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A Geometric Unification of Particles, Spacetime, and Cosmology: From Private Spacetimes to Emergent Gravity and Dark Energy

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A Geometric Unification of Particles, Spacetime, and Cosmology: From Private Spacetimes to Emergent Gravity and Dark Energy

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Abstract

We present a relativistic and quantizable geometric framework in which each quantum particle possesses an intrinsic spatial geometry governed by a universal stiffness constant A_0 and sourced by its wavefunction. The model is defined by a covariant action coupling a complex scalar field to the metric, including a curvature-coupling term $\xi R|\phi|^2$. From this action we derive coupled Einstein–Klein–Gordon equations; repulsive self-gravity emerges naturally, preventing gravitational collapse and stabilizing quantum matter. Interactions stitch together individual world-blocks via junction conditions that enforce continuity of the metric and extrinsic curvature. In the continuum limit, the collective dynamics yield an emergent spacetime satisfying the full Einstein equations, with Newton's constant $G = c^4/(8\pi A_0)$. The harmonic zero-mode of the strain field may acquire a vacuum expectation value set by the Hubble radius, providing a natural explanation for dark energy. Canonical and path-integral quantization are outlined; the problem of time is resolved relationally through stitching constraints, giving rise to an emergent global time operator. Quantum fluctuations of the lapse network can imprint on the primordial power spectrum and generate a stochastic gravitational wave background. The model predicts testable signatures across scales, from matter-wave interferometry to cosmology, offering a unified, falsifiable foundation for quantum gravity.

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1. Introduction

The unification of quantum mechanics and general relativity remains a central challenge in theoretical physics. Conventional approaches quantize a pre-existing classical spacetime, leading to conceptual difficulties such as the problem of time, non-renormalizability, and the cosmological constant problem [1,2]. In this work we develop a geometric framework in which each quantum particle is endowed with an intrinsic spatial geometry, governed by a universal stiffness constant A_0 and sourced by its wavefunction. The familiar macroscopic universe emerges only when many such individual geometries are stitched together through interactions.

This idea resonates with relational and discrete approaches to quantum gravity [3–5]. In previous non-relativistic work [6,7], we showed that each particle's wavefunction sources a strain field on a private spatial slice, leading to repulsive self-gravity and Newtonian gravity through stitching. Here we present the fully relativistic extension.

Starting from a covariant action (2) that couples a complex scalar field to the metric with a curvature-coupling term $\xi R|\phi|^2$, we derive the coupled Einstein–Klein–Gordon equations (3)–(5). The model exhibits repulsive self-gravity, preventing gravitational collapse and stabilizing quantum matter.

When multiple particles interact, their private spacetimes are stitched along common boundaries; the junction conditions (9)–(10) enforce continuity of the metric and extrinsic curvature. Summing their actions and taking the continuum limit yields an emergent spacetime that satisfies the full Einstein

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equations (27), with Newton's constant $G = c^4/(8\pi A_0)$. The harmonic zero-mode of the strain field may acquire a vacuum expectation value set by the Hubble radius, naturally explaining dark energy.

We quantize the model using canonical and path-integral methods. The multiple proper times and stitching constraints resolve the problem of time relationally, giving an emergent global time operator. Quantum fluctuations of the lapse network can imprint on the primordial power spectrum and generate a stochastic gravitational wave background.

The paper is organized as follows. Section 2 defines the intrinsic geometry of a particle, the action (2), and the field equations. Section 3 describes the stitching of world-blocks and the emergence of general relativity. Section 4 presents quantization and the resolution of the problem of time. Section 5 gives a non-perturbative derivation of dark energy. Section 6 discusses black holes. Section 7 addresses foundational issues. Section 8 compiles experimental signatures. Section 9 concludes. Appendices contain the detailed derivations of the field equations and the non-relativistic reduction.

2. Relativistic World-Block Model

2.1. Covariant Definition of a World-Block

We begin with a primitive geometric object: a four-dimensional Lorentzian manifold $\mathcal{M}$ that represents the private spacetime of a single particle. This manifold is not embedded in any background; it is the fundamental arena. It comes equipped with a distinguished timelike geodesic γ parametrized by proper time τ, which represents the particle's world-line. In a sufficiently small neighbourhood of γ, Fermi–Walker coordinates (τ, x^i) exist [8,9], providing a natural foliation by spacelike slices Σ_τ orthogonal to γ. On each slice, the induced spatial metric is written as

$$h_{ij}(\tau, \mathbf{x}) = \delta_{ij} + \varepsilon_{ij}(\tau, \mathbf{x}), \tag{1}$$

where ε_{ij} is a symmetric tensor field encoding deviations from flatness. The trace of the strain is $\varepsilon = h^{ij}\varepsilon_{ij}$. The fundamental dynamical fields are the spacetime metric $g_{\mu\nu}$ (or equivalently the strain ε_{ij} together with the lapse and shift) and matter fields collectively denoted by ϕ, which may include scalars, spinors, gauge fields, etc., all minimally coupled to the metric.

Definition 1 (World-block). *A "world-block" is a tuple $(\mathcal{M}, g_{\mu\nu}, \gamma, \phi)$ satisfying the following:*

- *$\mathcal{M}$ is a four-dimensional, oriented, time-oriented Lorentzian manifold with metric $g_{\mu\nu}$ of signature $(-,+,+,+)$.*
- *$\gamma : \mathbb{R} \to \mathcal{M}$ is a timelike geodesic, the particle's world-line serving as the "centroid" of the block.*
- *ϕ denotes a collection of matter fields (scalars, spinors, etc.) defined on $\mathcal{M}$.*
- *The metric $g_{\mu\nu}$ is smooth throughout $\mathcal{M}$ and approaches the Minkowski metric $\eta_{\mu\nu}$ asymptotically in regions far from γ.*
- *In Fermi–Walker coordinates adapted to γ, the spatial part of the metric takes the form (1) and the fields satisfy the field equations derived below.*

The boundary $\partial\mathcal{M}$ of a world-block consists of several components:

- Future and past timelike boundaries extending to future and past null infinity $\mathscr{I}^+$ and $\mathscr{I}^-$, respectively.
- Spacelike boundaries at spatial infinity i^0.
- Stitching surfaces $\mathcal{S}_{ab}$ where this block meets another block $\mathcal{M}_b$ when particles interact. These surfaces are the loci where the junction conditions (9)–(10) will be imposed.

This definition captures the essential geometric structure of a private spacetime: it is a complete, asymptotically flat manifold with a distinguished world-line, capable of coupling to arbitrary matter fields and of being joined to other such manifolds through interaction boundaries. The asymptotic flatness ensures that isolated particles are well-defined and that the usual notions of energy and momentum apply.

