

Double-Pulsar Drag Constrains Finite-Temperature Baths as the Source of Diósi–Penrose Decoherence

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Abstract

A gravitational decoherence rate of first order in Newton’s constant, the Diósi–Penrose rate $\Gamma = CE_G/\hbar$, cannot come from quantized linearized gravity for an equal-mass spatial superposition, nor, in a Gaussian treatment, from any translation-invariant state that obeys the spectrum condition. The natural remaining source is a stationary bath at finite temperature. We show that such a bath, coupled Hamiltonianly to the mass density, drags moving bodies: if its noise has the Diósi–Penrose form and is white and passive at the wavelengths a neutron star resolves, the friction rate is at least $(2\pi/3)(C\hbar G/k_B T_{\text{eff}})$ times the density, with T_{eff} the fluctuation–dissipation temperature. The orbital decay of the double pulsar PSR J0737–3039A/B then requires $T_{\text{eff}}/C > 3.3 \times 10^{16}$ K. For dissipative Diósi–Penrose models this raises the lower bound on the collapse temperature by a factor 3×10^{29} and excludes the proposed point $T_\beta \approx 3$ K by 10^{16} . What survives is Diósi–Penrose noise in the infinite-temperature limit, that is, noise without back-action, together with baths that behave differently at kilometre scales than in the laboratory. With Gaussian smearing R_0 of the mass density, data allow $R_0 > 4.5 \text{ \AA } C^{1/3}$ with no upper bound, and C is a free parameter of order one. For a uniform sphere of radius $a \gg R_0$ separated by $d \geq 2a$ the rate is $\Gamma = CGM^2[6/(5a) - 1/d]/\hbar$, independent of R_0 , so once other decoherence channels are subtracted one visibility measurement determines C ; for $1 \mu\text{g}$ of silica $\tau = 65 \text{ ps}/C$.

1 Introduction

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

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

with C a number of order one ($C = 1$ in Diósi’s normalization). The rate is first order in G and does not depend on any environment. Reviews of gravitational decoherence [4, 5] survey the models that produce it and the experiments that test it.

Where can such a rate come from? Treating the graviton field as an environment gives master equations whose rates are also first order in G , but they decohere superpositions of

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different *energies* [6, 7, 8]. For a spatial superposition of one body the two branches carry the same mass, the monopole of the source difference vanishes, and the field channel gives no rate (Section 4). The same holds, in a Gaussian treatment, for any translation-invariant state that obeys the spectrum condition. What remains is a stationary medium that breaks the spectrum condition: a bath at finite temperature. Black-hole horizons decohere superpositions through finite-temperature quanta, and the effect is reproduced by modelling the hole, or ordinary matter, as a system at finite temperature [9, 10, 11]; those quanta are free and on shell and lack Diósi–Penrose-form noise, so they fall outside the class bounded here. Dissipative Diósi–Penrose models [12, 13] postulate a bath with that noise explicitly.

This paper shows that, under conditions stated below, a finite-temperature bath cannot be the source. If the mass density couples to the bath through a Hamiltonian, the correlator that carries the noise also carries a dissipative response, and a body moving through the bath feels friction. Theorem 3.1 fixes that friction from two numbers, the noise strength C and the fluctuation–dissipation temperature T_{eff} . Binary pulsars measure friction very well. Provided the bath’s noise keeps the Diósi–Penrose form, is white over the drift band and is passive at the wavelengths a neutron star resolves, the orbital decay of the double pulsar PSR J0737–3039A/B requires

$$T_{\text{eff}}/C > 3.3 \times 10^{16} \text{ K} \quad (2)$$

(Corollary 3.3). For dissipative Diósi–Penrose models this raises the lower bound on the collapse temperature by a factor 3×10^{29} .

What survives is Diósi–Penrose noise in the infinite-temperature limit, that is, noise without back-action [2, 14], together with baths whose kilometre-scale behaviour differs from their laboratory-scale behaviour (Section 3.4). Sections 5 and 6 give the allowed parameter region of the survivor and a prediction that measures its coefficient C independently of its smearing length.

Section 2 fixes conventions. Section 3 contains the drag theorem, the pulsar bound and its consequences. Section 4 states precisely, as known results, why the rate cannot come from quantized linearized gravity or from a translation-invariant positive-energy state. Section 7 lists the loopholes that remain open.

2 Setting, the rate from noise, and the scaling test

The measured quantity. A rigid body of mass M is prepared in a superposition of two static positions a distance d apart. Both branches carry the same mass. The source is inertial. The two paths are recombined and the fringe visibility $|z(t)|$, with $z = 2\rho_{LR}$, is recorded as a function of the time t spent apart. A rate is a constant Γ with $|z(t)| \propto e^{-\Gamma t}$ at late times.

The rate from noise. Let the mass density couple to a potential operator, $H_{\text{int}} = \int \rho \Phi d^3x$. The branches differ through the which-path variable $B(t) = \int \Delta\rho(\mathbf{x}) \Phi(\mathbf{x}, t) d^3x$. For Gaussian fluctuations and pure dephasing,

$$-\ln |z(t)| = \frac{1}{2\pi\hbar^2} \int_{-\infty}^{\infty} d\omega S_B(\omega) \frac{1 - \cos \omega t}{\omega^2}, \quad S_B(\omega) = \frac{1}{2} \int d\tau e^{i\omega\tau} \langle \{\delta B(\tau), \delta B(0)\} \rangle. \quad (3)$$

Since $\int (1 - \cos \omega t) \omega^{-2} d\omega = \pi t$, the late-time slope is

$$\Gamma = \lim_{t \rightarrow \infty} \frac{-\ln |z(t)|}{t} = \frac{S_B(0)}{2\hbar^2}. \quad (4)$$

A rate requires a zero-frequency floor. If $S_B(\omega) = \mathcal{O}(\omega^2)$ as $\omega \rightarrow 0$, the integrand of Eq. (3) is bounded at small ω and $-\ln |z|$ saturates.

Diósi–Penrose noise. Classical white potential noise with

$$\langle \delta\Phi(\mathbf{x}, t) \delta\Phi(\mathbf{y}, s) \rangle = \frac{C\hbar G}{|\mathbf{x} - \mathbf{y}|} \delta(t - s), \quad S_\Phi(\mathbf{k}, \omega) = \frac{4\pi C\hbar G}{k^2} \text{ for all } \omega, \quad (5)$$

gives $S_B(0) = C\hbar G \iint \Delta\rho \Delta\rho |\mathbf{x} - \mathbf{y}|^{-1} = 2C\hbar E_G$, so Eq. (4) reproduces Eq. (1). At $C = 1$ this is the noise behind Diósi’s master equation $\dot{\hat{\rho}} = -(G/2\hbar) \iint [\hat{M}(\mathbf{x}), [\hat{M}(\mathbf{y}), \hat{\rho}]] |\mathbf{x} - \mathbf{y}|^{-1}$ [1, 2]; the normalization used with Penrose’s prescription in Refs. [15, 16] is larger by a factor 8π , as stated in Ref. [17], so it corresponds to $C = 8\pi$. We measure C in Diósi units throughout.

