

Spacelike-Record Gravity: A Candidate Theory of Diósi–Penrose Decoherence and What It Predicts

Marc Sperzel*

Independent Researcher

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Abstract

A companion paper showed that a Diósi–Penrose decoherence rate, $\Gamma = CE_G/\hbar$, requires a noise without back-action whose law is not Lorentz invariant. We write down a candidate theory that meets these conditions and work out what it predicts. On an assumed foliation by constant-mean-curvature slices, the classical part of geometry is the instantaneous constraint field sourced by a continuous record of the spacelike part of every constraint component, energy density, stress and momentum density of matter and of gravitational radiation, with a tidal term; the quantum sector couples to it only through these projected operators, and everything else is quantum. The vacuum stays dark at every order on flat backgrounds, the response does not drag, and least decoherence gives $C = 1$, at the floor $C \geq 1$ that complete positivity imposes for a Gaussian record with linear feedback, with stress and current superpositions decohering at least eight and four times faster than mass superpositions. Assuming a covariant mean field, the mean is linearized general relativity, and radiation reaction through 3.5PN order needs no record of the radiation. Data confine the one length to $1.93 \text{ \AA} < R_0 < 4.0 \times 10^{-5} \text{ m}$. The theory predicts a Gaussian cutoff of spontaneous X-ray emission near $\hbar c/R_0$, one kernel for decoherence and for the short-range force, null two-sphere entanglement tests and a moving-body rate $\gamma^2 \Gamma_{\text{rest}}$. The expansion rate is classical and the primordial tensors carry no classical stochastic component, but the force smearing lowers the scalar tilt of Starobinsky inflation at $N_* = 55$ to 3.6σ or more below Planck and 7.2σ or more below ACT DR6. The cosmological constant only decreases under the record’s heating, so its initial value is tuned, and $w + 1 \approx 1 \times 10^{-16}$. The costs are the assumed foliation, signalling suppressed by G in the preferred frame, an assumed Born rule, a record undefined at the light ring, and a burst-timing estimate that would raise the lower edge to $4.3 \text{ \AA}$.

1 Introduction

If a mass superposition loses coherence at the Diósi–Penrose rate [1, 2, 3]

$$\Gamma = \frac{C E_G}{\hbar}, \quad E_G = \frac{G}{2} \int \frac{d^3 k}{(2\pi)^3} \frac{4\pi}{k^2} e^{-k^2 R_0^2} |\Delta \rho_{\mathbf{k}}|^2, \quad (1)$$

*ORCID: 0009-0000-6252-3155. Email: me@marcsperzel.com

first order in G , with $\Delta\rho$ the difference of the branch mass densities and R_0 a Gaussian smearing length, two companion papers say what the noise behind it must be. It acts without back-action at kilometre scales: a reciprocal bath with the conditions of the drag theorem is bounded by the double pulsar to $T_{\text{eff}}/C > 3.3 \times 10^{16}$ K [4]. And its law is not Lorentz invariant: for local quantum field theory with a stable vacuum and a classical noise or a Gaussian record with an exactly Newtonian response, either the noise has a rest frame or the record’s kernels are not covariant [5]. Those are conditions, not a theory. They leave open whether any theory satisfies them while reproducing general relativity in the mean, keeping gravitational waves quantum and the vacuum dark, and surviving cosmology.

This paper writes down the least structure we have found that does, and works out what it predicts. The idea is one sentence:

The classical part of geometry is the instantaneous constraint field on constant-mean-curvature slices, sourced by a continuous record of the spacelike part of the full constraint source; everything else is quantum.

“Spacelike part” means the matrix elements of the source between states whose four-momenta differ by a spacelike vector. The full constraint source is every component of the stress tensor that the constraints see, energy density, transverse stress and momentum density, of matter and of the gravitational radiation on the slice, with a term linear in the radiation weighted by the background tidal field. The quantum sector couples to the classical field only through these projected operators. Least decoherence fixes the record’s strength. The theory has one free numerical constant, R_0 , and assumes, besides its rule, a foliation, a covariant mean field and the Born rule.

Given the foliation, the rule meets the companion theorem’s conditions in a minimal way. The classical field is instantaneous on the leaves, so its response is real and dissipationless, $T_{\text{eff}} = \infty$; the leaves are the time direction the theorem requires; the spacelike projection keeps the flat-space vacuum dark at every order. The coefficient is bounded below, $C \geq 1$, with $C = 1$ at least decoherence, so $\Gamma = E_G/\hbar$; stress and current superpositions then decohere at least eight and four times faster than mass superpositions of the same strength (Section 3). The mean field is linearized general relativity; radiation reaction through 3.5PN order needs no monitoring of the radiation; the record must contain the graviton energy and a tidal term (Section 4). Data confine R_0 to $1.93 \text{ \AA} < R_0 < 4.0 \times 10^{-5} \text{ m}$. The theory predicts a Gaussian cutoff of the spontaneous X-ray spectrum near $\hbar c/R_0$, one kernel for decoherence and for the short-range force, a null result in every two-sphere entanglement test with R_0 below the body size, and a moving-body rate $\gamma^2 \Gamma_{\text{rest}}$ in the frame of the slices (Section 5). In cosmology the expansion rate is classical and the primordial tensors carry no classical stochastic component; the force smearing turns off the metric response to inflaton modes for about 44 e -folds after horizon exit, which lowers the scalar tilt of Starobinsky inflation at $N_* = 55$ to 0.944–0.950, 3.6σ or more from Planck and 7.2σ or more from ACT DR6; the cosmological constant is an integration constant that the record’s heating drives down, so its initial value is tuned, and $w + 1 \approx 1 \times 10^{-16}$ today (Section 6).

The theory has costs, stated in Section 7: an assumed foliation, a preferred frame with ensemble-level signalling suppressed by G , an assumed Born rule, a dark vacuum shown only on flat backgrounds, an undefined record at the light ring, and an estimate from gamma-ray-burst timing that would raise the lower edge to $4.3 \text{ \AA}$. It is a candidate with tensions. Section 8

compresses it into three axioms.

Prior work. Penrose argued that a superposition of geometries is ill defined and decays at the rate $E_G/\hbar$ [3]; Diósi gave the master equation [1, 2]. Kafri, Taylor and Milburn modelled Newtonian gravity as a measurement-and-feedback channel [6]. Tilloy and Diósi sourced a classical Newtonian field by a continuous record of the mass density, recovered the Diósi–Penrose rate [7], and fixed $C = 1$ by a principle of least decoherence [8]. Oppenheim and collaborators built classical–quantum gravity, with the whole metric classical [9], and proved the decoherence–diffusion trade-off of which the floor $C \geq 1$ is the Newtonian specialisation [10, 11]. Pearle’s quasi-relativistic collapse model uses a number-conserving collapse operator that annihilates the vacuum, with a preferred ordering [13], after showing that spacelike noise excites the vacuum only beyond lowest order [12]. Anastopoulos and Hu argued that entanglement by Newtonian forces does not decide whether the true, radiative degrees of freedom of gravity are quantum [14]; Danielson, Satishchandran and Wald argued the converse for a quantum Newtonian field [15]. York showed that the constraint-determined part of the geometry and the conformal radiative data separate on constant-mean-curvature slices [16], and Gomes, Gryb and Kosłowski that general relativity in that gauge is dual to a theory with a preferred foliation and spatial conformal invariance, shape dynamics [17]. Martin and Vennin showed that continuous spontaneous localization during inflation is excluded by the cosmic microwave background unless the collapsed density contrast falls on super-horizon scales at least as fast as $(k/aH)^2$, the comoving density contrast being the natural exception [18].

What is new. Within the sources we have read: the rule with the full constraint source and with quantum constraint exchange at timelike momentum, and the Bondi-timing defect that forces the latter; radiation reaction through 3.5PN without monitoring the radiation; the tidal noise drive; the stress floor of least decoherence on the energy block and the vector floor; the moving-body rate; the inflation window and its tilt formula; the sign and size of the cosmological-constant drift; and the burst-timing estimate. The record–feedback hybrid, the floor, $C = 1$, the dark vacuum of the free-field projection, the York split and the shape-dynamics dictionary are from the cited work, and the minimal noise a non-entangling Newtonian model must carry has been bounded model-independently by Fabiano et al. [55]; Section 3 states how the floor here relates to that bound.

2 The rule

Signature is $(-+++)$, $q = (\omega, \mathbf{k})$ and $q^2 = k^2 - \omega^2/c^2$; q is spacelike for $q^2 > 0$. In Sections 2–4 we set $c = 1$ where no number is quoted.

2.1 The split on constant-mean-curvature slices

We assume that spacetime is foliated by slices of constant mean extrinsic curvature K (maximal slices, $K = 0$, for asymptotically flat data). Neither the existence of a preferred foliation nor its constant mean curvature is derived here, and constant-mean-curvature Cauchy surfaces need not exist in a general cosmological spacetime [62]. On each slice write the spatial metric in York’s conformal form, $g_{ij} = \psi^4 \hat{g}_{ij}$, with $\hat{g}_{ij} = \delta_{ij} + h_{ij}^{TT}$ about flat space and the trace-free momentum

split into a longitudinal and a transverse-traceless part. The Hamiltonian constraint becomes the Lichnerowicz–York equation for ψ , the momentum constraint an elliptic equation for the longitudinal momentum, and the condition $\partial_t K = \text{const}$ an elliptic equation for the lapse [16]. The conformal factor, the lapse and the shift are fixed on each slice by these equations once their sources are given; only h_{ij}^{TT} and its momentum propagate. In the post-Newtonian literature these are the ADM-TT coordinates [19]. At linear order in Poisson gauge,

$$ds^2 = -(1 + 2\Phi) dt^2 + 2V_i dx^i dt + [(1 - 2\Psi)\delta_{ij} + h_{ij}^{TT}] dx^i dx^j, \quad \partial_i V_i = 0, \quad (2)$$

the constraint data (Φ, Ψ, V_i) obey Poisson equations sourced by the energy density, the transverse current and the anisotropic stress, and $\square h_{ij}^{TT} = -16\pi G T_{ij}^{TT}$. Four of the six gauge-invariant components are constrained; two propagate.