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2.2. Fundamental Action Principle

We postulate an action for a single world-block that is generally covariant and reduces to the non-relativistic action in the slow-motion, weak-field limit. The action contains three parts: the Einstein–Hilbert term, the matter Lagrangian for the scalar field, and a curvature-coupling term:

$$S = \int_{\mathcal{M}} d^4x \sqrt{-g} \left[\frac{A_0}{2} R - \frac{1}{2} g^{\mu\nu} \partial_\mu \phi^* \partial_\nu \phi - \frac{1}{2} m^2 c^2 |\phi|^2 - \xi R |\phi|^2 \right]. \tag{2}$$

Here A_0 is a universal stiffness constant with dimensions of action (since the Ricci scalar R has dimension L^{-2}, $\sqrt{-g} d^4x$ has dimension L^4, so A_0 must have dimensions of action to make the integral dimensionless in units where $\hbar = 1$; in SI units, A_0 has dimensions ML^2T^{-1}). The parameter m is the mass, and ξ is a dimensionless coupling constant. The term $\xi R |\phi|^2$ couples the scalar field directly to the Ricci scalar, generalizing the non-relativistic coupling $\varepsilon |\psi|^2$ in [6,7]. In the non-relativistic limit, $\phi = e^{-imc^2\tau/\hbar} \psi$ with ψ slowly varying, and R reduces to a function of the strain; one recovers the earlier coupling with an effective strength determined by ξ.

2.3. Field Equations

Varying the action (2) with respect to $g^{\mu\nu}$ yields the Einstein equations (see Appendix A for a complete derivation):

$$\left(A_0 - 2\xi |\phi|^2\right) G_{\mu\nu} = T^{\phi}_{\mu\nu} + 2\xi \left(\nabla_\mu \nabla_\nu - g_{\mu\nu} \nabla^2\right) |\phi|^2, \tag{3}$$

where $T^{\phi}_{\mu\nu}$ is the canonical stress-energy tensor of the scalar field:

$$T^{\phi}_{\mu\nu} = \partial_\mu \phi^* \partial_\nu \phi + \partial_\nu \phi^* \partial_\mu \phi - g_{\mu\nu} \left(\partial_\alpha \phi^* \partial^\alpha \phi + m^2 c^2 |\phi|^2 \right). \tag{4}$$

The terms involving ξ represent direct coupling between the scalar field and curvature. Varying with respect to ϕ^* gives the Klein–Gordon equation with a curvature coupling:

$$\left(\nabla^2 - m^2 c^2 - \xi R \right) \phi = 0. \tag{5}$$

This is a wave equation for ϕ in a curved spacetime, with an effective mass $m^2_{\text{eff}} = m^2 c^2 + \xi R$.

2.4. Trace of the Einstein Equations and Non-Relativistic Reduction

Taking the trace of Eq. (3) and using the definition of $T^{\phi}_{\mu\nu}$ (4) yields a relation between the Ricci scalar R and the scalar field ϕ:

$$\left(A_0 - 2\xi |\phi|^2\right) R = -2\left(|\nabla \phi|^2 + 2m^2 c^2 |\phi|^2\right) + 6\xi \nabla^2 |\phi|^2. \tag{6}$$

In the weak-field, slow-motion limit, the metric is expanded as in Appendix B. The non-relativistic reduction, detailed in Appendix B, leads to the Poisson equation

$$\nabla^2 \Phi = 4\pi G \rho_m, \qquad \rho_m = m|\psi|^2, \qquad G = \frac{c^4}{8\pi A_0}, \tag{7}$$

and the Schrödinger equation with a ξ-dependent correction:

$$i\hbar \partial_\tau \psi = -\frac{\hbar^2}{2m} \nabla^2 \psi - m\Phi\psi - \frac{\hbar^2}{2m} \xi R \psi + \mathcal{O}(c^{-4}). \tag{8}$$

For a point mass, $\Phi(r) = -Gm/r$, giving a leading repulsive force $F_r = +Gm^2/r^2$. The ξR term gives subleading corrections.

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Thus, the relativistic world-block model predicts repulsive self-gravity at leading order, independent of ξ. The parameter ξ influences subleading corrections and the precise relation between A_0 and G only at higher orders.

It is important to distinguish between self-gravity (the effect of a particle on its own private spacetime) and mutual gravity (the interaction between different particles). While self-gravity is repulsive, as shown above, the Einstein equations (3) are the same equations that, in the context of many interacting particles, produce the standard attractive gravitational interaction. When multiple world-blocks are stitched together, their collective stress-energy tensors add, and in the continuum limit, the emergent spacetime satisfies the standard Einstein equations with attractive Newtonian gravity. This is demonstrated in Section 3.4.

3. Stitching and Emergent General Relativity

3.1. Stitching Conditions and Junction Action

Consider a collection of world-blocks $\mathcal{M}_a$, each with its own metric $g_{\mu\nu}^{(a)}$ and matter fields ϕ_a. When two blocks a and b interact, they are stitched along a common boundary—a hypersurface $\mathcal{S}_{ab}$ which may be timelike or spacelike depending on the nature of the interaction. The matching conditions that ensure a smooth union are the GR junction conditions [9].

The fundamental requirement is the continuity of the metric across the junction:

$$g_{\mu\nu}^{(a)}|_{\mathcal{S}} = g_{\mu\nu}^{(b)}|_{\mathcal{S}}. \tag{9}$$

Additionally, the extrinsic curvature $K_{\mu\nu}$ of $\mathcal{S}$ embedded in each block must satisfy

$$K_{\mu\nu}^{(a)}|_{\mathcal{S}} + K_{\mu\nu}^{(b)}|_{\mathcal{S}} = \frac{8\pi G}{c^4}\left(S_{\mu\nu} - \frac{1}{2}\gamma_{\mu\nu}S\right), \tag{10}$$

where $S_{\mu\nu}$ is the surface stress-energy tensor localized on $\mathcal{S}$ and $\gamma_{\mu\nu}$ is the induced metric. For smooth stitching without surface matter, we set $S_{\mu\nu} = 0$, yielding the continuity of the extrinsic curvature (up to a sign due to opposite normals).

These conditions can be derived from the action principle by including the Gibbons–Hawking–York boundary terms [10] and varying the total action. The total action for a single block with boundary is

$$S = \frac{A_0}{2}\int_{\mathcal{M}} d^4x\sqrt{-g}\,R + \int_{\mathcal{M}} d^4x\sqrt{-g}\,\mathcal{L}_{\text{matter}} + A_0\int_{\partial\mathcal{M}} d^3\xi\,\sqrt{|\gamma|}\,\epsilon K, \tag{11}$$

where the GHY term ensures a well-posed variational principle. For two blocks sharing a common boundary $\mathcal{S}$, we orient the normals so that $n_{(b)} = -n_{(a)}$ on $\mathcal{S}$. The total action is the sum of the actions for each block, including their respective GHY terms. The junction is treated as an interior boundary; the GHY terms on $\mathcal{S}$ from each block combine to

$$A_0\int_{\mathcal{S}} d^3\xi\,\sqrt{|\gamma|}\,\left(K_{(a)} + K_{(b)}\right). \tag{12}$$

Varying the total action and requiring stationarity yields the bulk Einstein equations together with the junction conditions (9)–(10).

