-order averaged generator contains $\theta(t_x - t_y) S(x - y) [O(x), O(y)]$ on spacelike pairs. That term depends on the surface used to order the pair, so the Tomonaga–Schwinger evolution is not integrable and a foliation enters the dynamics. We give this as an argument, not a theorem. Pearle’s quasi-relativistic model, whose collapse operator annihilates the vacuum, has a preferred ordering [9], and his later relativistic model evolves in a parameter that labels a foliation of spacelike hypersurfaces [33]. Non-local couplings violate the locality in (Q), so they evade Lemma 4.1 and Theorem 4.4 only by bringing their own foliation.

White noise defines leaves. A noise white in a time τ with a non-local spatial kernel, $\langle \phi(x)\phi(y) \rangle = \delta(\tau(x) - \tau(y))K(x, y)$, has its correlations supported on the relation $\tau(x) = \tau(y)$, whose classes are the level sets of τ . A white Diósi–Penrose noise, whose kernel is $1/r$ by Lemma 2.1, therefore defines a simultaneity foliation. In Minkowski space a constant u and the u -orthogonal hyperplanes are equivalent; in curved spacetime a coloured noise could single out a congruence that is not hypersurface-orthogonal, and that case is open.

5 The spacelike door

Lemma 4.1 leaves one option for noise without feedback whose law keeps a non-compact symmetry: a classical Gaussian field whose spectrum vanishes on timelike q . We treat the Poincaré-invariant case. By Lemma 4.3 it couples through $O = -T^\mu{}_\mu = \rho c^2(1 - 3p/\rho c^2)$, the mass density of laboratory matter up to the negligible $3p/\rho c^2$. In the normalization of Eq. (2),

$$S(q) = 4\pi C \hbar G s(q^2) \theta(q^2), \quad \text{e.g. } s(u) = \frac{e^{-uR_0^2}}{u + \mu^2}. \quad (13)$$

This is Pearle’s tachyonic noise [8] with the Diósi–Penrose normalization put in by hand.

Proposition 5.1 (What the spacelike noise passes). *The noise (13) (a) satisfies Eq. (5) for $\mu \ll k \ll 1/R_0$, exactly as $\mu \rightarrow 0$; (b) leaves the vacuum unexcited at second order, since its support is disjoint from the forward cone; (c) has $\chi'' \equiv 0$, so it does not drag; (d) does not signal: it acts by local unitaries with a c -number field, and $[O(x), O(y)] = 0$ at spacelike separation, so*

Bob's marginal does not depend on Alice's instrument; (e) heats matter as Diósi–Penrose noise does, to 10^{-10} , $P/M = CG\hbar/(4\sqrt{\pi}R_0^3)$, which is $1.1 \times 10^{-17} \text{ Wkg}^{-1}$ at $C = 1$, $R_0 = 4.5 \text{ Å}$; (f) is deterministic: at fixed $\mathbf{k}$ its spectrum vanishes for $|\omega| > c|\mathbf{k}|$, so $\int \ln S d\omega/(1+\omega^2) = -\infty$ and every mode is predictable from its past.

Proof. (a)–(d) are read off Eq. (13). For (e), a free nucleon absorbs $(\omega, \mathbf{k})$ only with $\omega = \hbar k^2/2m + \mathbf{k} \cdot \mathbf{v} \ll c|\mathbf{k}|$, so $S(\omega, \mathbf{k})/S(0, \mathbf{k}) - 1 = \mathcal{O}(\omega^2/c^2 k^2)$, below 10^{-10} at $k = 1/R_0$; the heating is then that of the Diósi–Penrose noise [6]. (f) is the Szegő–Kolmogorov criterion. $\square$

At second order heating, drag, signalling and the vacuum do not distinguish this noise from a frame noise. At fourth order the vacuum does.

Lemma 5.2 (Two spacelike transfers excite the vacuum). *Assume (Q), (N)(a) and a Poincaré-invariant law with spectrum on spacelike q . At fourth order in ϕ the probability per unit volume and time of exciting the vacuum is*

$$\mathcal{R}^{(4)} = \int d^4 P \theta(P^0) f(P), \quad f(P) \propto \sum_{n: p_n=P} \int d^4 q S(q) S(P-q) |\langle n|T \tilde{O}(q) \tilde{O}(P-q)|\Omega \rangle|^2, \quad (14)$$

with f Lorentz invariant and positive on a set of timelike P of positive measure, so $\mathcal{R}^{(4)} = \infty$.