The split depends on the slicing. A source moving uniformly with velocity $\mathbf{v}$ has a gauge-invariant transverse-traceless field whose share of the $v_i v_j$ part is $\frac{1}{2} \sin^4 \theta_k$, $4/15$ on average over directions. That field is off shell, $q^2 > 0$, since $q \cdot u = 0$ for the source’s four-velocity u ; on shell, $q^2 = 0$, only helicity ± 2 remains, and helicity is boost invariant (Section 4.1).

2.2 The spacelike part of an operator

Let $\bar{g}$ be a fixed, noiseless background metric and $\bar{H}$ the Hamiltonian of the quantum sector, matter and radiative data, on $\bar{g}$, with no coupling to the classical sector’s deviation from $\bar{g}$. For flat $\bar{g}$ let $P^\mu = (\bar{H}, \mathbf{P})$ be the generators of its translations, and $O(x) = e^{iP \cdot x} O(0) e^{-iP \cdot x}$ a Heisenberg operator. The projection below is defined with this P^μ , not with the noisy total dynamics.

Definition 2.1 (Spacelike part).

$$O^{\text{SL}}(x) = \int \frac{d^4 q}{(2\pi)^4} \theta(k^2 - \omega^2) e^{iq \cdot x} \tilde{O}(q), \quad \langle m | O^{\text{SL}}(x) | n \rangle = \theta\left((\mathbf{p}_m - \mathbf{p}_n)^2 - (E_m - E_n)^2\right) \langle m | O(x) | n \rangle, \quad (3)$$

with the open cone: matrix elements between states whose four-momenta differ by a null vector are dropped.

Proposition 2.2 (Dark vacuum). *Let $\bar{g}$ be flat and Ω the vacuum of $\bar{H}$. If the quantum sector obeys the spectrum condition, $O^{\text{SL}}(x)|\Omega\rangle = 0$ for every operator O and every x . Every product $O_1^{\text{SL}}(x_1) \cdots O_n^{\text{SL}}(x_n)$ annihilates Ω , so a record of O^{SL} with feedback through O^{SL} never excites the vacuum, at any order.*

Proof. $\langle n | O^{\text{SL}}(x) | \Omega \rangle$ carries θ of $(\mathbf{p}_n - \mathbf{p}_\Omega)^2 - (E_n - E_\Omega)^2$, and the spectrum condition puts $p_n - p_\Omega$ in the closed forward cone, where that argument is not positive. The product statement follows by applying the first statement to the rightmost factor. $\square$

For free fields $T_{\mu\nu}^{\text{SL}}$ is the number-conserving part of $:T_{\mu\nu}:$. A scattering term $a_p^\dagger a_p$ transfers $q = p' - p$ with $|\omega' - \omega| \leq |\mathbf{p}' - \mathbf{p}|$, spacelike except for collinear massless quanta; a pair term transfers $p + p'$ with $\omega + \omega' \geq |\mathbf{p} + \mathbf{p}'|$, timelike or null. The open cone drops the collinear set, of measure zero. This free-field core is Pearle’s quasi-relativistic collapse operator [13]. Normal ordering is not an alternative: a double commutator is blind to c-numbers, $[f + c, [f + c, \rho]] = [f, [f, \rho]]$, and the vacuum excitation of a monitored local density comes from its pair matrix

elements, which normal ordering keeps. No local operator annihilates the vacuum [20], so O^{SL} is non-local: in space, with Compton tails for massive fields and power-law tails for massless ones, and in time.

Two consequences fix where the projection acts. At $\mathbf{k} \rightarrow 0$, $T_{00}^{\text{SL}} \rightarrow H$, the total energy, so the passive gravitational mass is the total energy and active and passive masses agree. Every stationary source sits at $\omega = 0$, spacelike for $\mathbf{k} \neq 0$: stars, bound systems with their binding energy, thermal radiation through its cycle-averaged energy, acoustic waves with $\omega = c_s k < k$, and free streaming, whose transfers fill $|\omega| \leq k$. What is dropped are components of $\langle T \rangle$ oscillating faster than k , such as beats of a superposed atom at $k \lesssim 1/R_0$, which average to zero over any window longer than their period.

2.3 The model

Let n^μ be the unit normal of the leaves. The constraint source on a leaf is the stress tensor $\mathcal{T}^{\mu\nu}$ of matter and of the radiative data, whose energy density is

$$\mathcal{S} = n_\mu n_\nu T_{\text{matter}}^{\mu\nu} + \varepsilon_{TT} + \frac{1}{4\pi G} h_{ij}^{TT} \partial_i \partial_j \bar{\Phi}, \quad (4)$$

the matter energy density, the York energy density ε_{TT} of the radiative data on the slice, and the term of the Lichnerowicz–York operator linear in h^{TT} : the conformal Laplacian $\hat{g}^{ij} \partial_i \partial_j \psi$ contains $-h_{ij}^{TT} \partial_i \partial_j \bar{\psi}$, and with $\bar{\psi} = 1 - \bar{\Phi}/2$ at linear order this is the tidal term shown, weighted by the background tidal tensor $\partial_i \partial_j \bar{\Phi}$. Section 4 shows why each term is needed.

The record holds every component of $\mathcal{T}^{\mu\nu}$ that the constraints see. In Fourier space on the leaf, with $\hat{\mathbf{k}}$ the direction of $\mathbf{k}$, these are the energy density $\rho = \mathcal{S}$, the transverse half-trace of the stress $\sigma = \frac{1}{2}(\delta_{ij} - \hat{k}_i \hat{k}_j) \mathcal{T}^{ij}$, and the longitudinal and transverse momentum densities $j_L = \hat{\mathbf{k}} \cdot \mathbf{j}$ and $\mathbf{j}_T = \mathbf{j} - \hat{\mathbf{k}} j_L$. Write $X = (\rho, \sigma, j_L, \mathbf{j}_T)$. The instantaneous part of linearized graviton exchange between conserved sources (Eq. (11) below, with $\omega\rho = k j_L$) between sources X and X' is $X^\top \mathcal{K} X'$ with the matrix kernel

$$\mathcal{K}(\mathbf{k}) = \frac{8\pi G}{k^2} M, \quad M = \begin{pmatrix} 1/2 & 1 \\ 1 & 0 \end{pmatrix} \oplus \left(-\frac{1}{2}\right) \oplus (-2\mathbb{K}_T), \quad (5)$$

acting on (ρ, σ) , j_L and $\mathbf{j}_T$; its $\rho\rho$ entry is Newton’s $4\pi G/k^2$. In Poisson gauge the fields conjugate to ρ , σ and $\mathbf{j}_T$ are the Newtonian potential, the difference of the two scalar potentials that light and stresses feel, and the gravitomagnetic potential V_i .

Definition 2.3 (Spacelike-record gravity). (R1) *Quantum sector*. Matter, a local relativistic quantum field theory, and the radiative data (h^{TT}, π^{TT}) evolve under $\bar{H}$ on the background $\bar{g}$. They couple to the classical sector’s deviation from $\bar{g}$, mean and noise, only through (R3).

(R2) *Classical sector*. On each leaf, the constraint data are the instantaneous solution of the constraint equations with source the record

$$\varrho_t(\mathbf{x}) = \langle X(\mathbf{x}) \rangle_t + \delta\varrho_t(\mathbf{x}), \quad X = g_{R_0} * X^{\text{SL}}, \quad \langle \delta\varrho_t(\mathbf{x}) \delta\varrho_s^\top(\mathbf{y}) \rangle = \gamma^{-1}(\mathbf{x}, \mathbf{y}) \delta(t - s), \quad (6)$$

where X^{SL} is the spacelike part of every component, g_{R_0} a normalized Gaussian of width R_0 on the leaf, $\langle \cdot \rangle_t$ the conditional expectation given the record so far, t the leaf label,

and γ a positive matrix kernel. At linear order the classical field is $h_{\text{cl}} = \mathcal{K} * \varrho_t$; its ρ component is the Newtonian potential, $\Phi_{\text{cl}} = -G K * \varrho_{t,\rho}$ with $K(\mathbf{x}) = 1/|\mathbf{x}|$.

- (R3) *Feedback.* The classical field acts on the quantum sector through $H_{\text{fb}} = \int h_{\text{cl}}^\top X d^3x$, linearly, with a state-independent kernel, and only through the projected operators X .
- (R4) *Everything else is quantum:* transverse-traceless exchange at every q , and the constraint exchange of $\mathcal{T} - \mathcal{T}^{\text{SL}}$, which has timelike q , as a unitary Coulomb-type term.
- (R5) *Least decoherence.* γ minimizes the decoherence kernel first on the energy block and then on the remaining blocks (Section 3).

Averaging over records gives

$$\dot{\rho} = -\frac{i}{\hbar}[\bar{H} + V, \rho] - \iint L_{AB}(\mathbf{x}, \mathbf{y}) [X_A(\mathbf{x}), [X_B(\mathbf{y}), \rho]] d^3x d^3y, \quad L = \frac{\gamma}{8} + \frac{1}{2\hbar^2} \mathcal{K} \gamma^{-1} \mathcal{K}, \quad (7)$$

with $V = -\frac{1}{2} \iint X^\top \mathcal{K} X$, symmetrically ordered, the instantaneous pair interaction. The first term of L is the price of monitoring; the second is the noise of the classical field acting back on matter. On the energy block, with $f \equiv X_\rho$ and $\mathcal{K}_{\rho\rho} = GK$, this is the master equation of Tilloy and Diósi [7, 8], $L = \gamma/8 + (G^2/2\hbar^2) K \gamma^{-1} K$ and $V = -\frac{G}{2} \iint f K f$. It is of Gorini–Kossakowski–Lindblad–Sudarshan form for any positive kernel, whether or not $X_A(\mathbf{x})$ and $X_B(\mathbf{y})$ commute, so the dynamics is completely positive. Every noise term enters through the projected operators X , so by Proposition 2.2 the vacuum of $\bar{H}$ on flat $\bar{g}$ is a stationary state of Eq. (7), with no cutoff. On curved backgrounds the statement is not established (Section 7).

The vacuum is dark; excited states are not. A quantum of momentum p is heated per unit energy at κ times the rate per unit mass of non-relativistic matter, $P/M = G\hbar/(4\sqrt{\pi}R_0^3)$, with $\kappa \rightarrow 2/3$ for $pR_0 \ll \hbar$, $\kappa \rightarrow 4/3$ for $pR_0 \gg \hbar$, monotone in between, for a minimally coupled scalar in the density channel; spin, curvature-coupling and stress-channel factors are not computed.