The energy. E_G in Eq. (1) contains the same-branch terms; they make $E_G \geq 0$. For uniform spheres of radius a ,

$$E_G = GM^2 \left(\frac{6}{5a} - \frac{1}{d} \right) \quad (d \geq 2a), \quad (6)$$

which saturates at $(6/5)GM^2/a$ for $d \gg a$ (Appendix A). The point-mass expression GM^2/d is the magnitude of the cross term alone; it is not a limit of Eq. (1) for any mass distribution.

The scaling test. At fixed geometry, write $F(M, t) = -\ln |z|$ and define

$$R_M = \frac{F(2M, t)}{F(M, t)}, \quad R_t = \frac{F(M, 2t)}{F(M, t)}. \quad (7)$$

A rate of the form (1) has $R_M = 4$ and $R_t = 2$: the mass enters once through each coupling in the variance, G once through the noise strength, and the white spectrum makes the exponent linear in t . We apply this test below.

3 The drag theorem and the double-pulsar bound

3.1 The drag theorem

Let $H_{\text{int}} = \int \rho \Phi d^3x$ with Φ a bath operator. Per unit volume define the symmetrized noise and the dissipative response

$$S(\mathbf{k}, \omega) = \frac{1}{2} \int dt e^{i\omega t} \langle \{ \Phi_{\mathbf{k}}(t), \Phi_{-\mathbf{k}} \} \rangle, \quad \chi''(\mathbf{k}, \omega) = \frac{1}{2\hbar} \int dt e^{i\omega t} \langle [\Phi_{\mathbf{k}}(t), \Phi_{-\mathbf{k}}] \rangle. \quad (8)$$

Fix a wavenumber k_* and a speed v_{max} , and assume:

- (H) *Reciprocity*: the operator ρ couples to is the one that carries the noise, through a Hamiltonian (or Schwinger–Keldysh) coupling, so S and χ'' come from the same correlator.
- (S) *Stationarity*: the bath state is stationary, homogeneous and isotropic in a rest frame, at scales $\geq 1/k_*$.

- (D) *Diósi–Penrose noise*: $S(k, 0) = 4\pi C\hbar G/k^2$ for $k \leq k_*$.
- (W) *Whiteness over the drift band*: $\chi''(k, \omega)/\omega$ equals its $\omega \rightarrow 0$ value for $|\omega| \leq k_* v_{\max}$, $k \leq k_*$.
- (F) *Finite positive temperature*: $T_{\text{eff}}(k) \equiv S(k, 0)/(2k_B \partial_\omega \chi''(k, 0)) \in (0, \infty)$ for $k \leq k_*$; a KMS state at temperature T has $T_{\text{eff}} \equiv T$.
- (P) *Passivity*: $\omega \chi''(k, \omega) \geq 0$ at every k and ω . This holds in every KMS state at $T > 0$; for other states it is a separate assumption.
- (L) *Linear response*: the bath responds linearly to the body’s density, so the friction is linear in the velocity for all $|\mathbf{v}| \leq v_{\max}$.

Theorem 3.1 (Drag). *Under (H), (S), (D), (W), (F), (P) and (L), a body moving with velocity $\mathbf{v}$, $|\mathbf{v}| \leq v_{\max}$, relative to the rest frame feels $\mathbf{F} = -M\gamma\mathbf{v}$ with*

$$\gamma \geq \frac{2\pi}{3} \frac{C\hbar G}{k_B T_{\text{eff}}} \rho_{k_*}, \quad \rho_{k_*} \equiv \frac{1}{M} \int_{k < k_*} \frac{d^3 k}{(2\pi)^3} |\rho_{\mathbf{k}}|^2, \quad (9)$$

where T_{eff} is the largest value of $T_{\text{eff}}(k)$ for $k < k_*$.

Proof. Uniform motion gives $\rho_{\mathbf{k}}(t) = \rho_{\mathbf{k}} e^{-i\mathbf{k} \cdot \mathbf{v}t}$, so the body drives mode $\mathbf{k}$ at $\omega = \mathbf{k} \cdot \mathbf{v}$. By (H), (S) and (L), $\langle \Phi_{\mathbf{k}} \rangle = -\chi(\mathbf{k}, \mathbf{k} \cdot \mathbf{v}) \rho_{\mathbf{k}}$, and the self-force is

$$\mathbf{F} = - \int \rho \nabla \langle \Phi \rangle d^3 x = - \int \frac{d^3 k}{(2\pi)^3} \mathbf{k} \chi''(\mathbf{k}, \mathbf{k} \cdot \mathbf{v}) |\rho_{\mathbf{k}}|^2, \quad (10)$$

because χ' is even in ω and drops out against the odd factor $\mathbf{k}$. The power dissipated in mode $\mathbf{k}$, $(\mathbf{k} \cdot \mathbf{v}) \chi''(\mathbf{k}, \mathbf{k} \cdot \mathbf{v}) |\rho_{\mathbf{k}}|^2$, is non-negative by (P), so restricting the integral to $k < k_*$ can only lower the drag. There, (W) gives $\chi''(\mathbf{k}, \mathbf{k} \cdot \mathbf{v}) = (\mathbf{k} \cdot \mathbf{v}) \partial_\omega \chi''(k, 0)$; isotropy replaces $k_i k_j$ by $\delta_{ij} k^2/3$; and (F) with (D) gives

$$\partial_\omega \chi''(k, 0) = \frac{S(k, 0)}{2k_B T_{\text{eff}}(k)} = \frac{2\pi C\hbar G}{k_B T_{\text{eff}}(k) k^2}. \quad (11)$$