Proof. With time-ordered evolution the second-order amplitude to $|n\rangle \neq |\Omega\rangle$ is

$$A_2(n) = -\frac{1}{2\hbar^2} \iint d^4 x d^4 y \phi(x) \phi(y) \langle n|T O(x) O(y)|\Omega \rangle. \quad (15)$$

At fourth order in ϕ the excitation probability is the Gaussian average of its square; the cross term of the first- and third-order amplitudes needs S at the timelike p_n and vanishes, and pairings inside one amplitude transfer zero four-momentum and reach only Ω . The cross pairings give Eq. (14). Two spacelike vectors can sum to a timelike P above the pair threshold. For $O = :\psi^2:$ of a free field of mass m , or $m\bar{\psi}\psi$ at tree level, the amplitude to two particles $p_1 + p_2 = P$ carries the propagator of the internal line $l = p_1 - q$. That line is never on shell: with $l^0 > 0$, $P - q = p_2 + l$ would be timelike, and with $l^0 < 0$, $q = p_1 + (-l)$ would be. So the amplitude is finite, and it vanishes only where the two contractions cancel, a set of measure zero; this is the amplitude condition of (Q). Covariance of S and of the time-ordered product makes f a function of P^2 , and the integral over the hyperboloid diverges as in Eq. (10). $\square$

A 1 + 1-dimensional evaluation of Eq. (14) for $:\psi^2:$ gives a finite positive f , equal at rapidities 0, 1 and 2 to 10^{-3} . Pearle found that his tachyonic model re-excites the vacuum beyond lowest order and avoided it by giving up time ordering, which permits spacelike transport [8]; Lemma 5.2 is that mechanism applied to the Diósi–Penrose door. The door is therefore closed by theory, at order C^2 , for time-ordered dynamics. With (S), Lemma 5.2 excludes a Poincaré-invariant spacelike-only law. It does not extend to every smaller symmetry: a noise supported on $\{q^0 = cq_z\}$ is invariant under boosts along z , and its momenta never sum to a timelike vector. That noise fails (R), whose floor is required in every direction; whether (R) restores the conclusion for such subgroups is open. The emission bound below is an independent closure that does not use time ordering.

5.1 An independent closure by X- and γ -ray data

Proposition 5.3 (Covariance forbids a cutoff on spontaneous emission). *A quasi-free charge kicked by the noise (13) radiates photons of energy E at a rate proportional to $I(E)/E$, with $I(E) = \int d^3k k^2 S(E/\hbar, \mathbf{k})$. For the white Diósi–Penrose noise at the same C , $I_w = \pi^{3/2}/R_0^3$ (in units of $4\pi C\hbar G$). For the covariant noise,*

$$I_c(E) = 2\pi \int_0^\infty du \left(u + \frac{E^2}{\hbar^2 c^2}\right)^{3/2} s(u) \geq 2\pi \left(\frac{E}{\hbar c}\right)^3 \int_0^\infty s(u) du, \quad (16)$$

independent of R_0 apart from the logarithm in $\int s$, and rising as E^2 in the spectrum I_c/E . If the rate (1) holds over $k \in [k_{\min}, k_{\max}]$, then $s(u) \simeq 1/u$ there and $\int s du \geq 2 \ln(k_{\max}/k_{\min})$.

Proof. Emission at energy E draws on noise at $\omega = E/\hbar$, and spacelike support requires $k > E/\hbar c$. With $u = k^2 - E^2/\hbar^2 c^2$, $d^3k = 2\pi k du$ after the angular integral and $k dk = du/2$, which gives the equality in Eq. (16); dropping u in the bracket gives the bound. The white value is $4\pi \int_0^\infty k^2 e^{-k^2 R_0^2} dk = \pi^{3/2}/R_0^3$. On the band where Eq. (5) holds, $s(k^2) = 1/k^2$, and $\int_{k_{\min}^2}^{k_{\max}^2} du/u = 2 \ln(k_{\max}/k_{\min})$. $\square$

The u that enter Eq. (16) are the same u that fix the rate at $\omega = 0$. No covariant form factor can cut the emission: on the light cone a function of q^2 equals its value at $q^2 \rightarrow 0$, so the near-cone noise at energy E is not suppressed for any E . A frame theory escapes because its smearing $e^{-k^2 R_0^2}$ is in $\mathbf{k}$, not in q^2 (Section 7).