2.4 How the rule realises the theorem

Ref. [5] shows that a Diósi–Penrose rate requires noise without back-action and a law that is not Lorentz invariant, and names the escapes. The rule meets each requirement with one structure.

No back-action. The classical field responds to the record instantaneously and with a real kernel. Its dissipative part vanishes, $\chi'' = 0$, which is the member $T_{\text{eff}} = \infty$ of the class the drag theorem leaves [4]. Moving bodies are not dragged by the constraint sector; their radiation reaction is the quantum transverse-traceless sector’s, which is general relativity’s (Section 4.3).

A time direction. The theory assumes a foliation. The cone theorem shows why a classical field that is causal, non-radiating and Newtonian has nowhere else to live; it does not show that such a field exists, and existence is not derived.

Proposition 2.4 (Cone theorem). *Let D be a Lorentz-invariant distribution with $\text{supp } D \subset J^+$ (causal response) and $D(x) = D(-x)$ (no absorptive part, so no classical radiation). Then $\text{supp } D \subset \{0\}$, and $\tilde{D}(q)$ is a polynomial in q^2 : there is no Newtonian $1/k^2$.*

Proof. $D(x) = D(-x)$ puts $\text{supp } D$ in $J^+ \cap J^- = \{0\}$. A distribution supported at a point is a finite sum of derivatives of δ^4 , and Lorentz invariance leaves powers of $\square$. $\square$

With a foliation, J^+ is replaced by the half-space $t \geq 0$, whose intersection with its reflection is the leaf $t = 0$; the instantaneous kernel $\delta(t)K(\mathbf{x})$ is then allowed, and it is the only kind. A causal kernel with an absorptive part avoids the theorem but radiates: by Kramers–Kronig its static Newtonian value is the dispersion integral of the on-shell absorption, and the record noise, a causally filtered innovation, cannot be band-limited to $|\omega| < k$ (Paley–Wiener), so it pumps classical gravitational waves, at a rate that diverges for a covariant white record. The instantaneous kernel on leaves is therefore the one way, given a foliation, to have a classical Newtonian field with $\chi'' = 0$. Its kernels are not Lorentz covariant, as case (b) of the companion theorem requires (Theorem 2.2 of Ref. [5]).

A dark vacuum. X is non-local, so the model falls under the companion’s escape for non-local couplings, which bring their own ordering surface (Section 4.4 of Ref. [5]): here the leaves. Proposition 2.2 then keeps the vacuum dark at every order, including the fourth order at which a Lorentz-invariant spacelike noise re-excites it (Lemma 5.2 of Ref. [5]): there the noise is a c-number with spacelike support acting on a local operator, and two spacelike transfers sum to a timelike one; here every factor is projected with the exact spectrum.

The leaves are fixed by the classical geometry, which the record sources. In the language of Ref. [5] the time direction is borne by the state of the classical sector, the global geometric member of the state-borne branch, while the noise law on a given leaf does not depend on the quantum state.

3 The coefficient: a floor and least decoherence

Take first the energy block alone: record and feedback through $f = X_\rho$, with K the Gaussian-smeared Newton kernel. The decoherence rate of a superposition with smeared energy-density difference Δf is $\Gamma = \langle \Delta f, L \Delta f \rangle$ by Eq. (7), and in Diósi units $C = 1$ means $L = (G/2\hbar)K$, which gives $\Gamma = E_G/\hbar$ with E_G of Eq. (1).

Theorem 3.1 (Floor). *For every positive record kernel γ , with or without translation invariance,*

$$L(\gamma) \succeq \frac{G}{2\hbar} K \quad (\text{Loewner order}), \quad \text{with equality iff } \gamma = \frac{2G}{\hbar} K. \quad (8)$$

Hence $C \geq 1$ for every superposition, and $C = 1$ at a single record kernel, the same for all superpositions.

Proof. Put $A = \gamma/8$ and $B = (G/2\hbar)K$, both symmetric and $A \succ 0$, so that $L = A + \frac{1}{4}BA^{-1}B$. Then

$$(A^{1/2} - \frac{1}{2}A^{-1/2}B)^\top (A^{1/2} - \frac{1}{2}A^{-1/2}B) = A - B + \frac{1}{4}BA^{-1}B = L - B \succeq 0, \quad (9)$$

and the left side vanishes iff $A^{1/2} = \frac{1}{2}A^{-1/2}B$, i.e. $A = B/2$, i.e. $\gamma = 2GK/\hbar$. $\square$

The minimum is Tilloy and Diósi’s principle of least decoherence [8], which they state pointwise in Fourier space; the operator form needs no translation invariance. The inequality is the decoherence–diffusion trade-off of Oppenheim et al. [10, 11] specialised to a Newtonian field, proved by completing the square as for a standard quantum limit. With the gain ratio $x = \gamma\hbar/(2GK)$ on a mode, $C(x) = \frac{1}{2}(x + 1/x)$; at the minimum the two halves of L are equal and the classical field carries the source-independent noise $\langle \delta\Phi \delta\Phi \rangle = (G\hbar/2)K$.

Corollary 3.2. (a) Inefficient record. If a fraction $1 - \eta$ of the record is lost, $C(x, \eta) = \frac{1}{2}(x + 1/(\eta x))$ and $C_{\min} = 1/\sqrt{\eta}$. $C = 1$ needs $\eta = 1$ and $x = 1$.

(b) Two smearings. If the decoherence kernel is $Ce^{-k^2 R_D^2}/k^2$ and the force kernel $e^{-k^2 R_F^2}/k^2$ (translation invariance), the floor requires $C \geq 1$, fixed at $k \rightarrow 0$ for every smearing, and $R_D \leq R_F$.

(c) One kernel. At least decoherence $R_D = R_F$: the length that sets the decoherence of small bodies is the length that smears Newton's force.

Proof. (a) The monitoring term is unchanged, but only the fraction η of the record reaches the field, which raises the field noise by $1/\eta$: $L = (GK/4\hbar)(x + 1/(\eta x))$; minimize over x . (b) Pointwise in k the floor reads $Ce^{-k^2 R_D^2} \geq e^{-k^2 R_F^2}$ for all k ; at $k \rightarrow 0$ this is $C \geq 1$, and as $k \rightarrow \infty$ it needs $R_D \leq R_F$. (c) Equality in Eq. (8) makes the decoherence kernel equal to the realized force kernel. $\square$

Least decoherence is a selection rule. It picks one member of a family of consistent models; no inconsistency is derived for $C > 1$. It is the only principle of its kind we have found that gives a finite answer. Least information recorded about matter pushes $x \rightarrow 0$ and $C \rightarrow \infty$; maximal entropy production has no extremum, since $C(x)$ is unbounded at both ends; and least mutual information between record and matter fails in general, since stronger decoherence can freeze a driven state and produce less entropy (the quantum Zeno effect), and for mixed initial states the minimizer moves to $x \rightarrow 0$. What survives is the Loewner minimum of the rate at fixed state, which is the statement of Theorem 3.1.

Relation to minimal-noise bounds. Fabiano, Fujita, Matsumura and Carney classify the Galilean-invariant, time-local master equations that reproduce Newton's law in the mean, and show that any member that does not entangle two masses must inject a minimal noise, below which the interaction must be entangling [55]; Kafri, Taylor and Milburn [6] and Carney and Taylor [56] gave earlier versions. The Newtonian sector of this theory is a member of their class, with the Newton potential in the Hamiltonian, a single-body dissipator with kernel $e^{-k^2 R_0^2}/k^2$ that is not truncated to quadratic order in the positions, and a correlated two-body term. The two bounds differ in scope and in what they bound. Theirs holds for every model of the class and is a two-body statement: it compares the noise on each body with the entangling coupling Gm_1m_2/d^3 of a pair, so it is a criterion for entanglement. Theorem 3.1 holds only for a classical field sourced by a Gaussian record with linear feedback, and inside that class it bounds the decoherence of a single body by its own self-energy, $\Gamma \geq E_G/\hbar$, with no partner and for every superposition at once. The criterion of Proposition 5.1 has the structure of their two-state criterion (their Eq. 35), with the correlated decoherence of the record in the role of their correlated noise; for bodies larger than R_0 it shows that a model at the floor $C = 1$ stays above their threshold, so the theory is non-entangling as their result requires. Their bound does not imply $C \geq 1$: inside the Diósi–Penrose family, non-entanglement of two spheres with R_0 below their radii requires by Proposition 5.1 only $C \geq 5/\sqrt{119} = 0.458$, and that is all a non-entanglement requirement can give. Theorem 3.1 does not extend to their whole class: whether every non-entangling member of it has a single-body rate of at least $E_G/\hbar$ is open. The mathematics of Theorem 3.1 is Tilloy and Diósi's and the trade-off's; what this paper adds is its extension to the full constraint source below.

3.1 The full record: the energy block, the stress floor and the vector floor

With the matrix kernel of Eq. (5), write the decoherence kernel of Eq. (7) per mode as

$$L = \frac{2\pi G}{\hbar k^2} \Lambda(A), \quad \Lambda(A) = A + MA^{-1}M, \quad A = \frac{\hbar k^2}{16\pi G} \gamma \succ 0, \quad (10)$$

so that $C = 1$ for mass means $\Lambda_{\rho\rho} = 2M_{\rho\rho} = 1$. M is indefinite: its scalar block $M_s = \begin{pmatrix} 1/2 & 1 \\ 1 & 0 \end{pmatrix}$ on (ρ, σ) has determinant -1 , and its vector blocks are negative.

There is no least kernel. For $A = \text{diag}(a, a_s)$ on the scalar block, $\Lambda_{\sigma\sigma} = a_s + 1/a$ and $\Lambda_{\rho\rho} = a + 1/(4a) + 1/a_s$. So $\inf_A \Lambda_{\sigma\sigma} = 0$, approached as $a \rightarrow \infty$, $a_s \rightarrow 0$, where $\Lambda_{\rho\rho} \rightarrow \infty$. A smallest member of the family would minimize every diagonal entry at once; by the proposition below, minimizing $\Lambda_{\rho\rho}$ forces $\Lambda_{\sigma\sigma} \geq 8$. Least decoherence must therefore name a block, and a positive block needs a timelike vector. Rule (R5) names the energy density on the leaves, the block that carries Newton's law.

