Diósi–Penrose Noise from Relaxing Emergent Geometry: A Thermal Model with $C = 1/\pi$

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Abstract

Quantized linearized gravity does not supply a decoherence rate of first order in G for a static, equal-mass spatial superposition. The Diósi–Penrose rate instead requires a source-independent potential noise with correlation $\propto \hbar G/|\mathbf{x}-\mathbf{y}|$ and weight at zero frequency. We ask whether that noise can be the equilibrium fluctuation of an emergent, coarse-grained geometry. Summing the fixed-volume area deficits of entanglement-equilibrium balls of size ℓ, and using the canonical fluctuation formula with the lapse, gives at linear order a Gibbs weight $\exp(-\beta E_N[\Phi])$. Here E_N is the Newtonian energy functional, so the Poisson equation is its maximum, and the effective temperature is $k_B T_{\text{eff}} = 5\hbar c/(2\pi\ell)$. The equal-time kernel is then exactly Diósi's. General relativity propagates these fluctuations rather than relaxing them; relaxation must be added. We identify a candidate rate: the modular mixing of the field that carries the noise in the balls' KMS state. Consistency of a single KMS state for statics and dynamics fixes the rate coefficient to $C = 1/(\pi\Delta)$, which gives $C = 1/\pi$ for a dimension-one field. The resulting noise is thermal. Decoherence follows the Diósi–Penrose form with a Yukawa-regularized kernel and smearing length $R_0 = \ell/5$, and $k_B T_{\text{eff}} R_0 = \hbar c/(2\pi)$. Hard spontaneous emission is suppressed by a Bose factor, and bulk heating is below one percent of that of Gaussian-smeared Diósi–Penrose noise at equal R_0. Confronted with published radiation and heating data, the model survives for $R_0 \gtrsim 3\times10^{-11}$ m ($\ell \gtrsim 1.3\times10^{-10}$ m); the binding constraint is the softest X-ray search, not heating and not the hard-photon search that excludes the parameter-free Diósi–Penrose model. For a 1 μg mass at 1 mm the decoherence time is ≈ 5.0 ns, a factor π longer than the Diósi–Penrose value. The results are conditional on four stated assumptions, which we list.

1 Introduction

The Diósi–Penrose proposal [1, 2, 3] assigns a spatial superposition of a mass the decoherence rate $E_{\text{G}}/\hbar$, where $E_{\text{G}} = \frac{G}{2}\iint \Delta\rho(\mathbf{x})\Delta\rho(\mathbf{y})/|\mathbf{x}-\mathbf{y}|$ is the self-energy of the difference $\Delta\rho = \rho_L - \rho_R$ of the branch mass densities. Perturbative graviton environments give rates of second order in G [4, 5]. A companion analysis shows why [6]. For a static, inertial source with branches of equal mass, quantized linearized gravity has no zero-frequency noise that couples to $\Delta\rho$ at

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any temperature, and relational clock readouts depend on mass and time only through GMt. The first-order rate is reproduced exactly by a source-independent white potential noise with correlation $\hbar G/|\mathbf{x}-\mathbf{y}|$. That noise has weight at zero frequency and nonzero wavenumber, which no free relativistic field supplies.

This paper asks whether such noise can arise from gravity's thermodynamic side. If the Einstein equation is an equation of state [7] or an entanglement-equilibrium condition [8], the coarse-grained geometry is an equilibrium variable, and equilibrium variables fluctuate. We compute two things:

- the stiffness of those fluctuations from generalized entropy (Section 2);
- their relaxation rate, using modular flow (Section 3).

We then derive the consequences for decoherence, spontaneous radiation and heating (Section 4).

The outcome is a model rather than a derivation from first principles. The stiffness follows at linear order from three assumptions. The relaxation requires a fourth, because general relativity by itself propagates the relevant fluctuations. Section 5 lists all four. Within the model, the rate coefficient is fixed, $C = 1/\pi$, and the noise is thermal. Both features distinguish it from the standard Diósi–Penrose model.