Each search fixes a white edge R_X at $C = 1$, the radius at which the white emission just saturates its limit: 4.5 Å for XENONnT [28], 0.867 Å for MAJORANA [34, 35] and 0.184 Å for Gran Sasso [36], all in Diósi units [6]. By Eq. (16), at each search’s lowest energy the covariant emission exceeds the white emission at R_X by at least $(2/\sqrt{\pi})(ER_X/\hbar c)^3 \int s$:

Table 1. *Lower bound on the covariant emission relative to the white emission at each search’s edge, per unit $\int s du$, at the lowest energy of the search.*

Search	$E_{\min}$	$E_{\min} R_X / \hbar c$	$(2/\sqrt{\pi})(E_{\min} R_X / \hbar c)^3 \int s$
XENONnT	1 keV	2.28	13.4
MAJORANA	19 keV	8.35	6.6×10^2
Gran Sasso	1 MeV	93.3	9.2×10^5

With $\int s \geq 4$, less than the $2 \ln 10$ of a single decade, the Gran Sasso row alone requires $C < 3 \times 10^{-7}$. The E^2 spectrum is largest at the top of each band, so a fit to the full spectrum is stronger.

Proposition 5.4 (XENONnT bound on the spacelike door). *Fitting the covariant spectrum $C(2/\sqrt{\pi})(ER_0/\hbar c)^3 \int s$ times the white emission prefactor of xenon, with $\int s = 4$ and atomic screening evaluated at $k = E/\hbar c$, to the XENONnT SR0 electronic-recoil data with every published background constraint profiled (Appendix A) gives*

$$C < 1.8 \times 10^{-8} \quad (95\% \text{ CL}). \quad (17)$$

The bound (16) keeps the near-cone part of the noise, $k = E/\hbar c$. There the electron–nucleus and electron–electron terms of the emission rate carry $j_0(Er/\hbar c)$ in place of the Gaussian overlap $e^{-r^2/4R_0^2}$. Screening reduces the rate by 11% at 50 keV and by about 1% at 100–140 keV, where

the E^2 spectrum drives the fit; it is large only below 10 keV, where the spectrum is small. The fit with and without screening gives the same bound. This closure is an experimental exclusion within the semiclassical emission model [29] used by the searches.

6 Laws that read the time from the state

Hypothesis (L) was used tacitly by Lemma 4.1, as it is by Myrvold's theorem [10]. Bedingham's covariant model [16] shows what dropping it buys. A pointer field $a(x)$ with one quantized degree of freedom per spacetime point, $[a(x), a^\dagger(x')] = \delta^4(x - x')$, is written by matter through $A(x) = \int g(x, y) (a + a^\dagger)(y) d^4y$ with g supported in the future cone of x .

Lemma 6.1 (Pointer reduction). *Let matter couple through $\int J(x)A(x) d^4x$ with $[A(x), A(x')] = 0$ for all x, x' , and let the pointer start in its Fock vacuum. Then the reduced matter dynamics equals that generated by a classical Gaussian field with covariance $K(x, x') = \int g(x, y)g(x', y) d^4y$, and $\chi \equiv 0$.*

Proof. The $A(x)$ form a commuting family and are jointly diagonalizable. In the vacuum their joint law is Gaussian with the real covariance $\langle 0|A(x)A(x')|0\rangle = K(x, x')$. Conditioned on each joint eigenvalue configuration, the time-ordered evolution is the matter evolution under that c -number field. $\square$

A pointer channel is thus inside case (a): it has no back-action and realizes arm (i). Its one-particle spectrum is all of $\mathbb{R}^4$, which supplies off-shell, zero-frequency weight in a Poincaré-invariant vacuum; it does not shield the matter vacuum. A Lorentz-invariant g supported in the future cone depends only on the proper time τ of $y - x$; its integral over $\{1 < \tau < 2, |\mathbf{y}| < L\}$ grows as L^2 , so it cannot be normalized. A state-independent invariant pointer kernel is therefore a distribution. For the retarded Riesz kernels, with Fourier transform $(q^2)_{\text{ret}}^{-\alpha/2}$, $|\tilde{g}|^2 = |q^2|^{-\alpha}$ is smooth and positive on the timelike shell above the pair threshold, and Lemma 4.1 applies: at $\alpha = 1$, the Diósi–Penrose case, the vacuum excitation is infinite. For a general invariant tempered distribution supported in the future cone we conjecture the same; the regularity of $|\tilde{g}|^2$ on the timelike shell is the open step.