Proposition 3.3 (Stress floor). *For every $A \succ 0$, $\Lambda_{\rho\rho} \geq 1$, so $C \geq 1$ for mass with the full record. The minimum $\Lambda_{\rho\rho} = 1$ is attained, and it forces $\Lambda_{\sigma\sigma} \geq 8\Lambda_{\rho\rho}$.*

Proof. By Cauchy–Schwarz, $(e_\rho^\top M e_\rho)^2 \leq (e_\rho^\top A e_\rho)(e_\rho^\top M A^{-1} M e_\rho)$, so $\Lambda_{\rho\rho} \geq a + M_{\rho\rho}^2/a \geq 2M_{\rho\rho} = 1$ with $a = A_{\rho\rho}$, and equality holds iff $A e_\rho = M e_\rho$, i.e. $A_{\rho\rho} = \frac{1}{2}$, $A_{\rho\sigma} = 1$ and no other entry in the ρ column. Positivity of A then needs $a_s \equiv A_{\sigma\sigma} > 2$. Since $M e_\sigma = e_\rho$, $\Lambda_{\sigma\sigma} = a_s + (A^{-1})_{\rho\rho} = a_s + a_s/(a_s/2 - 1) = a_s + 2a_s/(a_s - 2)$, whose minimum over $a_s > 2$ is 8, at $a_s = 4$. $\square$

Proposition 3.4 (Vector floor). *Let P project onto a negative-definite block of M , with $B = PMP$ and $MP = PM$. For every $A \succ 0$, $P\Lambda P \succeq 2|B|$, with equality at $PAP = |B|$. The vector blocks decouple from the scalar block, so $\Lambda_{LL} \geq 1$ and $\Lambda_{TT} \succeq 4\mathbb{K}_T$ are attained together with $\Lambda_{\rho\rho} = 1$.*

Proof. Since M commutes with P , $PAP = \bar{A} + B(PA^{-1}P)B$ with $\bar{A} = PAP \succ 0$ on the range of P . By the Schur complement, $PA^{-1}P$ restricted to that range is the inverse of $\bar{A} - A_{PQ}A_{QQ}^{-1}A_{QP} \preceq \bar{A}$, so $PA^{-1}P \succeq \bar{A}^{-1}$ and $P\Lambda P \succeq \bar{A} + B\bar{A}^{-1}B$. Then $(\bar{A}^{1/2} + \bar{A}^{-1/2}B)^\top (\bar{A}^{1/2} + \bar{A}^{-1/2}B) = \bar{A} + 2B + B\bar{A}^{-1}B \succeq 0$ gives $P\Lambda P \succeq -2B = 2|B|$, with equality iff $\bar{A} = -B$ and A has no entries between the block and the rest. With $M_{LL} = -\frac{1}{2}$ and $M_{TT} = -2\mathbb{K}_T$ the floors are 1 and 4; M has no entries between the blocks, so a block-diagonal A reaches all three floors at once. $\square$

A superposition that differs in transverse stress decoheres at least eight times faster than a mass superposition with the same source amplitude, and one that differs in transverse momentum density, $\mathbf{j}_T/c$ against ρ , at least four times faster. For laboratory matter $\sigma \sim p \ll \rho c^2$ and $|\mathbf{j}| \sim \rho v$, so both are negligible; for radiation, $\sigma \sim \rho/3$, the stress channel raises the heating per unit energy by an order-one factor.

4 Consistency with general relativity

4.1 The mean field

Linearized one-graviton exchange between conserved sources is $W = 8\pi G T^{\mu\nu} P_{\mu\nu\rho\sigma} T'^{\rho\sigma}/q^2$, with $P_{\mu\nu\rho\sigma} = \frac{1}{2}(\eta_{\mu\rho}\eta_{\nu\sigma} + \eta_{\mu\sigma}\eta_{\nu\rho} - \eta_{\mu\nu}\eta_{\rho\sigma})$, so that the numerator is $T \cdot T' - \frac{1}{2}TT'$. Decompose in

the frame of q : $T^{00} = \rho$, $T^{0i} = j_L \hat{k}_i + j_{T,i}$, and $\Pi_{ij} = \Pi_{ij}^{TT} + (\text{vector}) + A \hat{k}_i \hat{k}_j + \sigma(\delta_{ij} - \hat{k}_i \hat{k}_j)$. Conservation, $\omega \rho = \mathbf{k} \cdot \mathbf{j}$ and $\omega j_i = k_j \Pi_{ij}$, gives $A = (\omega^2/k^2)\rho$ and

$$\frac{T \cdot T - \frac{1}{2}T^2}{q^2} = \frac{|\Pi^{TT}|^2}{q^2} + \frac{\frac{1}{2}(1 - \omega^2/k^2)\rho^2 + 2\rho\sigma - 2|\mathbf{j}_T|^2}{k^2}. \quad (11)$$

Only the transverse-traceless term has a pole. The rest is instantaneous, a polynomial in ω^2/k^2 over k^2 ; with $\omega\rho = kj_L$ it is the matrix kernel of Eq. (5). On shell the numerator is $|\Pi^{TT}|^2$ for every conserved source, so the radiative exchange is pure helicity 2 in every frame. The time-symmetric propagator $\frac{1}{2}(D_{\text{ret}} + D_{\text{adv}}) = \text{PV } 1/q^2$ is negative on timelike q , so it cannot be a decoherence kernel; the admissible classical kernel is $\theta(q^2)/q^2$, the spacelike support, on which retarded and advanced propagators agree.

The classical sector is sourced by the record of every component of X^{SL} through the matrix kernel $\mathcal{K}$, which is the full instantaneous part of Eq. (11); rule (R4) carries the constraint exchange of $\mathcal{T} - \mathcal{T}^{\text{SL}}$ and all transverse-traceless exchange in the quantum sector. The two pieces have disjoint q -support, so the linear exchange splits exactly, $W = W[\mathcal{T}^{\text{SL}}] + W[\mathcal{T} - \mathcal{T}^{\text{SL}}]$, and the mean metric is linearized general relativity, all of whose potentials, Φ , Ψ and V_i , are fed back. A record of the energy density alone would give the Newtonian potential and not the stress and current couplings that set $1 + \gamma$ and the gravitomagnetic field. Two assumptions enter: the gain of the classical sector equals the coupling of the transverse-traceless sector, a mean-field covariance that is equivalent to, not derived from, the statement that the boosted static field solves linearized general relativity, and is fixed empirically by the preferred-frame parameter, $|\alpha_2| < 4 \times 10^{-7}$ from the Sun's spin axis [63] and $|\hat{\alpha}_2| < 1.6 \times 10^{-9}$ from solitary pulsars [64]; and the equal-time Hamiltonian form of the quantum timelike exchange, whose cross term with the classical sector is not worked out here.

Why timelike constraint exchange must be quantum. Drop (R4)'s constraint term and source the classical field by $\mathcal{S}^{\text{SL}}$ alone. For a sudden spherical shell of radiation of energy E emitted at $t = 0$, a Fourier mode of the source is $M - E \theta(t)(1 - j_0(ckt))$, and the spacelike record keeps only its part with $|\omega| < ck$. The recorded mass inside a sphere of radius r , at $s = ct/r$, then differs from Gauss's law:

$s = ct/r$	−0.5 (before emission)	0.5	0.8	1.5
general relativity	0	0	0	1
spacelike source only	−0.10	0.10	0.15	≈ 1

in units of E : the mean monopole rings acausally at the ten-percent level before the shell arrives and before it is emitted. An emission of duration τ_e suppresses the timelike weight of a mode as $e^{-W^2/2}$ with $W = ck\tau_e$, below 10^{-8} at $W = 6$, so the defect lives at $r \gtrsim c\tau_e/3$, where the Bondi mass is defined [21]. It is not observable (for GW150914 the potential shift at the detector is $\sim 3 \times 10^{-22}$ and constant), but it is a defect of the mean field; rule (R4) removes it and restores Gauss's and Birkhoff's laws.

4.2 First post-Newtonian order

In ADM-TT variables $h^{TT} = \mathcal{O}(c^{-4})$ [19]. For a circular binary the near-zone quadrupole field obeys $h^{TT}/(\Phi/c^2) = 2v^2/c^2$, so its first contribution to the Hamiltonian is at 2PN. By this power

counting the Einstein–Infeld–Hoffmann dynamics involves the constraint sector alone. We do not compute it here: we assume that the classical sector solves the full non-linear constraint equations with the recorded source, in which case the first post-Newtonian dynamics is general relativity’s, with $\beta = \gamma = 1$, the perihelion advance of Mercury, $42.99''$ per century, the Shapiro delay and light deflection with $1 + \gamma = 2$ (which needs the recorded stress, Section 2.3), and the Nordtvedt parameter, since gravitational binding energy is part of the classical source. The Lichnerowicz–York non-linearity holds pathwise for a record with correlation time R_0/c , whose noise variance $\text{Var}(\Phi/c^2) \simeq \ell_P^2/(\sqrt{\pi}R_0^2) = 4 \times 10^{-51}$ at $R_0 = 1.93 \text{ \AA}$ shifts β by nothing measurable. Photons and gravitons are quantum fields on the same classical metric, so they share light cone and deflection, $c_{\text{GW}} = c$ and $\gamma_{\text{GW}} = \gamma_{\text{EM}}$, in agreement with GW170817 and GRB 170817A, $-3 \times 10^{-15} < \Delta v/c < 7 \times 10^{-16}$ [22].