The factor k^2 cancels, and Eq. (9) follows. $\square$

Remark 3.1 (Reading the bound). (i) In classical form $\gamma = (4\pi/3)G\rho_{k_*}/\kappa_{\text{FDR}}$ with $\kappa_{\text{FDR}} \equiv 2k_B T_{\text{eff}}/(C\hbar)$: $\hbar$ cancels, and the drag is the noise divided by the temperature. The rate fixes C ; the bound is on T_{eff}/C . (ii) Writing the momentum diffusion as $D_p = \frac{1}{2} \int k_i k_j S(k, 0) |\rho_{\mathbf{k}}|^2$, Eq. (9) is the Einstein relation $\gamma = D_p/(M k_B T_{\text{eff}})$. A Markovian model whose centre-of-mass momentum relaxes to a Maxwellian at T_{eff} is covered, whether or not a bath is written down, provided its friction is linear in the velocity up to orbital speeds and its noise satisfies (D), (W) and (P) at $k \leq k_*$. (iii) The bound is stated with ρ_{k_*} , not with its $k_* \rightarrow \infty$ limit $\int \rho^2/M$: when the noise resolves single nucleons the latter is dominated by granularity. At $k_* = 10/R$, $\rho_{k_*}/\bar{\rho} = 0.909$ for a uniform sphere of mean density $\bar{\rho}$, 1.42 for $\rho \propto 1 - r^2$ and 1.64 for an $n = 1$ polytrope. The uniform sphere is the conservative case. (iv) Treating a body as rigid with linear friction requires its centre-of-mass motion not to be thermal in the bath, $Mv^2 \gg k_B T_{\text{eff}}$, so that the fluctuating force is negligible against the drag. Many-body versions of dissipative collapse can differ: collective pair friction in a Fock-space formulation gives a stationary kinetic energy of order $2k_B T_\beta/N$ at collapse temperature T_β [18]. Under the stated condition such effects only lower the centre-of-mass temperature, and the bound tightens.

Lemma 3.2 (Orbital decay). *Two bodies with drags γ_1, γ_2 in a bound orbit lose orbital energy at the relative rate $\Gamma_{\text{orb}} = (\gamma_1 M_2 + \gamma_2 M_1)/(M_1 + M_2) \geq \min \gamma_i$, and $\dot{P}_b/P_b = -3\Gamma_{\text{orb}}$.*

Proof. In the centre-of-mass frame $\mathbf{v}_1 = (M_2/M)\mathbf{v}$ and $\mathbf{v}_2 = -(M_1/M)\mathbf{v}$, so the drag on the relative velocity is $-\Gamma_{\text{orb}}\mathbf{v}$. The velocity $\mathbf{w}$ of the centre of mass relative to the rest frame enters only as the constant differential acceleration $-(\gamma_1 - \gamma_2)\mathbf{w}$, which does no secular work on the orbital energy at linear order. With $\dot{E} = -\mu\Gamma_{\text{orb}}v^2$ and Kepler's law, $\dot{a} = -2\Gamma_{\text{orb}}a$ for the semi-major axis, and $P_b \propto a^{3/2}$ gives $\dot{P}_b/P_b = -3\Gamma_{\text{orb}}$, independent of the eccentricity to this order. $\square$

3.2 The double-pulsar bound

PSR J0737–3039A/B has $P_b = 0.1022515593$ d and $\dot{P}_b = -1.247920(78) \times 10^{-12}$; its orbital decay agrees with the quadrupole formula to 1.3×10^{-4} at 95% confidence [19]. Any additional decay therefore obeys

$$\Gamma_{\text{orb}} < \gamma_{\text{max}} = \frac{1.3 \times 10^{-4} |\dot{P}_b|}{3P_b} = 6.1 \times 10^{-21} \text{ s}^{-1}. \quad (12)$$

Both bodies are neutron stars. Model each as a uniform sphere of $M = 1.2 M_\odot$ and $R_{\text{NS}} = 14$ km, light and large and therefore of low density, $\bar{\rho} = 2.1 \times 10^{17} \text{ kg m}^{-3}$. Each star moves about the centre of mass at $300\text{--}325 \text{ km s}^{-1}$, so $v_{\text{max}} = 10^3 \text{ km s}^{-1}$ also covers a centre-of-mass velocity of several hundred km s^{-1} relative to the bath. With $k_* = 10/R_{\text{NS}}$, (W) and (L) are needed only up to $k_* v_{\text{max}} = 714 \text{ s}^{-1}$ at kilometre wavelengths.

Corollary 3.3 (Pulsar bound). *Every bath obeying the hypotheses of Theorem 3.1 at $k_* = 10/R_{\text{NS}}$ and $v_{\text{max}} = 10^3 \text{ km s}^{-1}$, and reproducing Eq. (1), satisfies*

$$\frac{T_{\text{eff}}}{C} > \frac{2\pi\hbar G (0.909 \bar{\rho})}{3k_B \gamma_{\text{max}}} = 3.3 \times 10^{16} \text{ K}, \quad (13)$$

equivalently $k_B T_{\text{eff}}/C > 2.8 \text{ TeV}$ or $\kappa_{\text{FDR}} > 8.6 \times 10^{27} \text{ s}^{-1}$.

The bound is an order-of-magnitude statement. Strong-field corrections (compactness ≈ 0.2) are of order one; a realistic stellar profile raises ρ_{k_*} (Remark 3.1) and strengthens it. A drag only adds decay, so a one-sided limit from the measured ratio $\dot{P}_b^{\text{GW}}/\dot{P}_{b,\text{GR}}^{\text{GW}} = 0.999963(63)$ [19] is about twice as strong; we quote the two-sided value. The two-sided band also excludes an active bath that pushes bodies: $|T_{\text{eff}}|/C > 3.3 \times 10^{16} \text{ K}$. Free fields at $T > 0$ are not constrained: their spectral weight sits on $|\omega| = ck$, so $S(k, 0) = 0$ and they fail (D). The method, binary-pulsar timing as a probe of friction from a medium, was used by Pani [20] for dark matter; we found no prior application to decoherence baths.

3.3 Consequences for known models

Dissipative Diósi–Penrose. Bahrami, Smirne and Bassi introduced dissipation into the DP model to tame its energy increase, and found that it does not reconcile the model with data unless the cutoff is large [12]. The linear-friction model of Ref. [13] adds friction at a collapse temperature T_β to the kernel $D_k = 4\pi\hbar G e^{-R_0^2 k^2}/k^2$. Its centre-of-mass damping $\Gamma = \hbar^2 \eta / (2M k_B T_\beta)$, with η the k^2 -weighted DP diffusion [21], is Eq. (9) at $C = 1$, $T_{\text{eff}} = T_\beta$.