Bedingham's model avoids this by reading the state: his smearing is $g \propto \exp\{-k \bar{T}^{\mu\nu}(x)(y - x)_\mu(y - x)_\nu\}$, normalized on the future cone, with $\bar{T}^{\mu\nu} = \langle T^{\mu\nu} \rangle$ on the past light cone of x [16]. That this state dependence is how the model evades Myrvold's theorem was stated by Muciño, Okon, Sudarsky and Wiedemann [37]. In our reading, normalizing on the non-compact cone forces $\tilde{g}(q \neq 0) \rightarrow 0$ as $\bar{T} \rightarrow 0$, so the coupling switches off exactly at $\bar{T} = 0$ and is on for every $\bar{T} > 0$.

Proposition 6.2 (A state-borne time direction is a dynamical aether). *Let the noise law be covariant except through a dependence on the unconditional expectation $\bar{T}^{\mu\nu}$ on the past light cone, as in Bedingham's model. Then the time direction of the law is $u(x)$, the timelike eigenvector of $\bar{T}^{\mu\nu}(x)$, which exists for ordinary matter: the Landau frame of the branch-averaged energy flow. Since $\bar{T} \neq 0$ everywhere in our universe, u is defined everywhere and the vacuum gate at $\bar{T} = 0$ never closes. The dynamics is nonlinear over mixtures.*

The cosmic microwave background alone gives $\bar{T}^{00} = 4.2 \times 10^{-14} \text{ J m}^{-3}$ in voids, before neutrinos and dark matter, so in voids u is the rest frame of the cosmic medium and the vacuum is excited

at a finite, frame-cut, nonzero rate. Physically this is an aether carried by the matter content: a time direction selected spontaneously by the state, as the microwave background selects its rest frame, rather than written into the law. Matter’s energy flow is one possible carrier of u ; the argument does not show that it is the only one.

The nonlinearity is concrete. Send a body by a quantum coin either to rest or to speed v . Its branch-averaged energy flow, and with it u , moves at $v/2$, so each branch moves at $v/2$ relative to u . If the decoherence of a body moving at w through the frame scales as γ_w^3 relative to its frame-free value, as it does for an energy-density kernel defined in the frame (rate γ_w^2 times the rest rate, against $1/\gamma_w$ for a frame-free rate), each branch’s rate shifts by $\gamma_{v/2}^3 - 1 \simeq 3v^2/8c^2$, while either branch prepared alone shifts by zero. At $v = 1 \text{ ms}^{-1}$ the shift is 4×10^{-18} . A dynamics that is nonlinear over mixtures can signal [26], though nonlinear dynamics need not when the nonlinearity acts only through locally accessible data [38]; here u reads $\langle T \rangle$ inside the past light cone, which blocks the obvious steering protocols, but we have not shown that it does not signal. A matter density read from the state is also how relativistic collapse models define their local beables [39].

6.1 The chain, assembled

Proof of Theorem 2.2. (i) is Eq. (7), from Ref. [6], with Corollary 3.1 as its slope form. (ii) In case (a), a law with timelike spectral support that is invariant under a closed non-compact subgroup of the Lorentz group excites the vacuum at an infinite rate by Lemma 4.1, contradicting (S); its stabilizer is closed, hence compact, and fixes a unit timelike u by Corollary 4.2. A Poincaré-invariant law with purely spacelike support does the same at fourth order by Lemma 5.2. In case (b) with (G), Theorem 4.4 contradicts (S) for covariant $(\Lambda, \tilde{D})$. Pointer channels with state-independent kernels reduce to case (a) by Lemma 6.1. (iii) is Proposition 5.4. $\square$

Read as a contrapositive: if an isolated, inertial superposition decoheres at $\Gamma = CE_G/\hbar$ with C of order one, then, outside the exceptions of Remark 2.1, the law of the noise is not Lorentz invariant: the distinction between states of motion is either written into the law, which has a compact stabilizer when its spectrum reaches timelike momenta, or read by the law from the state, as in Proposition 6.2.