4.3 Radiation reaction through 3.5PN order

Proposition 4.1 (Reaction without monitoring the radiation). *In the mean, the dissipative dynamics of a binary is general relativity’s through 3.5PN order, with no record of h^{TT} .*

Proof. The reactive, time-odd part of the exchange is $\pi\delta(q^2)$ times the numerator of Eq. (11), which on shell is $|\Pi^{TT}|^2$; at linear order every multipole radiates and is damped through the quantum transverse-traceless sector, which acts on quantum matter through a Hamiltonian. The remaining piece is the field-stress source of h^{TT} , half of the quadrupole formula. The Newtonian stress $t^{ij} = (4\pi G)^{-1}(\partial_i\Phi\partial_j\Phi - \frac{1}{2}\delta_{ij}|\nabla\Phi|^2)$ obeys $\partial_j t^{ij} = \rho\partial_i\Phi$, so after an integration by parts

$$\int t^{ij} d^3x = - \int \rho x^{(i}\partial^{j)}\Phi d^3x. \quad (12)$$

The coupling $h_{ij} \int t^{ij}$ becomes $h_{ij} \int \hat{\rho} x^{(i}\partial^{j)}\Phi_{\text{cl}}$: a quantum h^{TT} coupled to quantum matter with a classical coefficient, which needs no record of h^{TT} . Its mean part, with $\bar{\Phi}$, is a Hamiltonian; its noise part, with the record noise in Φ_{cl} , is feedback, and it is the tidal drive of Eq. (13) below. The rewriting is exact while the radiative multipole moments have compact support. Mass moments acquire non-compact contributions only from 2PN order, current moments from the first correction [23], so it holds for the Newtonian mass quadrupole (2.5PN reaction) and for the 1PN mass quadrupole, the Newtonian mass octupole and the Newtonian current quadrupole (3.5PN). $\square$

The rewriting fails first at 4.5PN, through the 1PN current and 2PN mass moments. Tails, the scattering of the waves off the classical monopole (1.5PN in the waveform, 4PN in the reaction), are a Hamiltonian for the quantum field on a classical background; non-linear memory is a vertex of the quantum sector alone. Both are general relativity’s in the mean.

4.4 Graviton energy and the tidal term

Graviton energy. The constraint equations are consistent with the evolution only if their source is conserved (the contracted Bianchi identity). Radiating matter alone is not: $dE_{\text{matter}}/dt = -L_{\text{GW}}$. The record must contain the energy density of the radiation, which on the leaves is the York energy density ε_{TT} of Eq. (4); Bondi’s mass loss [21] of the mean then follows by Gauss’s law. Logical consistency forces this term; no datum does.

The tidal term. On a curved background the Lichnerowicz–York operator contains $h_{ij}^{TT} \partial_i \partial_j \bar{\psi}$, linear in the radiation and of order one near a black hole; in binary dynamics it enters at 4.5PN. Without it in the record the constraint sector does not respond to h^{TT} , the coupled perturbation system is non-reciprocal, and its quasinormal spectrum is not general relativity’s. The size of that shift is not computed here; the requirement rests on consistency, and GW250114, which places the dominant quadrupolar frequency within a few percent of general relativity [24], is its empirical stake. Recording the tidal term is a linear record of $h_{ij}^{TT} \mathcal{E}_{ij}$, with $\mathcal{E}_{ij} = \partial_i \partial_j \bar{\Phi}/c^2$. Its back-action is a c-number noise drive of the radiation: the gradient part of the noise stress has no transverse-traceless part, so a uniform field does not radiate, and with $\langle |\Lambda^{TT} \mathcal{E}|^2 \rangle = \frac{2}{5} |\mathcal{E}|^2$ over directions the radiated power per unit volume is

$$P_V = \frac{2}{5} \frac{\hbar c^2 |\mathcal{E}|^2}{\sqrt{\pi} R_0}. \quad (13)$$

It is emitted at $k \sim 1/R_0$, is independent of G at fixed $\mathcal{E}$, and vanishes as $\hbar \rightarrow 0$. Outside the surface, at $R_0 = 1.93 \text{ \AA}$, it is $2 \times 10^{-32} \text{ W}$ for the Earth, $7 \times 10^{-13} \text{ W}$ for a neutron star, and $4 \times 10^{-13} \text{ W}$ for the GW150914 remnant outside its horizon, 10^{-62} of the ringdown peak power.

The decoherence of gravitational waves. Gravitons are decohered by the record as matter of the same energy density is: two radiation states whose energy densities on the leaves differ by $\Delta \varepsilon$ lose coherence at $E_G[\Delta \varepsilon/c^2]/\hbar$, and near a mass the tidal term also records h^{TT} linearly. The effect is real and small for the waves detectors see. The record noise reaches h^{TT} only through the tidal drive of Eq. (13), a c-number source at $k \sim 1/R_0$; in the detection band it gives a strain noise of $3 \times 10^{-52} \text{ Hz}^{-1/2}$ at 100 Hz near the Earth, so there is no classical diffusion of the strain at any detectable level, unlike in a theory with the whole metric classical [9, 11]. A wave of frequency ω crossing a distance L acquires a phase variance at most $(\omega/c)^2 L \ell_P^2 K(0)$, with $K(0) = 1/(\sqrt{\pi} R_0)$, and is scattered out of its mode with probability $\ell_P^2 \omega L / (3\pi R_0^2 c)$. At the lower edge these are below $2.1 \times 10^{-23} \text{ rad}$ and 5×10^{-32} for GW150914 and GW170817, against the radian a matched filter resolves.

The slicing also resolves the Weinberg–Witten obstruction to a covariant graviton energy in the record, and the foliation survives gravitational collapse; both are shown in Appendix A, with the one place the theory is incomplete, the record at the light ring.

5 Laboratory predictions

5.1 The rate, and the window allowed by data

At least decoherence the rate is $\Gamma = E_G/\hbar$ exactly, $C = 1$, with E_G of Eq. (1) built from the energy density on the leaves. For bodies larger than R_0 it does not depend on R_0 : a uniform sphere of mass M and radius a displaced by $d \geq 2a$ has $E_G = GM^2(6/(5a) - 1/d)$, and a $1 \mu\text{g}$ silica sphere decoheres in 62.8 ps as $d \rightarrow \infty$, 65.4 ps at $d = 1 \text{ mm}$ [4]. A single body in this saturated regime measures C ; Theorem 3.1 makes any $C < 1$ fatal to the theory, and least decoherence predicts $C = 1$.

The theory does not predict R_0 ; data confine it. Corollary 3.2(c) sets the decoherence smearing equal to the smearing of Newton’s force, so short-range tests of the inverse-square law bound R_0 from above, as Tilloy and Diósi [8] and Tilloy and Stace [58] noted for the smearing of their models, at about 10^{-4} m . The $52 \mu\text{m}$ torsion-balance test [29], matched to its Yukawa

bound at $38.6 \mu\text{m}$ in the infinite-slab approximation [5], gives $R_0 < 4.0 \times 10^{-5} \text{ m}$; the 120-fold hole pattern of the test bodies changes this by less than 1%. An earlier estimate with finite lateral size gave $4.4 \times 10^{-5} \text{ m}$; we have not reproduced it, and a fit to the full torque calculation is not done here. A model whose mean field is not carried by the record has no such edge. From below, the spacelike projection cuts the spontaneous X-ray emission of the Diósi–Penrose noise: a charge kicked by the record can emit a photon of energy E only if the transferred four-momentum is spacelike, $c|\mathbf{k}| > E/\hbar$, so the emission spectrum is multiplied by

$$S\left(\frac{ER_0}{\hbar c}\right), \quad S(x) = \operatorname{erfc} x + \frac{2x}{\sqrt{\pi}} e^{-x^2}, \quad (14)$$

for the self term (Section 7 of Ref. [5]). Without the cut, the X-ray searches of XENONnT [48, 49], MAJORANA [50, 51] and the Gran Sasso germanium detector [52] bound the Diósi–Penrose smearing, with XENONnT strongest at $4.5 \text{ \AA}$; with the cut only XENONnT binds. The refit of the XENONnT electronic-recoil spectrum with this cut, in Ref. [5], gives $R_0 > 1.93 \text{ \AA}$ at $C = 1$, the 95% limit constrained to the median sensitivity ($2.11 \text{ \AA}$ unconstrained). The allowed window is

$$1.93 \times 10^{-10} \text{ m} < R_0 < 4.0 \times 10^{-5} \text{ m}, \quad (15)$$

5.32 decades. At the lower edge the cutoff energy is $\hbar c/R_0 = 1.02 \text{ keV}$ and the heating is $P/M = G\hbar/(4\sqrt{\pi}R_0^3) = 1.4 \times 10^{-16} \text{ W kg}^{-1}$, 7×10^4 below the 10 pW kg^{-1} cryostat bound [54]. A Diósi–Penrose-strength X-ray excess with no Gaussian cutoff near $\hbar c/R_0$ would exclude the theory.

5.2 One kernel, two experiments

By Corollary 3.2(c) the same Gaussian kernel appears in the decoherence and in the force. For two point-like branches separated by d the decoherence energy and the pair potential are

$$E(d) = Gm^2 \left[\frac{1}{\sqrt{\pi} R_0} - \frac{\operatorname{erf}(d/2R_0)}{d} \right], \quad U(r) = -\frac{Gm_1m_2 \operatorname{erf}(r/2R_0)}{r}, \quad (16)$$

and the force falls short of Newton’s by

$$1 - \frac{F(r)}{F_N(r)} = \operatorname{erfc} \frac{r}{2R_0} + \frac{r}{\sqrt{\pi}R_0} e^{-r^2/4R_0^2}, \quad (17)$$

57% at $r = 2R_0$, 4.6% at $4R_0$ and 4×10^{-4} at $6R_0$ (Fig. 1b). A measurement of R_0 , by the X-ray cutoff energy if it lies near the lower edge or by the separation dependence of the rate of bodies smaller than R_0 , fixes the short-range force law with no free parameter.

5.3 Entanglement tests

Let two bodies A, B each be in two branches, with smeared density differences $\Delta f_A, \Delta f_B$, self-energies $E_{A,B} = U(\Delta f_{A,B}, \Delta f_{A,B})/2$ in the positive Newton form $U(f, g) = G \iint f g / |\mathbf{x} - \mathbf{y}|$, and mutual phase rate $\hbar\omega = |U(\Delta f_A, \Delta f_B)|$.