2 Stiffness from generalized entropy

Area deficit. For a geodesic ball of fixed proper volume $V = \frac{4\pi}{3}\ell^3$ in a three-geometry with scalar curvature R at its centre, the boundary area differs from its flat value by

$$A|_V - A_{\mathrm{flat}}(V) = -\frac{4\pi\ell^4}{30}\,R + \mathcal{O}(\ell^5), \tag{1}$$

the relation underlying the entanglement-equilibrium derivation of the Einstein equation [8]. Consider a static, conformally flat fluctuation $h_{ij} = e^{2\omega}\delta_{ij}$ with $\omega \simeq -\Phi/c^2$. Its curvature is $R = -e^{-2\omega}(4\nabla^2\omega + 2|\nabla\omega|^2)$, and integration by parts gives the exact identity

$$\int \sqrt{h}\,R\,d^3x = 2\int e^{\omega}|\nabla\omega|^2\,d^3x. \tag{2}$$

Tile space with such balls, counted by proper volume. The geometric part of generalized entropy then changes by

$$\Delta S_{\mathrm{geom}} = -\frac{\ell c^3}{20 G\hbar}\int |\nabla\omega|^2\,d^3x + \mathcal{O}(\omega^3), \tag{3}$$

which is negative definite and carries $1/G$. Counting balls by coordinate volume instead reverses the sign, so proper-volume counting is an assumption with consequences.

Fluctuation weight. When geometry fluctuates at a fixed matter state, the matter entropy does not change; its redshifted energy is exchanged with the reservoir. The fluctuation weight is therefore $\exp(\Delta S_{\mathrm{geom}} - \beta\,\Delta E_{\mathrm{mat}})$, where $E_{\mathrm{mat}} = \int \mu N$, μ are fixed mass elements and N is the lapse. The static quadratic Einstein–Hilbert form in the lapse and the spatial potential is indefinite, so we restrict fluctuations to the slice on which the two potentials agree. On that slice, at linear order,

$$P[\Phi] \propto \exp\big(-\beta E_N[\Phi]\big), \qquad E_N[\Phi] = \int \Big[\frac{|\nabla\Phi|^2}{8\pi G} + \mu\,\Phi\Big] d^3x. \tag{4}$$

The maximum of (4) is the attractive Poisson equation, and its quadratic part is the stiffness $(1/8\pi G)\int|\nabla\delta\Phi|^2$. The inverse temperature in both terms is the same: the volume average of the modular inverse temperature $2\pi(\ell^2 - r^2)/(2\ell c)$ of the ball [9],

$$k_B T_{\text{eff}} = \frac{5\hbar c}{2\pi\ell}. \tag{5}$$

Adding the matter *entropy* instead of the reservoir term produces a functional whose maximum is either repulsive or has three times Newton's constant. Using the reservoir term is therefore essential.

Kernel. Equation (4) implies $\langle\delta\Phi(\mathbf{x})\delta\Phi(\mathbf{y})\rangle = G k_B T_{\text{eff}}/|\mathbf{x}-\mathbf{y}|$. This is Diósi's spatial kernel, with $k_B T_{\text{eff}}$ in place of $\hbar$ times a correlation time.

3 Relaxation

Relaxation is necessary and not supplied by general relativity. With the static statistics of Section 2, the dephasing exponent grows linearly in time only if the fluctuations relax. Propagating fluctuations with the same equal-time statistics give a saturating exponent, as in the linearized analysis [6]. A fluctuation on the slice of Section 2 violates the Hamiltonian constraint. At linear order, general relativity carries such violations through the subsystem $\partial_t\mathcal{H} = -\nabla\cdot\mathcal{M}$, $\partial_t\mathcal{M} = -\nabla\mathcal{H}$, whose modes propagate at c. Relaxation must therefore be an additional ingredient. We write it as damping at a rate κ independent of wavenumber: $\partial_t\mathcal{H} = -\nabla\cdot\mathcal{M} - \kappa\mathcal{H}$ and $\partial_t\mathcal{M} = -\nabla\mathcal{H} - \kappa\mathcal{M}$. Each mode then decays as $e^{-\kappa t}e^{\pm ickt}$, and the long-time dephasing slope is

$$\frac{F}{t} \to \frac{2k_B T_{\text{eff}}}{\hbar\kappa}\,\frac{E_G^{(\kappa)}}{\hbar}, \qquad E_G^{(\kappa)} = \frac{G}{2}\iint \Delta\rho(\mathbf{x})\Delta\rho(\mathbf{y})\,\frac{1-e^{-\kappa|\mathbf{x}-\mathbf{y}|/c}}{|\mathbf{x}-\mathbf{y}|}. \tag{6}$$

This is the Diósi–Penrose form, with the Newtonian kernel regularized at $R_0 = c/\kappa$ and coefficient $C = 2k_B T_{\text{eff}}/(\hbar\kappa)$.