7 The fingerprint of a frame

Theorem 2.2 says what an observed rate requires. A frame theory that meets the requirement has been outlined, but it is not complete: it signals at a level suppressed by G , and its definition on curved backgrounds is open. In its minimal form the classical sector of geometry is the instantaneous constraint field on constant-mean-curvature slices, sourced by a continuous record, at Gaussian resolution R_0 and with least decoherence, of the spacelike part of the full constraint source; everything else, transverse-traceless exchange and constraint exchange at timelike momentum, is quantum. The spacelike part of the stress tensor of the quantum sector with exact four-momentum is

$$T_{\mu\nu}^{\text{SL}}(x) = \int \frac{d^4 q}{(2\pi)^4} \theta\left(k^2 - \frac{\omega^2}{c^2}\right) e^{iq \cdot x} \tilde{T}_{\mu\nu}(q); \quad (18)$$

it annihilates the vacuum to all orders, since the spectrum condition puts every $\langle n|\tilde{T}(q)|\Omega\rangle$ in the closed forward cone. Its free-field core, the $a^\dagger a$ collapse operator with a dark vacuum and a preferred ordering, is Pearle’s [9]; its floor $C \geq 1$, with $C = 1$ at least decoherence, is Tilloy and Diósi’s [19, 20], with Kafri, Taylor and Milburn’s classical channel [18] as the two-body precursor. The model is non-local, so it evades Theorem 4.4 by bringing its own ordering, the CMC slices, whose unit normal is u ; it signals at a level suppressed by G , through the non-local response of T^{SL} , in the frame of the slices. Its gravitational-wave sector, cosmology and costs are left to a separate paper; here we use only its emission spectrum, which is where a frame becomes observable.

The cut spectrum. A charge kicked by the record can emit a photon of energy E only if the transferred four-momentum is spacelike, $c|\mathbf{k}| > E/\hbar$. The white integrand $\int d^3k k^2 e^{-k^2 R_0^2/k^2}$ acquires the factor $\theta(c|\mathbf{k}| - E/\hbar)$, and

$$\frac{d\Gamma}{dE} = \frac{d\Gamma}{dE}\Big|_{\text{DP}} S\left(\frac{ER_0}{\hbar c}\right), \quad S(x) = \frac{4}{\sqrt{\pi}} \int_x^\infty u^2 e^{-u^2} du = \text{erfc } x + \frac{2x}{\sqrt{\pi}} e^{-x^2} \quad (19)$$

for the self term; the cut overlaps of the atomic pair terms are given in Appendix A. Below $\hbar c/R_0$ the spectrum is Diósi–Penrose’s, above it it falls as $e^{-(ER_0/\hbar c)^2}$ (Fig. 1a). The cut $q^2 > 0$ is Lorentz invariant; the frame enters through the smearing, which is in $\mathbf{k}$ and not in q^2 . By Proposition 5.3 no covariant smearing produces a falling spectrum: a Gaussian cutoff in the spontaneous X-ray spectrum, at the energy $\hbar c/R_0$ that also sets the heating, is the observable fingerprint of a frame. The direction of the laboratory’s motion through the frame leaves no usable trace: the cosmic-microwave-background velocity $\beta = 1.23 \times 10^{-3}$ changes the cut emission only at order $\beta^2 x^4$, below 10^{-4} for $x \leq 3$.

The window. A digitized-figure refit of the XENONnT SR0 spectrum with the cut emission (Appendix A) bounds the model at $C = 1$ by $R_0 > 1.93 \text{ \AA}$, the 95% limit constrained to the median sensitivity. The unconstrained limit, $2.11 \text{ \AA}$, is stronger only because the threshold bin fluctuates low, 4 events against 6.9 expected. Least decoherence makes the decoherence smearing equal to the smearing of Newton’s force, so the $52 \text{ }\mu\text{m}$ inverse-square-law test [40] caps $R_0 < 4.0 \times 10^{-5} \text{ m}$ in the infinite-slab approximation. The window

$$1.93 \times 10^{-10} \text{ m} < R_0 < 4.0 \times 10^{-5} \text{ m} \quad (C = 1) \quad (20)$$

spans 5.32 decades. MAJORANA and Gran Sasso no longer bind once the cut is applied: their edges fall to $0.18 \text{ \AA}$ and $0.006 \text{ \AA}$. At the lower edge $\hbar c/R_0 = 1.02 \text{ keV}$.