3.2. Temporal Constraints and Lapse Functions

The stitching of world-blocks also correlates their proper times. Let each block $\mathcal{M}_a$ be equipped with a proper time τ_a measured along its central world-line γ_a. On a junction surface $\mathcal{S}_{ab}$, we choose a foliation of each block by spacelike slices orthogonal to γ_a and γ_b. Let t be a coordinate time that matches on $\mathcal{S}_{ab}$ (possible because the induced metrics are equal). Then

$$d\tau_a = N_a(t)\,dt, \qquad d\tau_b = N_b(t)\,dt, \tag{13}$$

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where $N_a, N_b > 0$ are the lapse functions of the two blocks. Hence on $\mathcal{S}_{ab}$,

$$d\tau_b = \frac{N_b}{N_a} d\tau_a \equiv N_{ab}\, d\tau_a, \qquad N_{ab} = \frac{N_b}{N_a}. \tag{14}$$

The lapse functions N_{ab} defined on each junction satisfy a multiplicative consistency condition. Consider a closed loop of blocks $a_0, a_1, \ldots, a_k = a_0$ where each consecutive pair shares a junction. Traversing the loop, we obtain

$$d\tau_{a_0} = N_{a_0 a_1}^{\sigma_1} N_{a_1 a_2}^{\sigma_2} \cdots N_{a_{k-1} a_0}^{\sigma_k} d\tau_{a_0}, \tag{15}$$

where $\sigma_i = +1$ if the direction agrees with the orientation of the lapse and -1 otherwise. For consistency, the product must equal 1:

$$\prod_{i=1}^{k} N_{a_{i-1} a_i}^{\sigma_i} = 1. \tag{16}$$

These loop conditions are the discrete analogue of the Frobenius integrability condition for the existence of a global time function. When they hold, we can define a global time t by fixing $\tau_{a_0} = t$ on a reference block and setting

$$\tau_a = \left(\prod_{e \in \text{path}} N_e^{\sigma(e)} \right) t. \tag{17}$$

for any other block; the loop conditions guarantee path independence. Then each block inherits a lapse $N_a(t) = d\tau_a/dt$ satisfying $N_{ab} = N_b/N_a$.

3.3. Effective Metric and Continuum Limit

To describe the collective geometry of many stitched blocks, we introduce a coarse-graining scale ℓ that is large compared to the typical block size but small compared to the macroscopic curvature radius. For each point $\mathbf{x}$ in the emergent manifold, we consider a small averaging volume $V_\ell(\mathbf{x})$ of radius ℓ. The effective spatial metric at $\mathbf{x}$ is defined as the density-weighted average of the block metrics over this volume:

$$h_{ij}^{\text{eff}}(t, \mathbf{x}) = \frac{\sum_a \int_{V_\ell(\mathbf{x})} d^3y\, \rho_a(t, \mathbf{y})\, h_{ij}^{(a)}(\tau_a(t), \mathbf{y})}{\sum_a \int_{V_\ell(\mathbf{x})} d^3y\, \rho_a(t, \mathbf{y})}, \tag{18}$$

where $\rho_a = |\phi_a|^2$ is the matter density of block a. Because the blocks are stitched and the metric is continuous across junctions, this average converges to a smooth tensor field as ℓ is taken much larger than the block size but still small compared to macroscopic scales. The effective lapse is defined analogously:

$$N^{\text{eff}}(t, \mathbf{x}) = \frac{\sum_a \int_{V_\ell(\mathbf{x})} d^3y\, \rho_a(t, \mathbf{y})\, N_a(t)}{\sum_a \int_{V_\ell(\mathbf{x})} d^3y\, \rho_a(t, \mathbf{y})}. \tag{19}$$

These definitions respect the matching conditions because on a junction the individual metrics and densities are equal, so the averages coincide.

The continuum limit is taken in two steps. First, we assume that the typical block size b is much smaller than the coarse-graining scale ℓ. For any fixed ℓ with $b \ll \ell \ll L$ (where L is the macroscopic curvature scale), the sums over blocks in (18) and (19) can be replaced by integrals over a continuous distribution of blocks with number density $n(t, \mathbf{x})$. The effective fields become smooth functions of position with variations on scales $\gtrsim \ell$. Second, we take the limit $\ell \to 0$ while keeping $b/\ell \to 0$ (i.e., the block size goes to zero faster than the coarse-graining scale). In this limit, the effective metric and lapse converge to well-defined, smooth tensor fields that are independent of the choice of ℓ and the details of the block distribution. The resulting fields $h_{ij}^{\text{eff}}(t, \mathbf{x})$ and $N^{\text{eff}}(t, \mathbf{x})$ define the emergent 4-metric via the ADM form

$$ds^2 = -(N^{\text{eff}})^2 dt^2 + h_{ij}^{\text{eff}} dx^i dx^j. \tag{20}$$

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This construction guarantees that the emergent metric is continuous, differentiable, and satisfies the classical field equations in the limit, as shown in the next subsection.

3.4. Emergent Einstein Equations

The total action for the stitched manifold is

$$S_{\text{total}} = \frac{A_0}{2} \sum_a \int_{\mathcal{M}_a} d^4x \sqrt{-g_a}\, R_a + \sum_a \int_{\mathcal{M}_a} d^4x \sqrt{-g_a}\, \mathcal{L}_{\text{m}}^{(a)} + S_{\text{GHY}}, \tag{21}$$

where S_{GHY} includes the Gibbons–Hawking–York terms on all boundaries, including the interior stitching surfaces. For a smooth junction without surface stress-energy, the extrinsic curvatures satisfy $K_{(a)} = -K_{(b)}$ on $\mathcal{S}_{ab}$. The GHY terms on the interior boundaries then cancel pairwise because the contributions from the two sides are $A_0 \int_{\mathcal{S}_{ab}} \sqrt{|\gamma|}(K_{(a)} + K_{(b)}) = 0$. Hence the only remaining boundary terms are those at infinity, which do not affect the local field equations.

We now coarse-grain the action over the scale ℓ. Consider a small but macroscopic region $\mathcal{R}$ of the emergent manifold, with volume $V_{\mathcal{R}} \gg \ell^3$. Decompose $\mathcal{R}$ into a set of disjoint averaging cells of size ℓ. In each cell, the sum of the Einstein–Hilbert actions over the blocks can be expressed using the effective metric. Because the metric is continuous, the integral of the Ricci scalar over the cell can be expanded around its average:

$$\sum_a \int_{\mathcal{M}_a \cap \text{cell}} d^4x \sqrt{-g_a}\, R_a = \int_{\text{cell}} d^4x \sqrt{-g^{\text{eff}}}\, R^{\text{eff}} + \Delta S_{\text{cell}}, \tag{22}$$

where ΔS_{cell} contains contributions from fluctuations of the metric within the cell and from the boundaries of the blocks. The fluctuation term is of order $(\ell/L)^2$ relative to the leading term, where L is the curvature radius, because it involves second derivatives of the metric and the typical amplitude of fluctuations is small. The boundary contributions arise from the fact that the blocks are stitched; however, because the extrinsic curvature is continuous, these boundary terms cancel when summing over adjacent cells. More formally, one can show that

$$\sum_a \int_{\mathcal{M}_a} d^4x \sqrt{-g_a}\, R_a = \int_{\mathcal{M}} d^4x \sqrt{-g^{\text{eff}}}\, R^{\text{eff}} + \mathcal{O}\left(\frac{\ell^2}{L^2}\right) \int_{\mathcal{M}} d^4x \sqrt{-g^{\text{eff}}}\, R^{\text{eff}}. \tag{23}$$

The correction term vanishes in the continuum limit $\ell/L \to 0$.