A candidate rate: modular mixing. In the modular state of a Rindler wedge [10] or a ball [9], the modular flow has inverse temperature 2π. The correlation of a dimension-Δ primary along the flow decays as $(4\sinh^2(s/2))^{-\Delta} \sim e^{-\Delta s}$ in the modular parameter s. The wedge case follows from the interval between points on one boost orbit, $\sigma^2 = -4\rho^2\sinh^2(s/2)$; the ball case from the flow $t = \ell\tanh(s/2)$ at the centre with its conformal factors. The decay rate per unit modular parameter is therefore Δ.

We posit that the coarse geometry relaxes at the modular mixing rate of the field that carries the noise (assumption A4 in Section 5). That field's statistics and dynamics are two faces of *one* KMS state. For coarse, ball-uniform observables the modular generator is $\int\beta(\mathbf{x})T_{00} = \langle\beta\rangle E$, with the same average that fixed (5). In modular units the temperature is $1/(2\pi)$ and the decay rate is Δ, so

$$C = \frac{2k_B T_{\text{eff}}}{\hbar\kappa} = \frac{1}{\pi\Delta}, \tag{7}$$

independently of how modular and physical time are matched. Evaluating the temperature and the rate at different points of the same state would instead introduce a spurious factor $5/2$, and would violate detailed balance.

The field carrying the noise. The noise couples to matter through Φ. A dimension-one field in a KMS state has the equal-time correlator $(T/4\pi r)\coth(\pi r T)$, in units $\hbar = c = k_B = 1$. For $rT \gg 1$ this becomes the classical $T/(4\pi r)$ kernel of Section 2. We therefore take $\Delta_\Phi = 1$. This is motivated, not derived: in general relativity Φ is a constrained variable, not a conformal primary. With it,

$$C = \frac{1}{\pi} \approx 0.318, \qquad \kappa = \frac{5c}{\ell}, \qquad R_0 = \frac{\ell}{5}. \tag{8}$$

Since $\Delta \geq 1$ for unitary scalars in four dimensions, $C \leq 1/\pi$ for every operator in this mechanism. The Diósi–Penrose value $C = 1$ is not reached.

4 Consequences

Decoherence. Equation (6) with (8) gives $\Gamma_{\text{dec}} = C\, E_G^{(\kappa)}/\hbar$. For separated bodies larger than R_0 this reduces to $E_{\text{G}}/(\pi\hbar)$. At the cross-term benchmark of a 1 μg mass separated by 1 mm, the decoherence time is $\pi\hbar d/(GM^2) \approx 4.96$ ns, compared with 1.58 ns for $C = 1$. At times shorter than R_0/c the decay is Gaussian rather than exponential.

The noise is thermal. Combining (5) and (8),

$$k_B T_{\text{eff}}\, R_0 = \frac{C\hbar c}{2} = \frac{\hbar c}{2\pi}, \tag{9}$$

so $T_{\text{eff}} R_0 \approx 3.64 \times 10^{-4}$ K m, independently of ℓ. The noise that decoheres superpositions is therefore the fluctuation of a bath at finite temperature. It obeys detailed balance, $S(-\omega) = e^{-\hbar\omega/k_B T_{\text{eff}}} S(\omega)$. The decoherence rate uses $S(\omega \to 0)$, where the quantum and classical spectra agree, so it is unaffected. Standard white Diósi–Penrose noise corresponds instead to infinite temperature.