Corollary 3.3 gives $T_\beta > 3.3 \times 10^{16}$ K for $R_0 \ll R_{\text{NS}}$. Levitated-mechanics experiments exclude $T_\beta < 10^{-13}$ K for $R_0 < 1 \mu\text{m}$ [22, 21]; the pulsar bound is stronger by 3×10^{29} . It excludes the parameter point $R_0 = 10^{-15}$ m, $T_\beta \approx 3$ K (the temperature of the cosmic microwave background) discussed in Ref. [21] by 10^{16} . At $T_\beta > 3 \times 10^{16}$ K the dissipative model is plain DP for every laboratory observable, so the X-ray bounds of Section 5 apply to it unchanged.

Classical-channel and classical–quantum gravity. The models of Kafri, Taylor and Milburn [23] and of Tilloy and Diósi [14] treat the Newtonian field as classical and sourced by a measurement record. Their dissipative generalizations thermalize to a finite effective temperature [24] and fall under Theorem 3.1 through the Einstein relation; with Diósi–Penrose noise they require $T_{\text{eff}} > 3.3 \times 10^{16} C$ K. Without friction the response of the potential to the signal is instantaneous and real, $\chi'' = 0$, and there is no drag. The same holds for the Newtonian limit of classical–quantum gravity, where the potential is classical and diffusive and friction terms are possible but not included [25]. In that framework Oppenheim *et al.* prove a trade-off between the decoherence of the quantum system and the diffusion it induces in the classical field: long coherence times require strong diffusion relative to the coupling [26]. It is the classical–quantum analogue of Remark 3.1(ii). Both relations tie the noise that decoheres matter to a second, independently measurable coefficient: the trade-off to the diffusion of the metric, the Einstein relation to the friction on matter through T_{eff} . They constrain different things and are compatible. A frictionless classical–quantum model sits at $T_{\text{eff}} = \infty$, where the drag vanishes; the trade-off then bounds the diffusion of the potential, not a temperature.

Finite-temperature matter and horizons. A system at finite temperature decoheres a nearby superposition, and the decoherence by a black-hole horizon is reproduced by modelling the hole as such a system [11]; locally it is carried by very-low-frequency quanta in the laboratory [10]. These quanta are free and on shell, so they fail (D), and horizon decoherence is not in the class Theorem 3.1 bounds. The theorem applies to a medium that fills space with Diósi–Penrose-form noise at kilometre scales; such a medium drags neutron stars, and the pulsar bound excludes it unless its temperature exceeds Eq. (13).

3.4 The gap: noise without back-action

Corollary 3.3 leaves open the limit $T_{\text{eff}} \rightarrow \infty$ at fixed noise. In that limit $\chi'' \rightarrow 0$ while S stays finite, so the drag vanishes. Operationally it is noise without back-action: a classical potential field acting on matter, whose reduced dynamics is the random-unitary DP master equation with Hermitian Lindblad operators, which is how Diósi wrote the model [2]. For a bath with a finite-dimensional or type-I algebra the limit is its tracial state; a quantum field algebra of type III has no tracial state, and only the limit, or a classical noise field, realizes it. Classical stochastic gravitational fields acting on matter belong to this class [27]. The line between reciprocal and non-reciprocal noise is therefore not the sharp one; the sharp parameter is T_{eff} . Non-reciprocal DP is the $T_{\text{eff}} = \infty$ end of the dissipative family, and the pulsars allow only $|T_{\text{eff}}|/C > 3.3 \times 10^{16}$ K around it. This end, with every bath above that temperature, is escape (G1).

It is not the only escape. An ordinary-temperature bath can still hide where its kilometre-scale behaviour differs from its laboratory-scale behaviour:

- (G2) *an infrared cut*: the kernel departs from $1/k^2$ below some $k_{\text{IR}} \gg 1/R_{\text{NS}}$. The drag then comes only from the star's density structure above k_{IR} . For a sharp-edged uniform sphere that fraction is $0.955/(k_{\text{IR}}R_{\text{NS}})$, 6.8×10^{-7} at $k_{\text{IR}} = 100 \text{ m}^{-1}$; for a real star it depends on the crust and surface profile, which we have not modelled;
- (G3) *a slow bath*: a correlation time longer than $R_{\text{NS}}/(10v_{\text{max}}) = 1.4 \text{ ms}$ at kilometre wavelengths violates (W) and suppresses the drag, while the bath can still be white at laboratory k ;
- (G4) *a scale-dependent temperature*: $T_{\text{eff}}(k)$ finite at laboratory k but above $3.3 \times 10^{16} \text{ C K}$ at $k \lesssim 10/R_{\text{NS}}$;
- (G5) *active modes*: a non-KMS state with $\omega\chi'' < 0$ at $k > k_*$ violates (P) and can cancel the drag from $k < k_*$.

A bath that co-moves with each body (non-linear entrainment, violating (L) and (S) locally) also escapes, at the price of having no rest frame.

4 Why the rate cannot come from quantized linearized gravity or a translation-invariant positive-energy state

The statements in this section are known results or direct consequences of them. We state them precisely, with their hypotheses, because the drag theorem addresses what they leave open. Throughout, the source is a static, inertial rigid body with equal branch masses, and it enters as an external c-number density.

4.1 The field channel at linearized order

In a state that is KMS at inverse temperature β , the fluctuation–dissipation theorem gives $S_B(\omega) = \hbar J(\omega) \coth(\beta\hbar\omega/2)$, with J the dissipative spectral density of B , so

$$S_B(0) = \frac{2}{\beta} \lim_{\omega \rightarrow 0} \frac{J(\omega)}{\omega} : \quad (14)$$

a floor needs an Ohmic response. The vacuum restricted to a Rindler wedge [28], or for a conformal field to a causal diamond [29], is KMS for its modular flow; describing the vacuum in modular time adds no excitations. A free massless field obeys $\square\hat{\phi} = 0$, so in a translation-invariant state its two-point function is supported on $\omega = \pm c|\mathbf{k}|$, and weight at $\omega \rightarrow 0$ comes only from $k \rightarrow 0$. For two equal spheres with form factor $\tilde{\rho}(k)$,

$$\langle |\Delta\tilde{\rho}_{\mathbf{k}}|^2 \rangle = 2\tilde{\rho}(k)^2 \left(1 - \frac{\sin kd}{kd}\right) \xrightarrow{kd \rightarrow 0} \frac{1}{3}M^2k^2d^2. \quad (15)$$

For the most favourable model, a propagating field normalized so that its reorganization energy $\pi^{-1} \int_0^\infty J \omega^{-1} d\omega$ equals E_G (Appendix A),

$$J(\omega) = \frac{G\omega}{c} \langle |\Delta\tilde{\rho}_{\omega/c}|^2 \rangle \simeq \frac{GM^2d^2}{3c^3} \omega^3 \quad (\omega d/c \ll 1), \quad (16)$$

which is super-Ohmic. Only a difference of branch *energies* makes J Ohmic. That is the channel of the graviton master equations [6, 7, 8], whose rates are first order in G for superpositions

of different energies and vanish for the equal-mass spatial superposition considered here. In linearized general relativity the static density has no propagating scalar partner: its potential is fixed by the constraint, an instantaneous real kernel that dissipates nothing, and it fluctuates only through the matter that sources it, at order G^2 .