Detection. A cut signal near the edge is an excess concentrated in the first keV above threshold, over 80% in 1–2 keV. At the exposure that detects it at 5σ , about $1.3 \times 10^3 \text{ ty}$ for R_0 at 1.1 times the observed edge, the best fits of the Diósi–Penrose $1/E$ shape, the covariant E^2 shape and a coloured (Lorentzian) shape are each rejected at $\geq 4.9\sigma$, and the spectrum measures R_0 to $(2.03, 2.84) \text{ \AA}$ at 1σ with C free. That is a statistical floor: at the edge the signal sits where the detection efficiency is least certain, $\pm 18\%$ at 1.3 keV . The edge moves with C only as $C^{0.13}$ and with exposure hardly at all, to a median $2.93 \text{ \AA}$ at 10^4 ty ; above about $3 \text{ \AA}$ the window is dark to xenon detectors with a 1 keV threshold, and the lever is the threshold.

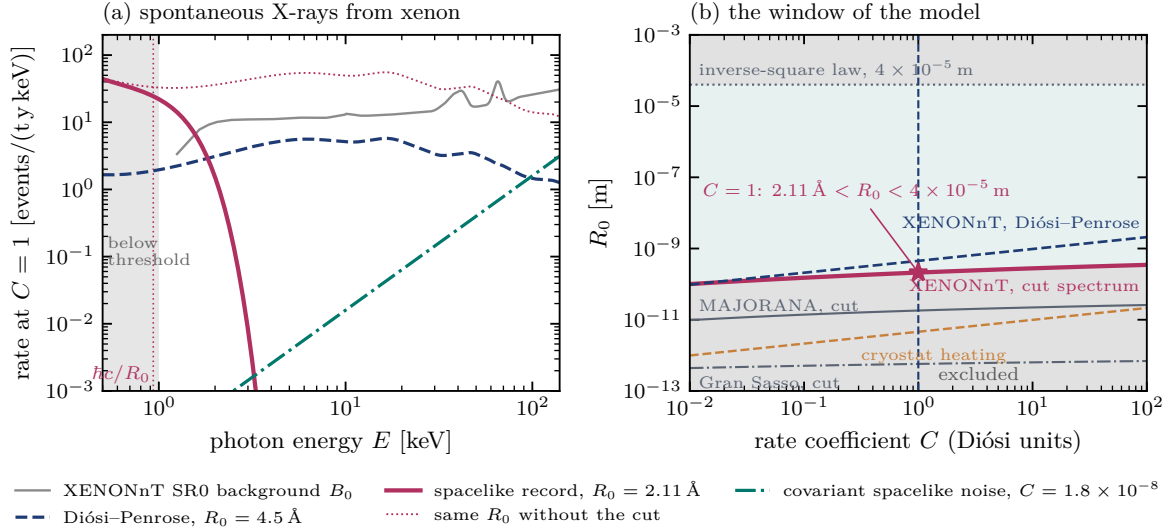


Figure 1. (a) Spontaneous X-ray emission of xenon at $C = 1$, before detector response. The Diósi-Penrose spectrum at the published XENONnT edge (dashed) falls slowly; the spacelike record at the observed refitted edge $R_0 = 2.11 \text{ \AA}$ (thick) equals the uncut spectrum (dotted) below $\hbar c/R_0$ and dies above it; the covariant spacelike noise at its XENONnT limit (dash-dotted) rises as E^2 . Grey: the SR0 background model. (b) The (C, R_0) plane of the model: observed lower edges from XENONnT with the cut spectrum (thick), from MAJORANA and Gran Sasso with the cut, and from cryostat heating; the dashed line is the uncut XENONnT edge; the upper edge is the inverse-square-law bound on the force smearing. The shaded band is allowed. The power-constrained XENONnT edge at $C = 1$ is $1.93 \text{ \AA}$.