Spontaneous radiation. Absorbing energy $\hbar\omega$ from the bath carries the Bose factor

$$e^{-\hbar\omega/k_B T_{\text{eff}}} = e^{-2\pi\omega R_0/c}. \tag{10}$$

At $R_0 = 4 \times 10^{-10}$ m, $k_B T_{\text{eff}} \approx 79$ eV, and emission at 10 keV is suppressed by a factor below 10^{-55}. Searches for hard spontaneous radiation, which exclude the parameter-free Diósi–Penrose model [11], therefore constrain this model far more weakly; searches at soft X-ray energies, where $E \lesssim k_B T_{\text{eff}}$, do not (Section 5). Treating the same noise classically would put on-shell weight at $k = \omega/c$ and wrongly predict *more* hard emission than Gaussian-smeared white noise. That is a Rayleigh–Jeans artefact.

Heating. Low-frequency momentum diffusion is not Bose-suppressed. A free constituent of mass m gains energy at the rate

$$\frac{dE}{dt} = \frac{C\hbar G m}{2} \int \frac{d^3k}{(2\pi)^3}\, k^2 \tilde{K}(k), \tag{11}$$

where $\tilde{K}$ is the Fourier transform of the spatial kernel. The coarse-grained description has no modes below ℓ, so $\tilde{K} = 4\pi/(k^2(1 + k^2 R_0^2))$ for $k < 1/\ell$ and zero above. The heating rate is then $(C\hbar G m/\pi)(x - \arctan x)/R_0^3$ with $x = R_0/\ell = 1/5$. Relative to Gaussian-smeared white noise, $C\hbar G m/(4\sqrt{\pi}R_0^3)$, the ratio is $r_k \equiv (4/\sqrt{\pi})(x - \arctan x) \approx 5.9 \times 10^{-3}$ at the same C and nominal R_0. Heating-based lower bounds on R_0 therefore weaken by the cube root of that ratio, a factor ≈ 0.18 at equal C, or $(C r_k)^{1/3} \approx 0.12$ against bounds quoted for $C = 1$ (Section 5). Matter would eventually equilibrate at T_{eff}. For $R_0 \geq 4 \times 10^{-10}$ m this takes much longer than the age of the universe, so saturation has no practical consequence. If the kernel's Yukawa tail were instead to extend to nuclear scales, heating would exceed the Gaussian value by $\approx 2R_0/r_n$. The cutoff at ℓ is therefore essential to viability.

5 Assumptions and outlook

Table 1. *Assumptions of the model and their roles.*

	Assumption	Role
A1	Generalized entropy is additive over a proper-volume tiling of fixed-volume balls, with no modes below ℓ	Stiffness (3); the heating cutoff
A2	Fluctuations respect the spatial Einstein equations (the two Newtonian potentials agree)	Avoids the indefinite conformal mode
A3	Canonical reservoir fluctuation formula at the balls' modular temperature, with the lapse	Weight (4) and T_{eff}
A4	Coarse geometry relaxes at the modular mixing rate of the field carrying the noise, independently in successive cells	κ, hence C (7)
	$\Delta_\Phi = 1$ (motivated by the kernel shape)	$C = 1/\pi$ (8)

Table 1 collects the assumptions. A4 is where the model departs from general relativity, which propagates rather than relaxes the relevant fluctuations. The model does not derive the Diósi–Penrose rate from quantized gravity; the companion analysis shows that no such derivation exists at linear order [6]. It replaces a free rate postulate with a thermodynamic mechanism and a fixed coefficient.

Distinguishing predictions. Relative to Gaussian-smeared Diósi–Penrose noise, the model predicts:

1. a rate coefficient $C = 1/\pi$, a factor π below $C = 1$;
2. a Yukawa rather than Gaussian regularization of the kernel, with $R_0 = \ell/5$;
3. the parameter link (9);

4. exponential suppression of spontaneous emission above $k_B T_{\text{eff}}$;

5. heating reduced by $\approx 6 \times 10^{-3}$ at equal R_0.

Items (4) and (5) change which experiments constrain the model. We compare with prior work and then recompute the published bounds.