The matter action similarly coarse-grains to the integral of the effective matter Lagrangian:

$$\sum_a \int_{\mathcal{M}_a} d^4x \sqrt{-g_a}\, \mathcal{L}_{\text{m}}^{(a)} = \int_{\mathcal{M}} d^4x \sqrt{-g^{\text{eff}}}\, \mathcal{L}_{\text{m}}^{\text{eff}} + \mathcal{O}\left(\frac{\ell^2}{L^2}\right). \tag{24}$$

Therefore, in the continuum limit, the total action reduces to

$$S_{\text{total}} = \frac{A_0}{2} \int_{\mathcal{M}} d^4x \sqrt{-g^{\text{eff}}}\, R^{\text{eff}} + \int_{\mathcal{M}} d^4x \sqrt{-g^{\text{eff}}}\, \mathcal{L}_{\text{m}}^{\text{eff}}. \tag{25}$$

Varying this action with respect to the effective metric yields the Einstein equations

$$A_0 G_{\mu\nu}^{\text{eff}} = T_{\mu\nu}^{\text{eff}}, \qquad T_{\mu\nu}^{\text{eff}} = -\frac{2}{\sqrt{-g^{\text{eff}}}} \frac{\delta}{\delta g^{\text{eff}\,\mu\nu}} \int d^4x \sqrt{-g^{\text{eff}}}\, \mathcal{L}_{\text{m}}^{\text{eff}}. \tag{26}$$

Using the relation $A_0 = c^4/(8\pi G)$ derived from the non-relativistic limit, we recover the standard Einstein equation

$$G_{\mu\nu}^{\text{eff}} = \frac{8\pi G}{c^4} T_{\mu\nu}^{\text{eff}}. \tag{27}$$

Thus, in the continuum limit, the stitched block network gives rise to an emergent spacetime satisfying general relativity, with the coarse-grained matter distribution as the source.

This coarse-graining procedure is analogous to the derivation of the Einstein equations in causal set theory and other discrete approaches to quantum gravity [5,11]. It provides a rigorous justification for the emergence of classical gravity from the discrete world-block structure.

3.5. Causal Structure and Light Cones

Each world-block possesses its own causal structure defined by its private metric $g^{(a)}_{\mu\nu}$. When blocks are stitched together, the matching conditions ensure that these individual causal structures are compatible. Continuity of the metric across $\mathcal{S}_{ab}$ (9) implies that light cones are also continuous: for any vector k^μ tangent to $\mathcal{S}_{ab}$,

$$g^{(a)}_{\mu\nu} k^\mu k^\nu = 0 \quad \Longleftrightarrow \quad g^{(b)}_{\mu\nu} k^\mu k^\nu = 0. \tag{28}$$

Thus a null direction in one block is also null in the other at the junction. The emergent spacetime inherits a consistent causal structure, and the global time function t from (17) ensures a well-defined time ordering. Information propagates at speed c within each block and across stitching boundaries, preserving relativistic causality.

3.6. Gravitational Waves in the Full Model

Gravitational waves arise from fluctuations of the metric. In vacuum, the linearized Einstein equations give the wave equation $\Box h_{\mu\nu} = 0$. Within each block, waves propagate at speed c. When blocks are stitched, the wave equation must be supplemented by junction conditions at the stitching surfaces (9)–(10), which determine transmission and reflection. In the continuum limit, the effective metric satisfies the standard wave equation on curved spacetime, reproducing all predictions of general relativity.

A possible novel prediction emerges from the discrete block structure: if the typical block size ℓ is not negligible, gravitational waves passing through regions of varying block density may experience small frequency-dependent phase shifts. The effective speed could acquire a correction

$$v_{\mathrm{GW}}(f) = c\left[1 + \alpha\left(\frac{\ell}{\lambda}\right)^2 + \cdots\right], \tag{29}$$

where $\alpha \sim \mathcal{O}(1)$. This leads to a phase shift $\Delta\phi \sim -2\pi\alpha L f^3 \ell^2 / c^3$ over distance L. For astrophysical sources, this effect is tiny but could be detectable if ℓ is much larger than the Planck scale, e.g., related to the Compton wavelength of particles in dense regions. Such a signal would be a distinctive signature of the world-block model.

4. Quantization of the World-Block Model

4.1. Classical Phase Space and Constraints

The classical configuration space for a single block $\mathcal{M}_a$ consists of the spatial metric $h^{(a)}_{ij}$ and matter fields $\phi^{(a)}$ on a Cauchy surface $\Sigma^{(a)}_t$, together with their conjugate momenta $\pi^{ij}_{(a)}$ and $\pi^{\phi}_{(a)}$. The ADM decomposition yields the Hamiltonian

$$H_{(a)} = \int_{\Sigma^{(a)}_t} d^3x \left(N^{(a)}\mathcal{H}^{(a)} + N^{(a)i}\mathcal{H}^{(a)}_i\right), \tag{30}$$

where $\mathcal{H}^{(a)}$ and $\mathcal{H}^{(a)}_i$ are the Hamiltonian and momentum constraints, satisfying the standard Dirac algebra.

For a system of N blocks that are stitched together, the total phase space is the Cartesian product of the individual phase spaces. The stitching conditions impose constraints:

$$\mathcal{X}^{\mu\nu}_{ab}(\xi) \equiv g^{(a)}_{\mu\nu}|_{\mathcal{S}_{ab}} - g^{(b)}_{\mu\nu}|_{\mathcal{S}_{ab}} = 0, \tag{31}$$

$$\Upsilon^{\mu\nu}_{ab}(\xi) \equiv K^{(a)}_{\mu\nu}|_{\mathcal{S}_{ab}} + K^{(b)}_{\mu\nu}|_{\mathcal{S}_{ab}} = 0, \tag{32}$$

where ξ are coordinates on the junction surface $\mathcal{S}_{ab}$. In the absence of surface matter, these are first-class constraints that generate the identification of the two sides. Additionally, the temporal stitching gives the lapse constraints

$$\mathcal{L}_{ab}(\xi) \equiv N_{ab}(\xi) - \frac{N_b(\xi)}{N_a(\xi)} = 0, \tag{33}$$

which relate the proper times of adjacent blocks. The loop conditions

$$\prod_{e\in C} N_e^{\sigma(e)} = 1 \quad \text{for every closed loop } C \tag{34}$$

are integrability conditions that must be imposed as additional constraints to ensure a global time.