These results agree with direct calculations. In the graviton vacuum, the decoherence functional of a spatial superposition that is opened and closed again approaches a constant set by $(mv/M_P)^2$ at late times rather than growing linearly [30]. The which-path information radiated by a superposition is set by the time taken to recombine it, not by the time it is held apart, $N \sim Q^2/T_{\text{rec}}^4$ in Planck units for a mass quadrupole Q [31], and in the inertial Minkowski vacuum it does not grow with the holding time [9].

Proposition 4.1 (Field channel). *In any translation-invariant state of the free linearized field, including states that are KMS for modular flow, $S_B(0) = 0$ for a static inertial source with equal branch masses. The dephasing exponent saturates.*

4.2 The clock channel

A clock at $\mathbf{x}_{\text{obs}}$ whose rate depends on the branch through gravitational redshift [32], as for the observer adjoined in the crossed-product construction [33, 34], gives a structureless source ($E = Mc^2$) a coherence of constant modulus that rotates at

$$\omega_{\text{rel}} = \frac{M |\Phi_L - \Phi_R|(\mathbf{x}_{\text{obs}})}{\hbar} = \frac{GM^2 d}{\hbar(r^2 - d^2/4)}, \quad (17)$$

on the axis at distance r from the midpoint; at $r = d$ this is $(4/3) GM^2/(\hbar d)$. It shifts the fringes and leaves their contrast intact. For any number of clocks with Hamiltonians H_k at $\mathbf{x}_k$ in any joint state, the branch-relative evolution is generated by $\sum_k H_k \Delta\Phi(\mathbf{x}_k)/c^2$ with $\Delta\Phi = GM\varphi(\mathbf{x})$, φ independent of G and M , so

$$z(t) = \left\langle \exp\left(\frac{i GMt}{\hbar c^2} \sum_k H_k \varphi(\mathbf{x}_k)\right) \right\rangle. \quad (18)$$

Proposition 4.2 (Clock channel). *For redshift-coupled clocks with source-independent statistics, the branch coherence depends on G , M and t only through GMt . Hence $R_M = R_t$, whereas Eq. (1) requires $R_M = 4$, $R_t = 2$.*

Finite clock-energy variances give $F \propto G^2 M^2 t^2$, the time-dilation dephasing of interfering clocks [35, 36]. Cauchy-like energy distributions give $F \propto GMt$, exponential at intermediate times but linear in M ; for a spectrum bounded below the decay cannot remain exponential at late times [37]. A factor M^2 appears only if the probe is a second massive body, the setting of gravitationally mediated entanglement [38, 39].

4.3 Translation-invariant states with positive energy

The obstruction in Proposition 4.1 does not need linearity, within a Gaussian treatment. Let the potential Φ (local or gravitationally dressed) and the bath state Ω satisfy:

- (A) Ω is invariant under spacetime translations $U(a) = e^{iP \cdot a}$;
- (B) the joint spectrum of P^μ lies in the closed forward cone $\bar{V}_+$;

- (C) $\Phi(x) = U(x)\Phi(0)U(x)^\dagger$, and the connected two-point spectral measure μ of Φ is locally finite;
- (D) the connected part has no atom at $p = 0$;
- (E) $|\Delta\tilde{\rho}(\mathbf{k})|^2 \leq M^2 k^2 d^2$, which holds for a rigid body since $|\Delta\tilde{\rho}|^2 = 4|\tilde{\rho}|^2 \sin^2(\mathbf{k} \cdot \mathbf{d}/2)$ and $|\tilde{\rho}| \leq M$.

Proposition 4.3 (Spectrum condition). *Under (A)–(E), the symmetrized two-point function of B has $S_B(0) = 0$. In a treatment that keeps the bath to second cumulant and the source as an external c -number, no such bath gives a decoherence rate.*

Proof. With $P|\Omega\rangle = 0$, (A) and (C), $\langle\Omega|\Phi(x)\Phi(0)|\Omega\rangle_c = \int e^{-ipx} d\mu(p)$ with $\text{supp } \mu \subset \bar{V}_+$ by (B). The zero-frequency noise of B is the density at $p^0 = 0$ of $|\Delta\tilde{\rho}(\mathbf{p})|^2 d\mu(p)$ and its mirror image, $S_B(0) = c_0 \lim_{\varepsilon \rightarrow 0} \varepsilon^{-1} \int_{0 \leq p^0 \leq \varepsilon} |\Delta\tilde{\rho}|^2 d\mu$, with c_0 a convention-dependent constant. On the support $|\mathbf{p}| \leq \varepsilon/c$, so by (E)

$$S_B(0) \leq c_0 \lim_{\varepsilon \rightarrow 0} \frac{M^2 d^2}{c^2} \varepsilon \mu(\{p^0 \leq \varepsilon\}) = 0 \quad (19)$$

by (C), with (D) removing an atom at the origin. $\square$

The content is the spectral support of Wightman functions [40], applied to which-path noise. (B) is an assumption about the bath state; in asymptotically flat gravity one would justify it by positivity of the total energy, which we do not derive here. The proposition does not reach higher cumulants or a quantized source. It covers the modular state of a diamond, because that state is the vacuum. Within the Gaussian treatment, a white floor therefore needs a state that breaks (A) or (B), for example one with a rest frame. Relatedly, Diósi identifies a hidden defect in a proposed Lorentz-invariant GKLS master equation and argues that no such Markovian equation should exist [41]. The escapes are the broken assumptions:

- *(B) fails* for every thermal, stationary non-equilibrium or infinite-temperature state: the generator of time translations in its GNS representation has spectrum $\mathbb{R}$. This is the class of Section 3. Relativistic collapse models escape the same way: a mediating field at each spacetime point [42], a scalar field whose energy spectrum takes all real values [43], or classical noise with a Lorentz-invariant correlation [44], which as a classical field need not obey (B).
- *Killing horizons* replace the positive generator of laboratory time by a boost generator with spectrum $\mathbb{R}$. They decohere at rates set by horizon data that vanish for an inertial laboratory without a horizon [45, 9, 10, 46].
- *(C) fails* only for an infrared accumulation of spectral weight at the tip of the cone; for a static displaced source we expect none, which is argued, not shown.
- *(D) fails* in the source's own Gauss-law sector, where boosts are spontaneously broken. We conjecture that the corrections are of order GM relative to the leading term, $G^2 M^3$ and higher, but have not shown it.
- *Non-unitary dynamics* (Penrose-type objective reduction [3]) lies outside the bath framework.