Proposition 5.1 (Entanglement criterion). *Under Eq. (7) at coefficient C , with the decoherence*

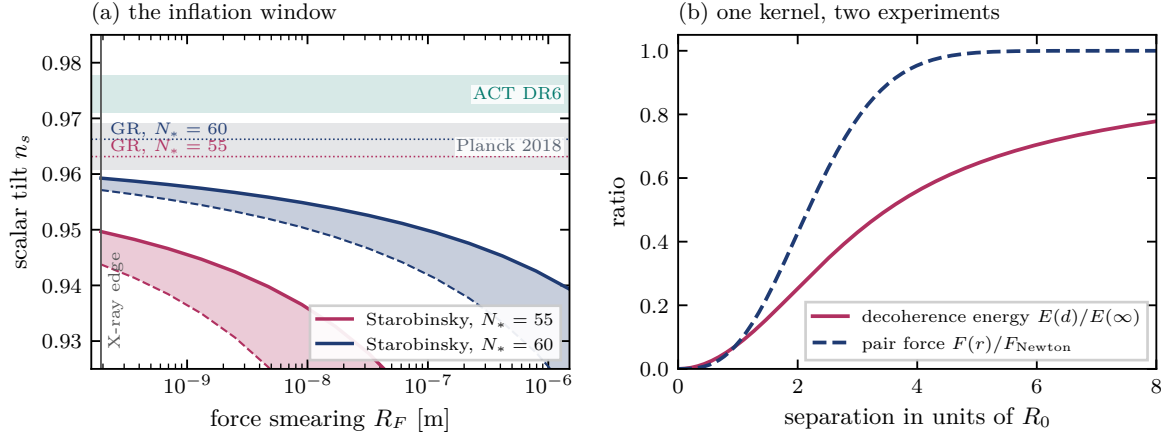


Figure 1. (a) Scalar tilt of Starobinsky inflation in the theory, from the full mode integration without metric back-reaction while the mode is inside the force smearing R_F , for $N_* = 55$ (maroon) and $N_* = 60$ (navy); each band spans the two lock-in prescriptions of Section 6.2 (solid: lock-in 2.2 e-folds early; dashed: nominal). Dotted lines: general relativity. Shaded: Planck 2018 and ACT DR6 at 1σ . The vertical line is the X-ray edge $R_F = 1.93 \text{ Å}$; the curves stop where fewer than three e-folds of inflation remain after the mode leaves R_F . (b) The one kernel of Eq. (16): decoherence energy of two point-like branches relative to its value at infinite separation, and the pair force relative to Newton's.

and force kernels equal, $|++\rangle$ becomes entangled at $t \rightarrow 0^+$ iff

$$C^2 < \frac{\eta}{1 - \eta}, \quad \eta = \frac{(\hbar\omega)^2}{4E_A E_B} \leq 1. \quad (18)$$

For uniform spheres with R_0 below their radii, touching with transverse arms and large displacement give $\eta = 25/144$, so no geometry of this kind entangles two spheres at $C \geq 5/\sqrt{119} = 0.458$.

Proof. The branch-pair decoherence rates are $CE_A/\hbar$ and $CE_B/\hbar$ for single flips and $C(E_A + E_B \mp U(\Delta f_A, \Delta f_B))/\hbar$ for double flips, and the non-local energy is $\frac{1}{4}\hbar\omega s_A s_B$ with $s = \pm 1$ labelling the branches. $|++\rangle$ is invariant under partial transposition, so at first order in t the partially transposed state can turn negative only on the three-dimensional complement of $|++\rangle$, where its generator has the eigenvalues 0 and $\frac{1}{4}[g_A + g_B \mp \sqrt{(g_A - g_B)^2 + (1 + C^2)\omega^2}]$, with $g_{A,B} = CE_{A,B}/\hbar$. The lower one is negative iff $4g_A g_B < (1 + C^2)\omega^2$, which is Eq. (18); $\eta \leq 1$ is Cauchy–Schwarz for the positive form U . For touching spheres of radius a with transverse arms of length $\Delta \rightarrow \infty$, $U(\Delta f_A, \Delta f_B) = 2Gm^2/(2a)$ and $E_A = E_B = \frac{6}{5}Gm^2/a$, so $\eta = (Gm^2/a)^2/(4 \cdot \frac{36}{25}(Gm^2/a)^2) = 25/144$. $\square$

At $C = 1$ this is the criterion of Trillo and Navascués [30], and it has the structure of the two-state criterion of Fabiano et al. [55] (Section 3). It bears on the proposals of Bose et al. and of Marletto and Vedral to witness gravitationally induced entanglement [60, 61]. A scan of two-branch geometries with unequal radii up to 10:1 does not exceed 25/144; numerically the condition at $t \rightarrow 0^+$ is also the condition at any later time. At $C = 1$ every two-sphere entanglement test with R_0 below the body sizes is therefore null, whatever carries the interaction, and the bodies decohere before any witness could form. The transverse-traceless sector cannot rescue it: it couples slow bodies through $\rho v_i v_j$ or support stresses, a coherent cross term of relative size $(v/c)^4 = 5 \times 10^{-49}$ for 250 μm in 1 s or $(gR/c^2)^2 \sim 10^{-44}$ for micron-sized bodies, while entanglement at $C = 1$ needs at least 0.49. For R_0 of order the separation the smeared feedback does entangle [30], which needs R_0 near the inverse-square-law edge. A positive witness

with R_0 below the body sizes would exclude the theory, as it would any classical constraint sector; it would not, by itself, show that gravitons are quantum [14].

5.4 Moving bodies and stress

The record monitors T^{00} in the frame of the leaves. A body moving with velocity $\mathbf{v}$ has $T^{00} = \gamma^2 \rho_0$, with ρ_0 its rest-frame density evaluated at contracted coordinates. For point-like branches the substitution $k_{\parallel} = \gamma k'_{\parallel}$ in Eq. (1) gives, in leaf time,

$$\Gamma_{\text{leaf}} = \gamma^2 \Gamma_{\text{rest}}(d_{\text{leaf}}), \quad (19)$$

with d_{leaf} the separation measured in the leaves. A Lorentz-covariant decoherence functional would give $\Gamma_{\text{rest}}/\gamma$, time dilation alone. For separations transverse to $\mathbf{v}$ the ratio is γ^3 ; at the speed of the cosmic-microwave-background dipole, 369.8 km s^{-1} , $\gamma^3 - 1 = 2.3 \times 10^{-6}$, with a sidereal anisotropy of the same order. Along $\mathbf{v}$ the contraction adds a dependence on d/γ that vanishes for $d \gg R_0$. The effect is 10^4 below a 2% measurement of C . It assumes that the leaves at our position move with the microwave background; constant-mean-curvature slices are uniform-expansion slices, which coincide with comoving ones only up to $\mathcal{O}(k^2/a^2 H^2)$, and the local velocity of their normal is not computed here.

Propositions 3.3 and 3.4 give two more channels: superpositions differing in transverse stress or in transverse momentum density decohere at least eight or four times faster than mass superpositions with the same source amplitude. Laboratory stresses and currents are too small for these to be tested; the stress channel enters cosmology through the heating of radiation.

6 Cosmology

6.1 Background and tensors

On constant-mean-curvature slices of a spatially flat Friedmann–Robertson–Walker universe the linearized Hamiltonian constraint is exactly Poisson’s equation. With $\delta K = 0$ and no background shear, ${}^3R + \frac{2}{3}K^2 - \sigma_{ij}\sigma^{ij} = 16\pi G\rho$ gives $\delta {}^3R = 16\pi G\delta\rho$, and with $\delta {}^3R = 4\nabla^2\psi/a^2$,

$$\frac{k^2}{a^2}\psi = -4\pi G\delta\rho \quad \text{at every } k, \text{ super-horizon included.} \quad (20)$$

The laboratory kernel and the cosmological constraint are one equation. The slices are uniform-expansion slices; they coincide with comoving slices only up to $\mathcal{O}(k^2/a^2 H^2)$. On them $\delta\rho \propto k^2\psi$ on super-horizon scales: the $p = 2$ class, which escapes Martin and Vennin’s exclusion of $p < 2$ [18]; the effect of collapse noise on inflationary perturbations is analysed in Ref. [59].

The expansion rate is classical. The leaves are labelled by $K = -3H$, York’s time, so the label of the foliation is the expansion rate. The Friedmann equation is the $\mathbf{k} \rightarrow 0$ constraint, sourced by the total energy, which commutes with the Hamiltonian: a record without back-action on what it records. Its noise in the potential, $S_{\Phi} = \hbar G 4\pi e^{-k^2 R_0^2}/k^2$, is a density noise $\propto k^2$, and on a Hubble patch over a Hubble time

$$\left(\frac{\delta H}{H}\right)^2 \simeq \frac{(\ell_P H/c)^2}{12\sqrt{\pi}} \left(\frac{c/H}{s}\right)^5, \quad s^2 = (c/H)^2 + R_0^2, \quad (21)$$

below 10^{-100} during inflation and below 10^{-122} today.

Tensors carry no classical stochastic component. The radiative sector is quantum linearized gravity on the classical background, so the tensor spectrum is the Bunch–Davies one, $P_t = 2H_*^2/(\pi^2 M_P^2)$, with $n_T = -2\epsilon_*$. The record noise reaches h^{TT} only through the tidal drive of Eq. (13), whose source vanishes on a homogeneous background: there is no classical stochastic tensor background at linear order, unlike in a theory with the whole metric classical. The tensors are not exempt from decoherence. While a mode is inside R_F the record does not see it; afterwards the record holds its energy density, quadratic in h^{TT} , and decoheres superpositions of different energy distributions as in Section 4.4. That rate for squeezed super-horizon tensors is not computed here.

Big-bang nucleosynthesis and the microwave-background spectrum are untouched. The heating per unit energy is $\kappa \Gamma_h$ with $\Gamma_h = P/(Mc^2) = 1.5 \times 10^{-33} \text{ s}^{-1}$ at the lower edge, and the injected fraction by recombination is below 10^{-18} even for $\kappa = 4.5$, an estimate of what the stress channel of Proposition 3.3 adds for radiation; the FIRAS bounds are $|\mu| < 9 \times 10^{-5}$ and $|y| < 1.5 \times 10^{-5}$.

6.2 The inflation window

The force smearing satisfies $R_F \geq R_D \geq 1.93 \text{ \AA}$ (Corollary 3.2 and the window). During inflation $c/H_* \ll R_F$, so every mode leaves the horizon deep inside the smearing and remains there for $N_R = \ln(H_* R_F/c)$ e -folds, 44.2 at $R_F = 1.93 \text{ \AA}$ for Starobinsky inflation [31] at $N_* = 55$. In that window the mode is super-Hubble but the classical sector does not respond to it, $e^{-k^2 R_F^2/a^2} \simeq 0$; the inflaton perturbation is a test field in an unperturbed background.