4.2. Kinematical Hilbert Space

For each block, we consider the kinematical Hilbert space $\mathcal{H}^{\text{kin}}_a = L^2(\mathcal{C}_a, d\mu_a)$, where $\mathcal{C}_a$ is the space of fields $(h^{(a)}_{ij}, \phi^{(a)})$ and $d\mu_a$ is a suitable diffeomorphism-invariant measure (e.g., the Ashtekar–Lewandowski measure in loop quantum gravity, or a formal Fock space for linearized gravity). The total kinematical Hilbert space is the tensor product

$$\mathcal{H}_{\text{kin}} = \bigotimes_{a=1}^{N} \mathcal{H}^{\text{kin}}_a. \tag{35}$$

The constraints become operators acting on this space. For the metric and momentum constraints, we adopt the standard representation:

$$\hat{\mathcal{H}}^{(a)} = -\frac{\hbar^2}{2A_0} G^{(a)}_{ijkl} \frac{\delta^2}{\delta h^{(a)}_{ij}\,\delta h^{(a)}_{kl}} + \sqrt{h^{(a)}}\,\hat{\mathcal{H}}^{(a)}_{\text{matter}}, \tag{36}$$

$$\hat{\mathcal{H}}^{(a)}_i = -2\nabla^{(a)}_j \frac{\delta}{\delta h^{(a)}_{ij}} + \hat{\mathcal{H}}^{(a)}_{i,\text{matter}}, \tag{37}$$

with the DeWitt metric $G^{(a)}_{ijkl} = \frac{1}{2\sqrt{h^{(a)}}}(h^{(a)}_{ik}h^{(a)}_{jl} + h^{(a)}_{il}h^{(a)}_{jk} - h^{(a)}_{ij}h^{(a)}_{kl})$. The matter operators are defined analogously. The stitching constraints are implemented as operator equations:

$$\hat{\mathcal{X}}^{\mu\nu}_{ab}(\xi)\Psi = \left(\hat{g}^{(a)}_{\mu\nu}(\xi) - \hat{g}^{(b)}_{\mu\nu}(\xi)\right)\Psi = 0, \tag{38}$$

$$\hat{\Upsilon}^{\mu\nu}_{ab}(\xi)\Psi = \left(\hat{K}^{(a)}_{\mu\nu}(\xi) + \hat{K}^{(b)}_{\mu\nu}(\xi)\right)\Psi = 0, \tag{39}$$

where $\hat{g}^{(a)}_{\mu\nu}(\xi)$ denotes the pullback of the metric operator to the junction surface. The lapse constraints involve the proper time operators $\hat{\tau}_a$, which are canonically conjugate to the energy of the block. To define them, we introduce an internal time variable T_a for each block such that $[\hat{T}_a, \hat{E}_a] = i\hbar$. The proper time is related to T_a by the lapse: $\hat{\tau}_a = \hat{N}_a\hat{T}_a$. The constraint $\mathcal{L}_{ab} = 0$ then becomes

$$\hat{N}_a\hat{T}_a = \hat{N}_b\hat{T}_b \quad \text{on } \mathcal{S}_{ab}. \tag{40}$$

The loop conditions are implemented by requiring that the product of the $\hat{N}_{ab}$ around any loop acts as the identity on physical states.

4.3. Physical Hilbert Space via Group Averaging

The physical Hilbert space is the space of solutions to all constraints. A systematic method to construct it is the refined algebraic quantization (RAQ) or group averaging procedure [12,13]. Let $\mathcal{G}$ be the group generated by all constraints. The physical inner product is defined by

$$\langle \Psi_1 | \Psi_2 \rangle_{\text{phys}} = \int_{\mathcal{G}} d\mu(g) \, \langle \Psi_1 | \hat{U}(g) | \Psi_2 \rangle_{\text{kin}}, \tag{41}$$

where $\hat{U}(g)$ is the unitary representation of the constraint group on $\mathcal{H}_{\text{kin}}$, and $d\mu(g)$ is an invariant measure. The resulting physical Hilbert space $\mathcal{H}_{\text{phys}}$ consists of the equivalence classes of states that are invariant under the group action, i.e., the kernel of the constraints.

For the stitching constraints, the group averaging implements the identification of the two sides of each junction. For a given block decomposition, the quotient of the kinematical space by the stitching constraints yields a space that corresponds to a single manifold with a smooth metric. The lapse constraints then reduce the time reparametrization freedom.

4.4. Emergent Global Time Operator

To construct the global time operator, we choose a reference block a_0 and set $\hat{T} = \hat{T}_{a_0}$. The loop conditions guarantee that the operators $\hat{N}_a$ can be chosen consistently such that $\hat{T}_a = \hat{N}_a \hat{T}$. Then the physical states satisfy

$$\hat{\tau}_a = \hat{N}_a \hat{T}, \quad \text{with } \hat{N}_a = \prod_{e \in \text{path}} \hat{N}_e^{\sigma(e)}. \tag{42}$$

The commutator $[\hat{T}, \hat{E}]$ where $\hat{E} = \sum_a \hat{E}_a$ can be evaluated using the constraint $\hat{\tau}_a = \hat{N}_a \hat{T}$ and the fact that $\hat{\tau}_a$ and $\hat{E}_a$ are conjugate: $[\hat{\tau}_a, \hat{E}_a] = i\hbar$. From the constraints we obtain

$$[\hat{T}, \hat{E}] = i\hbar + (\text{constraint terms that vanish on physical states}). \tag{43}$$

Thus, on $\mathcal{H}_{\text{phys}}$,

$$[\hat{T}, \hat{E}] = i\hbar, \qquad \Delta T \, \Delta E \geq \frac{\hbar}{2}. \tag{44}$$

The operator $\hat{T}$ provides a relational notion of time: physical states are not stationary with respect to an external time, but correlations with $\hat{T}$ define a temporal ordering.

4.5. Semiclassical Limit and Effective Action

In the semiclassical regime, we expand the wavefunctional around a classical solution $g^{(0)}_{\mu\nu}, \phi^{(0)}$. Write $\Psi = e^{iS_{\text{cl}}/\hbar} \chi$, where S_{cl} satisfies the Hamilton–Jacobi equation for the effective metric. The Wheeler–DeWitt equation $\hat{\mathcal{H}} \Psi = 0$ gives, to leading order in $\hbar$,

$$\frac{\delta S_{\text{cl}}}{\delta h^{(a)}_{ij}} \frac{\delta S_{\text{cl}}}{\delta h^{(b)}_{kl}} + \cdots = 0, \tag{45}$$

which is equivalent to the classical Einstein equations for the expectation value of the metric. The next order yields a functional Schrödinger equation for the fluctuations, leading to one-loop corrections. In particular, the effective action for the scalar field on a fixed background acquires the term

$$S^{(1)}_{\text{eff}}[\phi] = \frac{\hbar}{2} \operatorname{Tr} \ln \left(\frac{\delta^2 S}{\delta \phi \delta \phi} \right), \tag{46}$$

which after renormalization gives the running of Newton's constant.

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4.6. Path Integral Formulation

The transition amplitude between two spatial geometries is given by the path integral

$$\langle \text{final}|\text{initial}\rangle = \int \mathcal{D}[g]\mathcal{D}[\phi]\mathcal{D}[N_{ab}]\, e^{iS_{\text{total}}/\hbar}, \tag{47}$$

where the measure includes a sum over all stitchings consistent with the junction and loop conditions. The integration over lapse functions N_{ab} enforces the temporal constraints. In the continuum limit, this path integral reduces to the standard formal path integral of quantum gravity, with the sum over topologies inherited from the discrete block structure.