Table 1. Published 95% lower bounds on R_0 , converted to $C = 1$ by Eq. (21). The Gran Sasso, MAJORANA and XENONnT searches look for spontaneous X- and γ -ray emission from germanium or xenon; the others bound heating or force noise.

Search	published R_0 (m)	C_{pub}	R_0 at $C = 1$ (m)
XENONnT, 1–140 keV [17]	4.5×10^{-10}	1	4.5×10^{-10}
MAJORANA, 19–100 keV [16, 48]	2.54×10^{-10}	8π	8.7×10^{-11}
Gran Sasso, 1–3.8 MeV [15]	0.54×10^{-10}	8π	1.8×10^{-11}
Cryostat heat leak, 10 pW/kg [47]	4.6×10^{-12}	1	4.6×10^{-12}
Neptune internal heat, 20 pW/kg [47]	3.7×10^{-12}	1	3.7×10^{-12}
LISA Pathfinder force noise [49]	4.0×10^{-14}	1	4.0×10^{-14}
Neutron-star heating [50]	10^{-13}	1/2	1.3×10^{-13}

5 The survivor and its allowed region

Definition 5.1 (Surviving rate clause). The full mass density of matter, nuclei and electrons, convolved with a normalized Gaussian $\propto e^{-r^2/2R_0^2}$, sees classical, symmetric, white potential noise with the correlation (5) and does not act back on it ($T_{\text{eff}} = \infty$). The master equation is Diósi’s with coefficient C , and $\Gamma = CE_G/\hbar$ with E_G evaluated on the smeared density.

There is no friction and no detailed-balance factor, so every bound on DP noise applies directly. Two parameters remain, C and R_0 . The analysis does not fix C ; it is a free parameter of order one.

Why a density smearing. The smearing must give the momentum diffusion a finite k^2 moment. Smearing the full density by $e^{-k^2 R_0^2/2}$ per factor gives the kernel $(4\pi/k^2) e^{-k^2 R_0^2}$ and, for any constituent of mass m , nucleus or electron,

$$\frac{d\langle p^2 \rangle}{dt} = 4\pi C \hbar G m^2 \int \frac{d^3 k}{(2\pi)^3} e^{-k^2 R_0^2} = \frac{C \hbar G m^2}{2\sqrt{\pi} R_0^3}, \quad \frac{P}{M} = \frac{1}{2m^2} \frac{d\langle p^2 \rangle}{dt} = \frac{CG \hbar}{4\sqrt{\pi} R_0^3}, \quad (20)$$

a heating rate per unit mass independent of m . A cutoff of the kernel instead, $(1 - e^{-r/R_0})/r$ with Fourier transform $4\pi/[k^2(1 + k^2 R_0^2)]$, falls only as k^{-4} ; for constituents of size r_N its k^2 moment is $2R_0/r_N$ times the Gaussian one and diverges for point electrons. The same divergence appears if only nuclei are smeared. We therefore smear the full mass density with a Gaussian, the form used in the published bounds.

Converting the bounds. Spontaneous radiation and heating are both proportional to C/R_0^3 . A lower bound R_{pub} published at normalization C_{pub} becomes

$$R_0 > R_{\text{pub}} (C/C_{\text{pub}})^{1/3}. \quad (21)$$

Table 1 lists the bounds at $C = 1$. The normalization of each was checked against an independent statement: Eq. (20) reproduces the heating figures of Ref. [47], and the XENONnT statement that it improves on MAJORANA by a factor of five [17] holds only for the corrected MAJORANA value [48] read at $C = 8\pi$, where the ratio is 5.2.

Proposition 5.2 (Allowed region). *The surviving rate clause is consistent with the data of Table 1 if and only if*

$$R_0 > 4.5 \times 10^{-10} \text{ m } C^{1/3}, \quad (22)$$

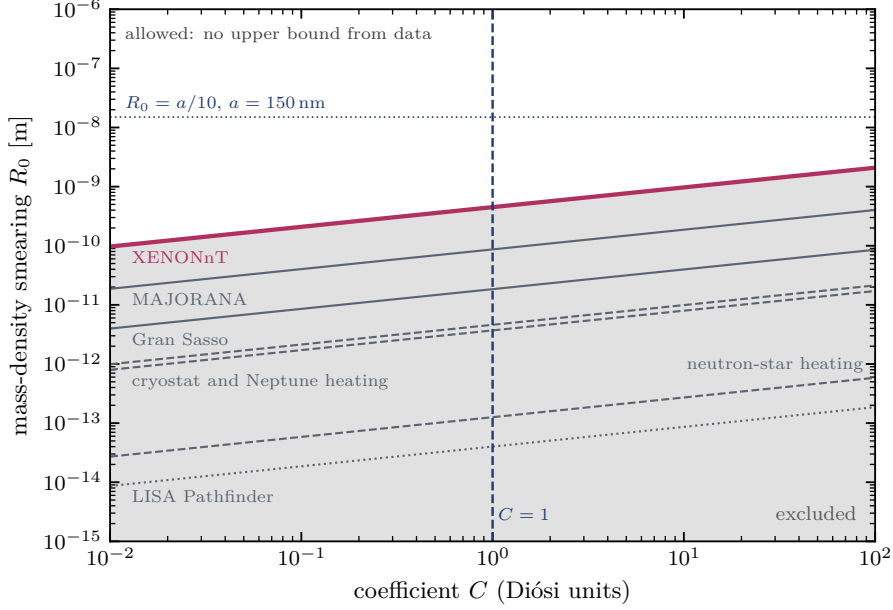


Figure 1. The (C, R_0) plane of the surviving rate clause. Each line is a published lower bound converted with Eq. (21); XENONnT (thick) binds at every C , and the grey region is excluded. Nothing bounds R_0 from above. The dashed line marks Diósi's normalization $C = 1$. Below the dotted line, $R_0 < a/10$ for a sphere of radius $a = 150$ nm, the rate of Section 6 is independent of R_0 to better than 3%.

set by XENONnT at every C . No published noise observable bounds R_0 from above.