Proposition 6.1 (The window). *Let N count e -folds before the end of inflation, $N_x = N_* - N_R$, and $F = \exp(-2 \int_{N_x}^{N_*} \epsilon dN)$. In slow roll, inside the window*

$$\frac{d \ln \zeta}{dN_{\text{fwd}}} = -2\epsilon, \quad P_\zeta \rightarrow F^2 P_\zeta, \quad (22)$$

and after the mode leaves R_F it is conserved. The spectrum is shifted by

$$\Delta n_s = 4[\epsilon_* - \epsilon(N_x)(1 - \epsilon_*)] \leq 0, \quad n_T = -2\epsilon_*, \quad r = \frac{16\epsilon_*}{F^2} = -\frac{8n_T}{F^2}. \quad (23)$$

Proof. For $u = \dot{\phi}$ the background equation gives $\ddot{u} + 3H\dot{u} + V''u = -3\dot{H}u$. Without metric back-reaction the long-wavelength perturbation obeys $\delta\ddot{\phi} + 3H\delta\dot{\phi} + V''\delta\phi = 0$. With $\delta\phi = uT$, $u\ddot{T} + (2\dot{u} + 3Hu)\dot{T} - 3\dot{H}uT = 0$, and in slow roll $\dot{T} = (\dot{H}/H)T$, so $d \ln T/dN_{\text{fwd}} = -\epsilon$. Since $\zeta = -H \delta\phi/\dot{\phi} = -HT$ and $d \ln H/dN_{\text{fwd}} = -\epsilon$, $d \ln \zeta/dN_{\text{fwd}} = -2\epsilon$. Integrating over the window gives F . With $\ln k = -N_*$ and $dN_x/dN_* = 1 - dN_R/dN_* = 1 - \epsilon_*$, $d \ln F^2/d \ln k$ gives Eq. (23). After the window the recorded density sources Eq. (20): frozen super-horizon modes commute with themselves at different times up to $(k/aH)^3$, so they are recorded without back-action, and fluid densities are number-conserving, so both are in the spacelike record; ζ is then conserved by local energy conservation alone [32]. The tensor spectrum is unaffected. $\square$

The full mode integration without back-reaction inside R_F agrees with Eq. (23) to 7%. Where the classical response switches on is not unique: a flat-gauge formulation locks in at $e^{-k^2 R_F^2/a^2} \sim 1$, a comoving-density formulation 2.2 e -folds earlier, where its Poisson coefficient passes through

Table 1. *Starobinsky inflation at $R_F = 1.93 \text{ \AA}$. Two values: lock-in nominal / 2.2 e-folds early. Planck 2018: $n_s = 0.9649 \pm 0.0042$ [33]; ACT DR6 with Planck, lensing and DESI BAO (P-ACT-LB): $n_s = 0.9743 \pm 0.0034$ [34].*

N_*	N_x	Δn_s	n_s (GR: 0.963, 0.966)	from Planck	from ACT DR6
55	10.8	-0.0194 / -0.0135	0.9438 / 0.9497	5.0σ / 3.6σ	9.0σ / 7.2σ
60	15.9	-0.0091 / -0.0070	0.9571 / 0.9593	1.9σ / 1.3σ	5.1σ / 4.4σ

zero. Continuity of ζ through lock-in is assumed; a well-posed linear theory of the hybrid on a time-dependent smearing is not available.

Table 1 and Fig. 1a give the numbers. The lower edge is the most favourable point of the window: larger R_F moves N_x closer to the end of inflation, where ϵ is larger. The theory does not exclude inflation; it selects. Models whose ϵ stays flat until the end, with $\epsilon(N_x) - \epsilon_* \lesssim 2 \times 10^{-3}$, keep their tilt; Starobinsky and plateau models in their usual form are pushed 3.6σ or more from Planck and 7.2σ or more from ACT DR6 at $N_* = 55$, and at $N_* = 60$ remain 4.4σ or more from ACT DR6. If the mode is still inside R_F at the end of inflation, which for Starobinsky needs $R_F \gtrsim 3 \times 10^{-6} \text{ m}$ at $N_* = 55$ and $2 \times 10^{-8} \text{ m}$ at $N_* = 50$, the evolution without gravity through the end of inflation gives $F = 0.40$, independent of k , and a perturbative reheating inside the window multiplies ζ by a further $3/4$: $\delta\rho/\rho$ stays constant for both fluids while $\zeta = \delta\rho/(3\rho_\phi + 4\rho_r)$ goes from a third to a quarter of it. The amplitude is absorbed in the potential, so $r = 16\epsilon_*/F^2 \approx 0.04$ at $N_* = 55$, at the BICEP/Keck bound $r < 0.036$ [35].

A comoving or conformally invariant smearing, natural in shape dynamics, would remove the window, at the price of an R_0 that is not a fixed length: its laboratory value would differ from its inflationary one by $e^{N_{\text{tot}}}$.

6.3 The cosmological constant

The equilibrium axiom gives the trace-free Einstein equation, in which Λ is an integration constant; with the Bianchi identity, $\partial_\nu \Lambda = 8\pi G J_\nu$, where J_ν is the non-conservation of the matter stress tensor. The record heats matter and radiation, $J_0 < 0$, so Λ only decreases:

$$\dot{\rho}_\Lambda = -\Gamma_h(\rho_m + \kappa\rho_r)c^2, \quad w+1 = \frac{\Gamma_h(\rho_m + \kappa\rho_r)}{3H\rho_\Lambda}. \quad (24)$$

Today $w+1 = 1.1 \times 10^{-16}$ at the lower edge, positive at every redshift and rising to the past. The theory predicts no phantom crossing and the opposite sign of w_a to the DESI DR2 preference for $w_0 > -1$, $w_a < 0$, which reaches 2.8 – 4.2σ depending on the supernova sample [36]. If that preference holds, it is not this heating, and dark energy would need a separate sector.

The integrated drift is large. In the radiation era $\rho_r = 3/(32\pi G t^2)$, and the drift accumulated after t_1 is

$$\frac{\Delta\rho_\Lambda(t > t_1)}{\rho_\Lambda} = \frac{3\kappa\Gamma_h}{32\pi G t_1 \rho_\Lambda}, \quad (25)$$

dominated by the earliest heated epoch. At $t \geq 1 \text{ s}$ the radiation is hard on the scale of the record, $pR_0/\hbar \approx 2 \times 10^3$ at 1 s , so $\kappa = 4/3$ and the drift from $t_1 = 1 \text{ s}$ is $158\rho_\Lambda$ at the lower edge, or about $530\rho_\Lambda$ if the stress channel raises κ to 4.5 . Earlier epochs add more, as $1/t_1$, but the flat-space heating rate holds only while $c/H \gg R_0$. Extrapolating it down to $c/H = R_0$, at $t_1 = R_0/(2c) = 3 \times 10^{-19} \text{ s}$, gives about $5 \times 10^{20}\rho_\Lambda$; this is an extrapolation, and the epoch

$c/H < R_0$ is not treated. The initial value Λ_0 must cancel the drift to that precision. This adds no free parameter, since Λ was an integration constant, but it closes two doors: $\Lambda_0 = 0$ gives a negative effective cosmological constant and recollapse at scale factor below 0.2 even with the drift from $t_1 = 1$ s alone; and the accumulated non-conservation of a collapse dynamics, proposed as the origin of dark energy in unimodular gravity [37], cannot produce a positive Λ from $\Lambda_0 = 0$ here, since the heating only lowers Λ . Dark energy is not explained by the theory.

7 Costs

A preferred frame, and signalling suppressed by G . The record of a non-local f responds at the level of the ensemble, instantaneously in leaf time and at every distance: an intervention by Alice changes $\langle f \rangle$ at Bob’s position through the spectral projection, with Compton-suppressed tails for massive fields and power-law tails for massless ones, and the classical field passes the change on at order G . The foliation forbids closed causal loops but not this signal. Its size is not computed here. A covariant version would be worse: if the non-local response existed in every frame, relaying it through a frame moving at v would return it at $t = -vL < 0$, a closed signalling loop. A preferred-frame signal suppressed by G is the price of a dark vacuum with a non-local record.

An assumed foliation. The leaves are postulated, and so is their constant mean curvature. The cone theorem shows only that a causal, non-radiating, Newtonian classical field needs them; it does not produce them, and constant-mean-curvature Cauchy surfaces do not exist in every cosmological spacetime [62].

The Born rule is assumed. The probability of a record history is the norm of the unnormalized conditional state, which is the Born rule for continuous measurement. Given it, outcome frequencies are $|c_{\pm}|^2$. The theory does not derive it: with the ostensible record measure, every superposition ends 50:50, and the final ensemble depends on how a mixed state was decomposed, which with an entangled partner is a signal. Linearity over mixtures selects the Born rule, as in Gisin’s argument [38]; the theory assumes it.

The record on curved backgrounds. Definition 2.1 uses the four-momentum of a flat-space quantum sector. In Friedmann–Robertson–Walker spacetime it is defined on transition frequencies, keeping the components with $\omega_j^2 < k_{\text{phys}}^2$ without Fourier-analysing their slowly varying envelopes; fluid densities and frozen super-horizon modes are then in the record, and the cone boundary is crossed only at horizon exit, deep inside the smearing whenever $c/H < R_0$. At the light ring of a black hole the tidal transitions sit at the cone, $3M\omega_{220} = 1.12$, and the projection has no definition there. The theory’s mean is general relativity’s by construction, but its record during merger and ringdown is not defined. The dark vacuum of Proposition 2.2 is established for a flat background only.