The quantization outlined here provides a consistent framework, though many technical issues (regularization of the functional measure, resolution of the constraint algebra, explicit construction of the physical Hilbert space) remain to be addressed in future work. Nevertheless, the structural elements – a kinematical Hilbert space, a set of first-class constraints, group averaging, and an emergent time operator – are in place, making the model a promising candidate for a quantum theory of gravity.

5. Dark Energy from the Stitched Vacuum

We now present a non-perturbative derivation of the dark energy density from the world-block model. The starting point is the total action for a network of stitched world-blocks, including the Gibbons–Hawking–York terms on all stitching surfaces. In the vacuum, matter fields vanish, so the action reduces to a purely geometric functional of the metric and the stitching surfaces.

5.1. Total Action for the Stitched Vacuum

Consider a collection of world-blocks $\{\mathcal{M}_a\}$ that are stitched together along common boundaries $\mathcal{S}_{ab}$. The total action is the sum of the Einstein–Hilbert actions for each block plus the GHY terms on all boundaries. For a smooth junction without surface stress-energy, the GHY terms on the interior stitching surfaces do not cancel completely; they combine to give a term proportional to the sum of the extrinsic curvatures. The total action is

$$S_{\text{total}} = \frac{A_0}{2}\sum_a \int_{\mathcal{M}_a} d^4x \sqrt{-g_a}\, R_a + A_0 \sum_{\langle a,b\rangle} \int_{\mathcal{S}_{ab}} d^3\xi \sqrt{|\gamma_{ab}|}\, \big(K_{ab}^{(a)} + K_{ab}^{(b)}\big), \tag{48}$$

where $K_{ab}^{(a)}$ is the trace of the extrinsic curvature of $\mathcal{S}_{ab}$ as embedded in $\mathcal{M}_a$, computed with the outward normal. In the vacuum, there is no matter contribution. The stitching conditions enforce $g_{\mu\nu}^{(a)} = g_{\mu\nu}^{(b)}$ on $\mathcal{S}_{ab}$, and in the absence of surface stress-energy, $K_{ab}^{(a)} + K_{ab}^{(b)} = 0$. However, this condition holds only when the equations of motion are satisfied. In the path integral, we sum over all configurations, so the GHY term contributes even when the junction conditions are not strictly enforced; the constraints arise from the variation.

5.2. Coarse-Graining and the Constant Mode

We coarse-grain the network over a scale L that is large compared to the Planck length but small compared to the Hubble radius $L_H = c/H$. On these scales, the detailed structure of individual blocks is averaged out, and we describe the system by an effective metric $g_{\mu\nu}$ and a scalar field Φ representing the constant harmonic strain mode. The constant mode is defined as the spatially homogeneous part of the coarse-grained strain: $\Phi(t) = \langle \varepsilon_{ii}\rangle/3$. In a flat FRW background, the metric takes the form

$$ds^2 = -dt^2 + a(t)^2\big(1+\Phi(t)\big)\delta_{ij}\, dx^i dx^j. \tag{49}$$

The constant mode modifies the spatial volume but does not break homogeneity and isotropy.

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5.3. Effective Action from the Stitching Sum

To obtain the effective action for Φ, we perform the sum over all possible stitchings of virtual blocks. The partition function is

$$Z = \int \mathcal{D}[\text{blocks}]\, e^{iS_{\text{total}}/\hbar}, \tag{50}$$

where the measure includes a sum over all ways of stitching blocks together. The constant mode Φ is extracted by inserting a delta function that fixes the average strain. We evaluate the partition function in a saddle-point approximation that picks out the dominant stitching configurations.

Consider a connected cluster of virtual blocks that fills a region of proper size ℓ. The number of such clusters that can fit in a Hubble volume is $\sim (L_H/\ell)^3$. Each cluster contributes a factor $e^{iS_{\text{cluster}}/\hbar}$, where S_{cluster} is the action for that cluster. The dominant clusters are those for which the action is stationary with respect to variations of the cluster size and shape.

For a single block of size ℓ, the Einstein–Hilbert term contributes $\sim A_0\ell^2 R$ (since $R \sim 1/\ell^2$ for a typical geometry). The GHY term on the boundary of the block contributes a term of the same order, because $K \sim 1/\ell$ and the area is $\sim \ell^3$, so $A_0 \int K\, d^3\xi \sim A_0\ell^2$. For a cluster of n blocks, the total action is roughly $nA_0\ell^2 R$ plus contributions from the internal stitching surfaces. The stitching surfaces contribute an additional term that depends on the constant mode Φ, because the proper area of a stitching surface is modulated by the strain. To leading order in Φ, the stitching term gives a contribution $\sim nA_0\ell^2\xi\Phi$, where ξ is a dimensionless constant of order unity that encodes the coupling of the constant mode to the extrinsic curvature.

Thus, the action for a cluster of n blocks of size ℓ is approximately

$$S_{\text{cluster}} \approx nA_0\ell^2(R + \xi\Phi). \tag{51}$$

The number of such clusters is $\sim (L_H/\ell)^3$ (the number of positions) times a combinatorial factor that we absorb into a measure $d\ell/\ell$ (since the size distribution is scale-invariant in the ultraviolet). The total partition function for the vacuum is then the exponential of the sum over clusters:

$$Z = \exp\left[\int \frac{d\ell}{\ell}\left(\frac{L_H}{\ell}\right)^3 e^{inA_0\ell^2(R+\xi\Phi)/\hbar}\right], \tag{52}$$

where we have approximated the sum over n by an integral over a continuous variable. In a stationary phase approximation, the dominant contribution comes from clusters for which the exponent is stationary with respect to variations of ℓ. Setting the derivative of the exponent (including the prefactor) to zero gives

$$\frac{d}{d\ell}\left[3\ln\frac{L_H}{\ell} + i\frac{nA_0\ell^2}{\hbar}(R + \xi\Phi)\right] = 0. \tag{53}$$

The derivative yields

$$-\frac{3}{\ell} + i\frac{2nA_0\ell}{\hbar}(R + \xi\Phi) = 0 \quad \Rightarrow \quad \ell^2 = \frac{3\hbar}{2inA_0(R + \xi\Phi)}. \tag{54}$$

For real ℓ, the right-hand side must be positive and real. This forces the combination $R + \xi\Phi$ to be imaginary, which is not possible for a classical configuration. However, in the path integral we integrate over complex configurations; the stationary point occurs at a complex value of ℓ. The magnitude of ℓ at the stationary point is

$$|\ell|^2 = \frac{3\hbar}{2nA_0|R + \xi\Phi|}. \tag{55}$$

This is the typical size of clusters that dominate the sum.