Proof. All bounds scale as Eq. (21), so their order does not depend on C , and XENONnT is the largest at $C = 1$. For the upper side: the kernel $e^{-k^2 R_0^2}/k^2$ is pointwise non-increasing in R_0 , so every positive functional of the noise, E_G , heating, force noise and emission, is non-increasing in R_0 . A null result for such an observable bounds R_0 only from below. $\square$

At the edge of the region the heating is $P/M = G\hbar/[4\sqrt{\pi}(4.5 \text{ \AA})^3] = 1.1 \times 10^{-17} \text{ W kg}^{-1}$, independent of C and six orders below the cryostat floor. The leading power law R_0^{-3} is used down to atomic scales; refined atomic emission rates [51] would move Eq. (22) by the cube root of the change in rate. Two candidate upper bounds do not apply. Inverse-square-law tests constrain a smearing of the Newtonian *force*; here the mean field is that of general relativity and only the noise is smeared. A requirement that the noise also make macroscopic bodies classical within a fixed time [52] would cap R_0 ; the survivor is a decoherence clause, not a solution of the measurement problem, and we do not impose it. Figure 1 shows the region.

6 The sharp prediction: a measurement of C

Every non-interferometric bound measures C/R_0^3 and cannot separate C from the free R_0 . An interferometric rate can, in one regime, once the other decoherence channels are accounted for.

Proposition 6.1 (Saturated rate). *For a uniform sphere of mass M , radius a and density ρ , with interatomic spacing $\lesssim R_0 \ll a$, in a superposition with separation $d \geq 2a$,*

$$\Gamma = \frac{C}{\hbar} GM^2 \left(\frac{6}{5a} - \frac{1}{d} \right) \left[1 - \mathcal{O}(R_0^2/a^2) \right], \quad \Gamma_{\text{sat}} \equiv \lim_{d \rightarrow \infty} \Gamma = \frac{6C}{5\hbar} G \left(\frac{4\pi\rho}{3} \right)^{1/3} M^{5/3}. \quad (23)$$

The rate does not depend on R_0 .

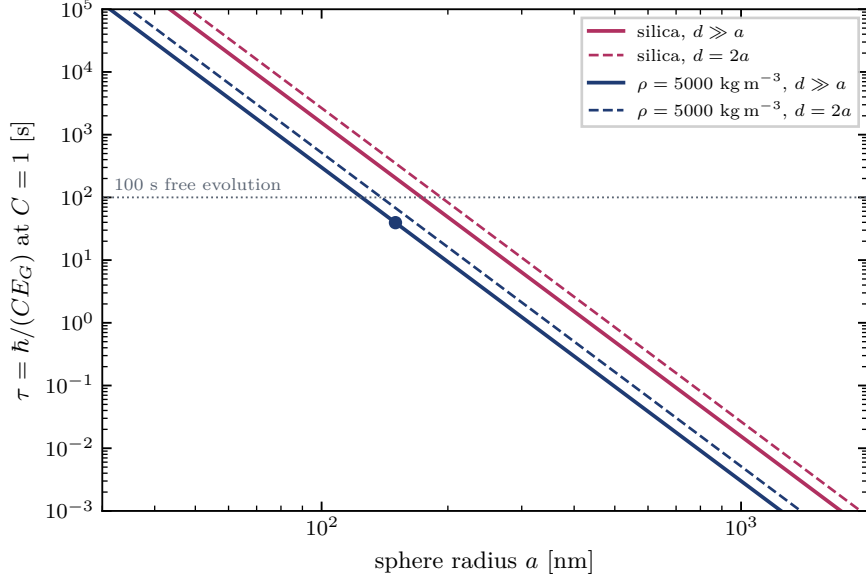


Figure 2. Decoherence time $\tau = \hbar/(CE_G)$ at $C = 1$ against sphere radius, for silica and for $\rho = 5000 \text{ kg m}^{-3}$, at $d \gg a$ (solid) and $d = 2a$ (dashed), in the saturated regime $R_0 \ll a$ where it does not depend on R_0 . The horizontal line is a free-evolution time of 100 s. The time scales as $1/C$.

Proof. Equation (6) gives the unsmeared energy (Appendix A). Smearing multiplies the kernel by $e^{-k^2 R_0^2} = 1 - k^2 R_0^2 + \dots$, and the $k^2 R_0^2$ term turns $4\pi/k^2$ into the contact term $-4\pi R_0^2 \delta^3(\mathbf{x} - \mathbf{y})$. For a uniform ball $\int \rho^2 d^3x = 3M^2/(4\pi a^3)$, so the self term $(G/2) \cdot 2 \cdot (6/5)M^2/a$ changes by $-G \cdot 3M^2 R_0^2/a^3$, a relative change $-(5/2)(R_0/a)^2$. The cross term between disjoint balls is unchanged up to terms exponentially small in $(d - 2a)/R_0$. With $a = (3M/4\pi\rho)^{1/3}$ the limit $d \rightarrow \infty$ gives Γ_{sat} . $\square$

Physical content. Once the spheres no longer overlap, the rate is set by the self-energy of each branch's density, $\frac{6}{5}GM^2/a$, minus the cross term GM^2/d ; the distance enters only through that cross term. The separation dependence is $E_G(d)/E_G(\infty) = 1 - 5a/6d$: 0.583 at $d = 2a$, within 10% of saturation only for $d \geq 8.33a$. Below $d = 2a$ the overlapping-sphere formula of Appendix A applies and tends to $GM^2 d^2/(2a^3)$. In the allowed region (22) R_0 exceeds the interatomic spacing of about $2.5 \text{ \AA}$ for $C > 0.17$, so the continuum density applies; granularity then adds a relative correction of order $\bar{m}a/(MR_0)$, with $\bar{m}$ the mean nuclear mass, which is below 10^{-11} for the benchmark below.

Benchmarks. For $1 \text{ }\mu\text{g}$ of silica ($\rho = 2200 \text{ kg m}^{-3}$, $a = 47.7 \text{ }\mu\text{m}$) at $d = 1 \text{ mm}$, $E_G = 24.2 GM^2/d$ and

$$\tau = \frac{\hbar}{CE_G} = \frac{65.4 \text{ ps}}{C}, \quad (24)$$

rising to $108 \text{ ps}/C$ at $d = 2a$ and falling to $62.8 \text{ ps}/C$ as $d \rightarrow \infty$. The smearing correction is $(5/2)(R_0/a)^2 = 2 \times 10^{-10}$ at $R_0 = 4.5 \text{ \AA}$. A space platform of the MAQRO class [53, 54], with free evolution of order 100 s, reaches the saturated regime: a sphere of $a = 150 \text{ nm}$ and $\rho = 5000 \text{ kg m}^{-3}$ has $\tau = 39.5 \text{ s}/C$ for $d \gg a$ and $67.8 \text{ s}/C$ at $d = 2a$. A silica sphere needs $a \geq 173 \text{ nm } C^{-1/5}$ for $\Gamma \times 100 \text{ s} \geq 1$ at $d \gg a$. Figure 2 shows the rate against radius.