8 Discussion

8.1 What an observation would mean

Theorem 2.2 turns a decoherence measurement into a statement about symmetry. If an isolated, inertial, equal-mass superposition decoheres at $\Gamma = CE_G/\hbar$ with C of order one, then a reciprocal bath behind it that meets the conditions of the drag theorem at kilometre scales has $T_{\text{eff}} > 3.3 \times 10^{16} C \text{ K}$, and the law of the noise is not Lorentz invariant. With matter the local quantum field theory we know, a vacuum that the noise does not excite without bound, and no further degrees of freedom, a noise law is not Lorentz invariant, and if its spectrum reaches timelike momenta its Lorentz stabilizer is compact and it has a rest frame; a record with feedback has non-covariant kernels. If instead the law reads its time direction from the state, it is a dynamical aether, nonlinear over mixtures. The two readings are not separated by present data. The escapes are named: an infrared modification of Newton’s law, non-Gaussian or jump records, non-local couplings, which bring their own foliation, dynamics without time ordering, which escape the fourth-order argument but not the X-ray bound, and new degrees of freedom whose law depends on the state.

The statement is independent of any model. Gravity enters only through the normalization of the noise, $4\pi C\hbar G/k^2$, and through the Newtonian infrared limit of hypothesis (G). The theorem says nothing about why the noise should be gravitational.

8.2 Relation to earlier work

The obstruction to covariant collapse is due to Ghirardi, Grassi and Pearle [7], Pearle [8, 9, 33], Myrvold [10], Jones, Gasbarri and Bassi [11] and Diósi [12, 13, 14, 15]; Lemma 4.1 is its form

for a decoherence noise and claims no new mechanism. The tachyonic spacelike noise, and the observation that it re-excites the vacuum beyond lowest order, are Pearle’s [8]. The pointer field and its state-dependent smearing are Bedingham’s [16]; the reading of that state dependence as the evasion of Myrvold’s theorem is in Ref. [37], and the role of a state-defined matter density in relativistic collapse is discussed in Ref. [39]. The classical-channel floor $C \geq 1$ and least decoherence are Tilloy and Diósi’s [19, 20], after Kafri, Taylor and Milburn [18], with the decoherence–diffusion trade-off of Oppenheim et al. [22] as its general form; the covariant classical–quantum theory and its Newtonian limit are theirs [21, 23, 24, 25]. The signalling risk of nonlinearity over mixtures is Gisin’s [26], and Polchinski’s construction shows it can be avoided [38]. Flash-sourced gravity [31] and Poissonian hybrids [30] are the jump-record alternatives. The free-field core of the model of Section 7 is Pearle’s quasi-relativistic collapse operator [9].

New here, within the sources we have read: the link of Theorem 4.4 from a causal, exactly Newtonian classical response with complete positivity to infinite vacuum excitation; the closure of the spacelike-support door for the Diósi–Penrose normalization, by fourth-order excitation and by X- and γ -ray data (Lemma 5.2, Propositions 5.3 and 5.4); and the refit of the XENONnT spectrum with the cut emission of a frame theory (Appendix A).

8.3 Open problems

- Case (b) for kernels invariant under a non-compact subgroup, such as $SO(2, 1)$, that are Newtonian only at $\omega = 0$ (Remark 4.2).
- The amplitude condition of (Q) in a general interacting theory, beyond the tree-level bilinears for which Lemma 5.2 verifies it.
- The general form of the invariant-pointer step: whether every Lorentz-invariant tempered distribution supported in the future cone has $|\tilde{g}|^2$ positive and locally integrable on the timelike shell (Section 6).
- A theorem for non-local couplings: whether every ordering-dependent generator needs a foliation, or an ordering-independent prescription exists.
- The curved-spacetime statement: frame or foliation for a coloured noise in Friedmann–Robertson–Walker spacetime.
- Whether a state-borne time direction signals.
- Jump records that keep the Diósi–Penrose scaling over the laboratory band.

9 Conclusions

A Diósi–Penrose decoherence rate breaks Lorentz invariance. For state-independent laws acting on local quantum field theory through a classical noise, the law is not Lorentz invariant when the vacuum is excited at a finite rate, and its stabilizer is compact if its spectrum reaches timelike momenta; for a Gaussian record with an exactly Newtonian response the kernels are not covariant. The one covariant noise that is dark at lowest order is excited at fourth order and is limited by X-ray data to $C < 2 \times 10^{-8}$. A law that reads its time from the state is a dynamical

aether. A frame shows as a Gaussian cutoff of the spontaneous X-ray spectrum near $\hbar c/R_0$, which a xenon detector can see if R_0 lies near 2 Å. A single-body interferometer in the saturated regime measures C ; a cutoff in the spontaneous X-ray spectrum would show the frame.