Burst timing. The record noise acts on the energy density of every species. Quanta with $\hbar c/E \ll R_0$ are in the geometric-optics regime, $\dot{\boldsymbol{\theta}} = -c\nabla_{\perp}\phi_n$, and their direction diffuses at

$$\frac{d\langle\theta^2\rangle}{dl} = \frac{\ell_P^2}{3\sqrt{\pi}R_0^3}, \quad (26)$$

independent of energy above $\hbar c/R_0$, $6.9 \times 10^{-42} \text{ m}^{-1}$ at the lower edge. For paths that reach the observer (a Brownian bridge in angle) the excess delay has a spread $0.0527 DL^2/c$. For a burst at $z \approx 0.9$ (comoving distance in flat Λ CDM, cosmological factors of order one omitted) this is 11 s at $R_0 = 1.93 \text{ \AA}$, and 2 ms for GRB 170817A at 40 Mpc. The short GRB 090510 delivered a 31 GeV photon within its first second [39]; a spread below 1 s needs $R_0 > 4.3 \text{ \AA}$, and millisecond structure would need $R_0 > 2.0 \text{ nm}$. This is an estimate, not a bound: the conversion of a light-curve width into a limit on R_0 needs the burst’s intrinsic variability, its redshift and the cosmological factors. If it holds, it removes the lower part of the window, 1.9–4.3 $\text{\AA}$, which is also the part an X-ray detector can reach.

What else is open. The equal-time Hamiltonian form of the quantum timelike constraint exchange, its cross term with the classical sector, and its complete positivity; gauge invariance of the split beyond linear order, where classical and operator diffeomorphisms mix; continuity of ζ through lock-in (Section 6.2); the spin and curvature-coupling factors of the residual heating; and the velocity of the leaves at our position.

8 The theory in three axioms

The rule of Definition 2.3 compresses into three statements.

1. **Record.** On a foliation of spacetime, part of geometry is classical: a Markovian, diffusive record of the spacelike part of the constraint source at resolution R_0 (Gaussian profile, open cone), fed back linearly through a state-independent kernel as the gravitational field. Matter and the rest of geometry, helicity two, are quantum.
2. **Equilibrium.** Every small causal diamond is in entanglement equilibrium in the mean over the record, in every frame [40].
3. **Least decoherence.** The averaged decoherence kernel is least on the energy block of the foliation, in the sense of Section 3.1.

Axiom 2 gives the trace-free Einstein equation in the mean, with Λ an integration constant, and fixes the record’s gain to G ; its clause “in every frame” is the covariant mean field, $\lambda = 1$, carried and not derived. Axiom 3 is Tilloy and Diósi’s principle [8] in the form of Theorem 3.1. R_0 is the only free numerical constant. It is not the only free choice: Axiom 1 assumes the foliation and its constant mean curvature, the spacelike projection with the open cone, the Markovian, diffusive, linear and state-independent class in which the floor is proved (which excludes jump records by postulate), and the Gaussian profile; Axiom 3 assumes the energy block; and the background assumes the Born rule, local quantum field theory with the spectrum condition, and an adiabatic vacuum in expanding spacetimes. At linear order the classical set is York’s constrained data and the quantum set is the shape of shape dynamics [17]; the theory is that

dictionary with the Lichnerowicz–York equation sourced by the record instead of by operators, and the dictionary motivates the leaves without deriving them.

9 Discussion

9.1 Relation to earlier work

Penrose’s argument concerns stationary superpositions [3], whose fields have only spacelike support; it is about exactly the sector the theory makes classical, and it leaves the radiative sector quantum. In the theory nothing superposes geometries: the constraint geometry is single-valued by construction, and complete positivity then forces matter to decohere at no less than $E_G/\hbar$ (Theorem 3.1). The master equation is Diósi’s [1, 2], with the smearing length as the mass-density resolution [57], the record and feedback are Tilloy and Diósi’s [7], after Kafri, Taylor and Milburn [6], and $C = 1$ is their least decoherence [8]. The floor is the decoherence–diffusion trade-off of Oppenheim et al. [10, 11], and the minimal-noise bounds of Refs. [56, 55] are its model-independent counterpart for entanglement (Section 3); the difference from their classical–quantum gravity [9, 53] is that here only the constraint sector is classical, so gravitational waves are quantum, do not diffuse, and do not decohere stresses through a classical radiative field. The dark vacuum of the free-field projection is Pearle’s [13], and the observation that spacelike noise excites the vacuum only beyond lowest order is his [12]; what is added is the projection on the exact interacting spectrum, which gives the X-ray cutoff, its use for the full constraint source, and the gravitational feedback. Relativistic collapse dynamics has a long history: Ghirardi, Grassi and Pearle found that it excites the vacuum [43], Diósi formulated relativistic continuous measurement [44], Tumulka built a relativistic flash model [45], Bedingham a covariant model with a pointer field [46], and Pearle a relativistic model that evolves on a foliation [47]. Myrvold showed that a Lorentz-invariant Markovian collapse with a stable vacuum needs non-standard degrees of freedom [41]; the theory keeps the vacuum stable by giving up locality and Lorentz invariance of the record, and its classical field is the non-standard degree of freedom. Piccione and Bassi built a Poissonian hybrid with a stable vacuum [42]; jump records lie outside the class of Theorem 3.1. Anastopoulos and Hu argued that Newtonian entanglement is agnostic about the nature of gravity’s true degrees of freedom [14]; the theory is a concrete case with a classical constraint sector and a quantum radiative one, and in it two-sphere entanglement tests are null (Proposition 5.1). Danielson, Satishchandran and Wald argued that Newtonian entanglement implies graviton entanglement [15]; their argument starts from Newtonian entanglement, which the theory does not produce, and its radiative sector is quantum in any case. York’s split [16] and the shape-dynamics dual [17] supply the kinematics. Martin and Vennin’s $p = 2$ condition [18] is met on the leaves by Eq. (20); the inflation window of Section 6.2 is a separate effect of the force smearing.

9.2 What would decide

The theory makes R_0 the one number that links the laboratory to the sky, and several measurements can each exclude it:

1. A single-body interferometer in the saturated regime measuring $C < 1$ (Theorem 3.1), or $C \neq 1$ if least decoherence is kept.

2. A Diósi–Penrose-strength spontaneous X-ray excess without the Gaussian cutoff of Eq. (14); or, if a cut excess is found, a cutoff energy $\hbar c/R_0$ inconsistent with the short-range force law of Eq. (17).
3. Gravitationally induced entanglement of two bodies with R_0 below their sizes.
4. A measured scalar tilt that excludes every inflation model with $\epsilon(N_x) - \epsilon_* \lesssim 2 \times 10^{-3}$, or a tensor ratio inconsistent with $r = -8n_T/F^2$.
5. A classical stochastic tensor background, or diffusion of gravitational-wave strain above the tidal-drive level of Section 4.4.
6. Millisecond structure in the arrival times of hard photons from cosmological bursts, if the estimate of Section 7 survives a full treatment, for R_0 below 2 nm.

10 Conclusions

A theory in which the classical part of geometry is the record of the spacelike part of the constraint source on constant-mean-curvature slices, and everything else is quantum, meets the conditions that a Diósi–Penrose rate imposes, given an assumed foliation. Its constraint field is instantaneous, so it does not drag; its leaves are the time direction a rate requires; its projection keeps the flat-space vacuum dark. Least decoherence fixes $C = 1$. In the mean it is general relativity, provided the record holds every constraint component, including the graviton energy and a tidal term, the timelike constraint exchange is quantum, and the mean field is covariant ($\lambda = 1$, assumed); radiation reaction through 3.5PN order needs no record of the radiation. Data confine its one length to $1.93 \text{ \AA} < R_0 < 4.0 \times 10^{-5} \text{ m}$; it predicts a Gaussian X-ray cutoff near $\hbar c/R_0$, one kernel for decoherence and the short-range force, null entanglement tests, and a γ^2 moving-body rate. Its cosmology has a classical expansion rate and no classical stochastic tensor background, but it shifts the scalar tilt of plateau inflation beyond the data at $N_* = 55$, needs a tuned initial cosmological constant, and says nothing about dark energy. It costs a preferred frame with signalling suppressed by G , an assumed Born rule and an undefined record at the light ring, and a burst-timing estimate threatens the lower part of its window. It is a candidate with tensions that experiment can resolve.

A Weinberg–Witten, collapse and ringdown

A.1 Weinberg–Witten and the slice energy

A massless particle of helicity $|h| > 1$ has no Lorentz-covariant conserved $T_{\mu\nu}$ with non-vanishing forward one-particle matrix elements [25]. In the Breit frame, $\mathbf{p}' = -\mathbf{p}$, every rank-two contraction of helicity-two polarization tensors vanishes, since each contains a factor $e_+ \cdot e_+ = 0$; a covariant tensor’s matrix element in any frame is a boost of its Breit-frame element, so it vanishes at every spacelike q , including the forward limit, where the Bondi mass needs the value ω . The record uses the York energy density on the leaf, whose one-graviton elements are $(\omega\omega' + \mathbf{p} \cdot \mathbf{p}')(2\sqrt{\omega\omega'})^{-1} \cos^4(\theta/2) \rightarrow \omega$ in the forward limit. It is not a tensor, which is how the theorem is evaded, and it needs the slicing the theory already has.

The open cone records the cycle-averaged energy of a wave. Scattering transfers between graviton modes are spacelike, null only for collinear modes; pair transfers, which carry the oscillation of the energy density at $(2\omega, 2\mathbf{k})$, are timelike or null. The record keeps the Isaacson average [26]. A finite packet's envelope has $q^2 = \Delta k_\perp^2 (1 - \Delta k_\parallel/k) > 0$ and is kept; only a strictly planar wave is dark.

A.2 Strong fields

Static configurations, neutron-star structure, tidal Love numbers and Kerr exteriors, sit at $\omega = 0$ and are classical constraint problems, general relativity's up to the noise variance above. Maximal slices of Schwarzschild end on the limit surface $r = 3M/2$ [27], so the foliation exists through collapse and never reaches the singularity; the collapsing lapse freezes the interior's record in leaf time. Ringdown is general relativity's in the mean once the tidal term is recorded: on the classical Schwarzschild background Leaver's continued fraction gives $M\omega_{220} = 0.37367 - 0.08896i$ [28]. The transitions that carry the tidal term exchange (ω, k) with the classical sector. In the near zone $\omega/(ck) \sim 2v/c < 1$, spacelike as the rule requires; at the light ring $3M\omega_{220} = 1.12$, at the cone, where the record's cone has no definition on a curved background. The theory is incomplete in the merger-ringdown regime; it is not inconsistent there, since its mean is general relativity's by construction.

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