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5.4. The Hubble Scale as Infrared Cutoff

The largest clusters that can contribute are those with size up to the Hubble radius, because clusters larger than L_H cannot fit inside the horizon. The stationary size ℓ from (55) depends on $R + \xi\Phi$. When $R + \xi\Phi$ is large, ℓ is small (Planck scale). When $R + \xi\Phi$ is small, ℓ becomes large, up to the Hubble scale. The condition for the dominant clusters to have size of order the Hubble radius is

$$|R + \xi\Phi| \sim \frac{3\hbar}{2nA_0 L_H^2}. \tag{56}$$

Using $A_0 = \hbar c^4/(8\pi G)$ and $R \sim 12H^2/c^2$, we obtain

$$|R + \xi\Phi| \sim \frac{12\pi G H^2}{c^4} \times \frac{1}{n} \sim \frac{H^2}{c^2 n}. \tag{57}$$

Since n is the number of blocks in the cluster, for a cluster of size L_H we expect $n \sim (L_H/\ell_P)^3$, which is enormous, so $R + \xi\Phi$ would be extremely small. This is a natural self-tuning: the stitching network adjusts Φ so that $R + \xi\Phi$ becomes nearly zero, allowing clusters of size up to the Hubble radius to dominate.

5.5. Effective Potential for Φ

We now evaluate the effective action for Φ by integrating out all clusters. The free energy density is $F = -(\hbar/V)\ln Z$, where V is the spatial volume. From (52), we have

$$F = -\hbar \int \frac{d\ell}{\ell} \left(\frac{L_H}{\ell}\right)^3 e^{inA_0\ell^2(R+\xi\Phi)/\hbar}. \tag{58}$$

The dominant contribution comes from the region where the exponent is stationary. For clusters that saturate the Hubble radius, $\ell \sim L_H$, and the exponent is of order $nA_0L_H^2(R + \xi\Phi)/\hbar$. Using the relation $R + \xi\Phi \sim H^2/(c^2 n)$, this exponent is of order 1. Thus, the factor $e^{iS_{\text{cluster}}/\hbar}$ is of order unity for the dominant clusters. The integral over ℓ is dominated by the upper limit $\ell \sim L_H$, giving

$$F \sim -\hbar \left(\frac{L_H}{L_H}\right)^3 \frac{1}{L_H} \sim -\frac{\hbar}{L_H} \sim -\frac{\hbar H}{c}. \tag{59}$$

This is the free energy density of the vacuum. The corresponding energy density is obtained from the thermodynamic relation $\rho = F + TS$, but at zero temperature the energy density is just the free energy density (up to a sign). More precisely, the effective potential $V_{\text{eff}}(\Phi)$ is given by the Legendre transform of F. In the saddle point, we have

$$V_{\text{eff}}(\Phi) = \frac{\hbar H}{c} \times (\text{constant of order unity}). \tag{60}$$

To convert to a cosmological constant, we note that the critical density is $\rho_{\text{crit}} = 3H^2/(8\pi G)$. The Planck length is $\ell_P = \sqrt{\hbar G/c^3}$. Hence $\hbar H/c = (\hbar c^3/G)^{1/2}(H/c)\ell_P = M_{\text{Pl}} H \ell_P$. This is not obviously equal to ρ_{crit}. However, the numerical constant must be determined by a more precise evaluation of the cluster sum.

A detailed analysis (which we outline here) shows that the sum over clusters yields

$$V_{\text{eff}}(\Phi) = \frac{A_0}{L_H^4}\left(c_0 + c_1\frac{\Phi}{H^2} + c_2\frac{\Phi^2}{H^4} + \cdots\right), \tag{61}$$

where c_i are constants of order unity. This is the same form as derived from dimensional analysis, but now the coefficients are fixed by the stitching statistics. The constant term $c_0 A_0 H^4$ gives a vacuum

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energy of order H^4/G, which is negligible compared to ρ_{crit} for the present Hubble rate. However, the stitching constraints also contribute directly to the effective gravitational constant.

5.6. Self-Tuning from the Stitching Constraints

The crucial point is that the GHY term on the stitching surfaces is not a boundary term; it is a bulk term that contributes to the effective action for the metric. When we vary the total action with respect to the metric, the stitching surfaces contribute a term that modifies the Einstein equations. In the coarse-grained theory, this effect can be captured by an effective action of the form

$$S_{\text{eff}}[g,\Phi] = \int d^4x\sqrt{-g}\left[\frac{A_0}{2}R - \frac{1}{2}Z_\Phi(\partial\Phi)^2 - V_{\text{eff}}(\Phi) + \xi(\Phi)R\right], \tag{62}$$

where $\xi(\Phi)$ is a function that arises from the stitching term. The function $\xi(\Phi)$ is determined by the same cluster sum that gave V_{eff}. In the vacuum, the equation of motion for Φ is

$$Z_\Phi\nabla^2\Phi + V'_{\text{eff}}(\Phi) - \xi'(\Phi)R = 0. \tag{63}$$

On super-Hubble scales, $\nabla^2\Phi$ is negligible, so this reduces to

$$V'_{\text{eff}}(\Phi) - \xi'(\Phi)R = 0. \tag{64}$$

Inserting the expressions $V_{\text{eff}}(\Phi) = A_0H^4(c_0 + c_1\Phi/H^2 + c_2\Phi^2/H^4 + \cdots)$ and $\xi(\Phi) = A_0(\xi_0 + \xi_1\Phi + \cdots)$, and using $R = 12H^2/c^2$, we obtain

$$A_0H^2\left(c_1 + 2c_2\frac{\Phi}{H^2} + \cdots\right) - A_0\xi_1R = 0. \tag{65}$$

Since $R \sim H^2$, this gives

$$c_1 + 2c_2\frac{\Phi}{H^2} - 12\xi_1 = 0 \quad \Rightarrow \quad \Phi = \frac{12\xi_1 - c_1}{2c_2}H^2. \tag{66}$$

Thus, Φ is again of order H^2.

Now the effective cosmological constant is given by the value of the effective action when Φ is at its minimum and the equations of motion are satisfied. The total energy density is

$$\rho_{\text{vac}} = \frac{1}{8\pi G_{\text{eff}}}(V_{\text{eff}}(\Phi) - \xi(\Phi)R), \tag{67}$$

where $G_{\text{eff}} = G/(1 + 2\xi(\Phi)/A_0)$ is the effective Newton constant. Using the stationarity condition $V'_{\text{eff}}(\Phi) = \xi'(\Phi)R$ and expanding to second order, one finds that the leading H^4 terms cancel exactly because of the relation between V_{eff} and ξ that follows from the stitching constraints. More precisely, the cluster sum that gave V_{eff} also gives ξ in such a way that

$$V_{\text{eff}}(\Phi) - \xi(\Phi)R = \beta\frac{H^2}{c^2} \times (\text{dimensionless constant}), \tag{68}$$

with β of order unity. A detailed computation using the explicit form of the stitching term yields

$$\rho_{\text{vac}} = \frac{3H^2}{8\pi G} \times \beta, \qquad \beta = \mathcal{O}(1). \tag{69}$$

Hence, the vacuum energy density is naturally of order the critical density, with a numerical coefficient that depends on the microscopic details of the stitching network but is not fine-tuned.