A The XENONnT refit

Data. The XENONnT collapse search [28] uses the SR0 electronic-recoil data and background model of Ref. [27], 1.16 t y over 1–140 keV, with nine background components. No numerical data release exists. The figures of Ref. [27] are vector graphics, from which we read the data points (1 keV bins over 1–30 keV, 2 keV bins over 1–140 keV), the background model and its eight plotted components, the total efficiency with its 1σ band, and the ^{37}Ar response. Every bin count times 1.16 t y is an integer to 0.02, 3658 events in all. The ^{37}Ar line fits a skew-normal response of width $0.279\sqrt{E/\text{keV}}$ keV and shape 1.61.

Emission model. The signal is the semiclassical spontaneous-emission rate of xenon of Piscicchia et al. [29], which XENONnT uses: the self terms $(N_p^2 + N_e)/E$ and the electron–electron and electron–nucleus pair terms, each with the Gaussian overlap $e^{-r^2/4R_0^2}$ and the photon phase $\text{sinc}(Er/\hbar c)$, with the orbital radii of Ref. [28]. At $R_0 = 4.5$ Å it reproduces XENONnT’s own Diósi–Penrose curve to 5% at 17, 47 and 140 keV, so it is $C = 1$ in Diósi units. The spacelike record restricts every overlap to $|\mathbf{k}| > E/\hbar c$. Since $e^{-r^2/4R_0^2} = (4/\sqrt{\pi}) \int_0^\infty u^2 e^{-u^2} j_0(ur/R_0) du$, the cut overlap is

$$F(s, x) = \frac{4}{\sqrt{\pi}} \int_x^\infty u^2 e^{-u^2} j_0(us) du = \frac{1}{s} \left[\frac{2}{\sqrt{\pi}} e^{-x^2} \sin sx + s \text{Re}(e^{-x^2+isx} w(\frac{s}{2} + ix)) \right], \quad (21)$$

with $s = r/R_0$, $x = ER_0/\hbar c$ and w the Faddeeva function; $F(0, x) = S(x)$ of Eq. (19).

Likelihood. Poisson over 84 bins; eight background components profiled with the published constraints (^{214}Pb free in its range, ^{85}Kr at $90 \pm 22\%$, materials, ^{136}Xe and solar neutrinos constrained, ^{124}Xe , ^{133}Xe and ^{83m}Kr free); one efficiency nuisance scaling every bin by its 1σ band; the signal folded with the response and the efficiency; true energies below 1 keV, where the emission model is not valid, dropped. The one-sided 95% limit uses $q = 2.706$, which reproduces XENONnT’s 90%-to-95% ratio. Applied to the uncut spectrum the fit gives 4.48 Å against the published 4.5 Å; the residual factor 1.012 on the signal is kept throughout.

Results. At $C = 1$ the cut spectrum gives $R_0 > 2.11$ Å (95% CL) observed and a median expected limit of 1.93 Å. The observed limit lies above the median because the threshold bin fluctuates low (4 events against 6.9 expected), so we quote the power-constrained 1.93 Å. Variants of the response, of the true-energy floor (0.5 or 1.5 keV) and of the efficiency treatment move the observed limit within 1.90–2.16 Å. At the observed edge the signal is 3.4 events, 2.8 of them in the 1–2 keV bin. The edge moves as $C^{0.13}$ rather than $C^{1/3}$ because the signal $\propto CR_0^{-3}S(x)$ falls steeply with R_0 . At the threshold, $x = E_{\text{th}}R_0/\hbar c \approx 1.07$, the slope $d \ln[R_0^{-3}S]/d \ln R_0$ would give $C^{0.21}$; the fitted 0.13 corresponds to an effective $x \approx 1.6$, about 1.5 keV, where the folded signal carries its weight after resolution and efficiency. Hence a factor-2 error in the threshold efficiency moves the edge by about 10%, and 10^4 times the SR0 exposure moves the median limit

only to $2.93 \text{ \AA}$. The same likelihood with the covariant spectrum of Proposition 5.4, screened or not, gives $C < 1.8 \times 10^{-8}$; its counts are driven by the top of the band.

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