What a measurement determines. In the saturated regime $C = \hbar\Gamma_{\text{grav}}/E_G$, with E_G known from M , a and d , where Γ_{grav} is the measured visibility decay rate after the known decoherence channels (residual-gas collisions, black-body emission, scattering and absorption, trap and recoil noise) have been subtracted. With those subtracted, one measurement determines C . A rate below the prediction means either a smaller C or $R_0 \gtrsim a$; the two are separated by the radius dependence, $\Gamma \propto a^5 \rho^2$ at fixed C in the saturated regime against a suppression once R_0 approaches a . Distinguishing two values of C needs Γ_{grav} to the fractional precision of their difference. A residual rate of second order in G , or none, excludes the survivor at the tested (C, R_0) .

7 Discussion

7.1 What remains open

- *The survivor* (Definition 5.1): noise without back-action, or a bath at $|T_{\text{eff}}|/C > 3.3 \times 10^{16}$ K (escape G1). Neither quantized linearized gravity, nor a translation-invariant positive-energy state in a Gaussian treatment, nor a finite-temperature bath obeying the hypotheses of Theorem 3.1 at kilometre scales supplies it.
- *Baths that differ at kilometre scales* (G2–G5): an infrared cut of the kernel, a correlation time longer than 1.4 ms, a temperature that rises toward long wavelengths, or active modes above k_* . The infrared cut is bounded only through the high- k structure of neutron stars, which we have not modelled; closing these escapes needs that structure, or bodies resolved at other scales.
- *Horizons*: real rates set by horizon data [45, 9, 10, 46], absent for an inertial laboratory without a horizon.
- *Beyond the Gaussian treatment*: higher cumulants, a quantized source, non-tempered infrared sectors and the Gauss-law sector of the source (Section 4) are not excluded by any argument here.
- *Objective reduction* in Penrose’s sense [3]: outside the bath framework.

Corollary 3.3 rests on a uniform-sphere neutron-star model and on treating the star as a rigid body with linear friction up to 10^3 km s^{-1} ; both are conservative within the condition of Remark 3.1(iv), and strong-field corrections are of order one. We have found no earlier bound on decoherence-bath friction from orbital decay, but searched only arXiv and the general literature.

7.2 Consequences

1. Dissipative Diósi–Penrose models need $T_\beta > 3.3 \times 10^{16}$ K, a bound 3×10^{29} above the laboratory one; at such temperatures they coincide with the non-dissipative model.
2. A rate of first order in G observed for an isolated, inertial, equal-mass spatial superposition would not confirm a derivation from quantized gravity or from a thermal medium. By Corollary 3.3 and Section 4 it would signal noise without back-action, one of the escapes G2–G5, a horizon, physics beyond the Gaussian treatment (higher cumulants, a quantized source, non-tempered infrared or Gauss-law sectors), or non-unitary reduction.

3. In the saturated regime of Proposition 6.1, once the other decoherence channels are subtracted, one visibility measurement gives $C = \hbar\Gamma_{\text{grav}}/E_G$. For 1 μg of silica the prediction is $\tau = 65 \text{ ps}/C$; a MAQRO-class platform with 100 s of free evolution reaches the regime with spheres of 150–175 nm radius.
4. Radiation and heating searches can raise the floor (22) on R_0 but cannot exclude the survivor, since no noise observable bounds R_0 from above.

7.3 Emergent-gravity programmes

Programmes in which the Einstein equation is an equilibrium condition of diamond entanglement [55, 56], such as the quantum-geometric correspondence that motivated this work, sometimes look to the coarse-grained geometry for a decoherence mechanism. The results above constrain that hope. If the geometry is described by the vacuum, Proposition 4.3 gives no Gaussian white floor; if it is described as a medium at finite temperature, the medium drags and Corollary 3.3 requires $T_{\text{eff}}/C > 3.3 \times 10^{16} \text{ K}$. A rate linear in G in such a programme is therefore an additional postulate of the Diósi type, with C not fixed, unless the medium uses one of the escapes G2–G5.

A The energy E_G of two uniform spheres

With $\int d^3k (2\pi)^{-3} 4\pi k^{-2} f(\mathbf{k}) = (2/\pi) \int_0^\infty dk \langle f \rangle$ for any angle average $\langle f \rangle$, Eq. (1) reads

$$E_G = \frac{G}{\pi} \int_0^\infty dk \langle |\Delta \tilde{\rho}_{\mathbf{k}}|^2 \rangle, \quad (25)$$

with the angle average (15) and the form factor of a uniform sphere, $\tilde{\rho}(k) = 3M(\sin ka - ka \cos ka)/(ka)^3$. The spectral density (16) has reorganization energy $\pi^{-1} \int_0^\infty J \omega^{-1} d\omega$ equal to (25), which fixes its normalization; its small- ω expansion gives the ω^3 law. In position space, with $\iint \rho \rho |\mathbf{x} - \mathbf{y}|^{-1} = (6/5)M^2/a$ for each ball (twice the self-energy $\frac{3}{5}GM^2/a$ divided by G) and $\iint \rho_L \rho_R |\mathbf{x} - \mathbf{y}|^{-1} = M^2/d$ for disjoint balls (Newton’s shell theorem),

$$E_G = \frac{G}{2} \left[2 \cdot \frac{6M^2}{5a} - \frac{2M^2}{d} \right] = GM^2 \left(\frac{6}{5a} - \frac{1}{d} \right) \quad (d \geq 2a). \quad (26)$$

For overlapping balls, $s = d/a < 2$, the cross term follows from the overlap integral of two uniform balls, and

$$E_G = \frac{GM^2}{a} \left(\frac{s^2}{2} - \frac{3s^3}{16} + \frac{s^5}{160} \right) \xrightarrow{s \rightarrow 0} \frac{GM^2 d^2}{2a^3}. \quad (27)$$

Both branches agree at $s = 2$, where $E_G = 0.7 GM^2/a$. Quadrature of (25) reproduces both forms to better than one part in 10^6 at $s = 0.3, 1.2, 3$ and 10 .

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