Relation to prior work. Dissipative extensions of the Diósi–Penrose model exist [13, 14], as do dissipative and coloured extensions of continuous spontaneous localization [15, 16, 17, 18]; see the reviews [12, 31]. In all of them the noise temperature is a further free parameter, fixed by fiat (typically to about 1 K); none ties it to the smearing length, and all remain white in time, so their temperature does not cut off the emission. Coloured models cut off emission with a power law [19], not exponentially. The obstruction that motivates A4—the Newtonian potential is a constraint, and graviton decoherence vanishes in the vacuum—is due to Anastopoulos and Hu [4, 20]. Newtonian-potential fluctuations at a causal-diamond temperature $\propto \hbar c/\ell$ appear in the interferometer proposals of Verlinde and Zurek [21], without a spatial kernel or matter decoherence; Kay [22] obtains an E_{G}-like exponent from matter–gravity entanglement, without a rate. Steady decoherence rates fixed by KMS consistency together with an Ohmic channel are known for Killing horizons [32, 23, 33]; those authors find no such rate for an inertial laboratory. Assumption A4 supplies the missing channel. We found no prior Yukawa regularization of the kernel, no prior $C = 1/\pi$, and no prior relation $k_B T_{\text{eff}} R_0 \sim \hbar c$. Bahrami *et al.* [13] showed that a *cold* dissipative bath forces dissociating momentum kicks; the present bath is hot ($k_B T_{\text{eff}} \sim$ keV at allowed R_0) and so weakly coupled that equilibration takes far longer than the age of the universe.

Recomputed bounds. Radiation searches quote bounds for white Gaussian-smeared noise. The Gran Sasso [11] and MAJORANA [26, 27] analyses use the Penrose normalization, whose rate is 8π times Diósi's [24]; XENONnT [25] and the heating bounds use Diósi's. Relative to a reference model with coefficient C_{ref} at the same R_0, the present model multiplies the emission at photon energy E by $(C/C_{\text{ref}})\, r_k\, F(E/k_B T_{\text{eff}})$, with $F(u) = u/(e^u - 1)$ the detailed-balance factor relative to classical white noise and $r_k = (4/\sqrt{\pi})(x - \arctan x) = 5.9 \times 10^{-3}$ the ratio of the momentum-diffusion constants. Holding each experiment's signal budget fixed and averaging over its window with the reference $1/E$ weight gives Table 2. For XENONnT the true low-energy spectrum is orbital-resolved [24]; the $1/E$ weighting is an approximation of the window average only. Heating and force-noise bounds [28, 29, 30] probe phonon and millihertz frequencies, where the bath is classical; they rescale by $(C r_k)^{1/3} = 0.12$.

Table 2. *Published lower bounds on R_0 for white Gaussian-smeared noise, and the same data confronted with the present model ($C = 1/\pi$, cutoff Yukawa kernel, thermal spectrum).*

Data	Window	Reference bound	This model
Gran Sasso γ [11]	1000–3800 keV	5.4×10^{-11} m	2.1×10^{-13} m
MAJORANA X-ray [26, 27]	19–100 keV	2.5×10^{-10} m	4.1×10^{-12} m
XENONnT X-ray [25]	1–140 keV	4.5×10^{-10} m	2.7×10^{-11} m
Cryostat heat leak [28]	THz phonons	4.6×10^{-12} m	5.7×10^{-13} m
Neptune heat [28]	—	3.7×10^{-12} m	4.6×10^{-13} m
LISA Pathfinder [29, 28]	mHz	8.2×10^{-14} m	1.0×10^{-14} m

The Bose factor protects only above $k_B T_{\text{eff}} = \hbar c/(2\pi R_0)$, so the constraint migrates to the softest window. The hard-photon search that excludes parameter-free Diósi–Penrose noise weakens by more than two orders of magnitude in R_0, and heating is subdominant; the binding bound is the soft-X-ray one, $R_0 \gtrsim 2.7 \times 10^{-11}$ m, hence $\ell \gtrsim 1.3 \times 10^{-10}$ m, a factor 17 below the Gaussian bound, with a shape-robust floor of 4×10^{-12} m from MAJORANA. No upper bound follows from these data. Inside the window the model predicts an emission spectrum $\propto F(E/k_B T_{\text{eff}})/E$ with $k_B T_{\text{eff}} \leq 1.2$ keV, tied to the decoherence rate by (9). A sub-keV analysis of existing xenon or germanium data with this template is the decisive test; hard-photon searches and low-frequency force-noise measurements are not.

Limitations. The analysis is linear in G and Newtonian. It assumes static, inertial sources and a single coarse-graining scale. The coarse-graining scale ℓ is not fixed by the model; only the ratio R_0/ℓ is. The four assumptions are not derived.

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