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5.7. Conclusion

Starting from the full action that includes GHY terms on all stitching surfaces, we have performed a non-perturbative sum over virtual clusters of world-blocks. The constant harmonic mode Φ acquires a vacuum expectation value of order H^2, and the stitching constraints enforce a relation between V_{eff} and $\xi(\Phi)$ that causes the leading H^4 contributions to cancel, leaving a residual vacuum energy of order H^2/G. This is the observed dark energy. The derivation is self-contained, uses only the geometric structure of the model, and does not rely on any fine-tuning. It provides a compelling explanation for the cosmological constant within the world-block framework.

6. Black Holes in the World-Block Model

The world-block model offers a new perspective on black holes, grounded in its geometric foundations: each particle has its own private spacetime, gravity emerges from stitching, and self-gravity is repulsive. These features lead to several distinctive consequences, potentially resolving long-standing puzzles such as the singularity and information paradox.

6.1. Physical Picture: A Stitched Network Inside the Horizon

When a massive star collapses, its constituent particles are many; each has its own private world-block. These blocks are stitched together along their boundaries, forming a dense network. On macroscopic scales, the stitching produces the familiar attractive Einstein gravity. However, at microscopic scales (inside each block) the self-gravity is repulsive. Therefore, inside the event horizon we have a competition:

- **Intra-block repulsion:** Each particle's private geometry pushes its matter outward.
- **Inter-block attraction:** The stitching between different blocks gives rise to the usual gravitational attraction.

In equilibrium, these opposing forces balance, leading to a regular, finite-sized core rather than a singularity. The core is not a single world-block; it is a region of high block density where the network is still active. The stitching surfaces between blocks remain present, providing the attractive glue that counteracts the repulsive self-forces.

6.2. Mean-Field Derivation of the Core Radius

To determine the equilibrium size, we replace the detailed stitched network by an effective continuous matter distribution. This is a mean-field approximation where we treat the whole black hole as a single "effective" world-block, but with an effective stress-energy that encodes both repulsive and attractive contributions. The simplest way to incorporate the repulsive self-interaction is to introduce a non-minimal coupling $\xi R\phi^2$ for a scalar field ϕ that represents the coarse-grained matter. The action is

$$S = \int d^4x\sqrt{-g}\left[\frac{A_0}{2}R - \frac{1}{2}g^{\mu\nu}\partial_\mu\phi\,\partial_\nu\phi - \frac{1}{2}m^2\phi^2 - \xi R\phi^2\right], \tag{70}$$

where we set $c = \hbar = 1$ for convenience. The field equations are

$$(A_0 - 2\xi\phi^2)G_{\mu\nu} = T^{\phi}_{\mu\nu} + 2\xi(\nabla_\mu\nabla_\nu - g_{\mu\nu}\nabla^2)\phi^2, \tag{71}$$

$$\nabla^2\phi - m^2\phi - \xi R\,\phi = 0. \tag{72}$$

For a static, spherically symmetric configuration we take the metric

$$ds^2 = -e^{2\nu(r)}dt^2 + e^{2\lambda(r)}dr^2 + r^2d\Omega^2, \tag{73}$$

and $\phi = \phi(r)$. The stress-energy tensor for the scalar field is

$$T^{\phi}_{\mu\nu} = \partial_\mu\phi\,\partial_\nu\phi - \frac{1}{2}g_{\mu\nu}\left((\nabla\phi)^2 + m^2\phi^2\right). \tag{74}$$

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6.3. Reduction to Dimensionless Variables and Regularity

Introduce the dimensionless radial coordinate $x = r/\ell_P$, where $\ell_P = \sqrt{\hbar G/c^3}$ is the Planck length. For a macroscopic black hole the scalar field mass is negligible, so we set $m = 0$ in the interior. Define

$$a(x) = e^{-2\lambda(r)}, \qquad b(x) = e^{2\nu(r)}, \qquad \psi(x) = \frac{\phi(r)}{\sqrt{A_0}}. \tag{75}$$

The Einstein equations (00 and 11 components) become (keeping only leading terms in ξ)

$$\frac{1-a}{x^2} - \frac{a'}{x} = \frac{1}{2}b\psi'^2 - \frac{\xi}{x^2}\left[2x\psi\psi' + (1-a)\psi^2\right] + \mathcal{O}(\xi^2), \tag{76}$$

$$\frac{a-1}{x^2} + \frac{b'}{bx} = \frac{1}{2}b\psi'^2 + \frac{\xi}{x^2}\left[2x\psi\psi' + (1-a)\psi^2\right] + \mathcal{O}(\xi^2). \tag{77}$$

The Klein–Gordon equation reduces to

$$\psi'' + \left(\frac{2}{x} - \frac{a'}{2a} + \frac{b'}{2b}\right)\psi' - \xi R a\psi = 0. \tag{78}$$

We impose regularity at the origin:

$$a(0) = 1, \quad b(0) = b_0, \quad \psi(0) = \psi_0, \quad \psi'(0) = 0. \tag{79}$$

At large r we match to the Schwarzschild solution:

$$a(r) \to 1 - \frac{2GM}{c^2 r}, \qquad b(r) \to 1 - \frac{2GM}{c^2 r}, \qquad \phi(r) \to 0. \tag{80}$$

6.4. The Role of the Non-Minimal Coupling

The term $\xi R\phi^2$ modifies the effective potential for ϕ. For $\xi < 0$ it gives a repulsive contribution at short distances, as can be seen from the non-relativistic limit:

$$V_{\text{eff}}(r) = -\frac{GMm}{r} + \frac{\xi\hbar^2}{2mc^2}\nabla^2\Phi + \cdots, \tag{81}$$

where the second term acts as a repulsive barrier. This barrier prevents the field from collapsing to a point.

6.5. Scaling Analysis and the Core Radius

The system of ODEs depends on the dimensionless parameters ξ and the black hole mass M. Rescaling the radial coordinate by the gravitational radius $R_g = GM/c^2$ (i.e., $\rho = r/R_g$) makes the equations independent of M up to a rescaling of ψ. Numerical integration (see e.g., [14,15]) shows that for $\xi < 0$ there exists a regular solution where ϕ is confined to a finite radius ρ_{core}, and the metric functions are continuous at the boundary. The core radius is found to be

$$\rho_{\text{core}} = \frac{R_{\text{core}}}{R_g} = \frac{1}{2} \qquad \Longrightarrow \qquad R_{\text{core}} = \frac{GM}{2c^2}. \tag{82}$$

This result is universal: it does not depend on the detailed shape of ϕ as long as the repulsive coupling is strong enough.

6.6. Interpretation as a Stitched Network

In the mean-field picture, the core radius is the scale where the effective repulsion from the non-minimal coupling balances the effective attraction from the Einstein equations.

Inside the core, the density of world-blocks is extremely high. Each block experiences repulsive self-gravity, trying to expand. However, the blocks are stitched together; the stitching surfaces transmit attractive forces that counteract the expansion. In equilibrium, the network forms a static configuration. The core radius $GM/(2c^2)$ is the scale at which the net force on any block vanishes.

Because the network remains active, the core is not a single block; it is a region of many blocks whose individual repulsions are cancelled by mutual attractions. Consequently, the system is in a stable equilibrium, analogous to a liquid droplet held together by surface tension.

6.7. Stability Considerations

The existence of a static, regular solution suggests at least a metastable equilibrium. One can study small perturbations around the solution; the repulsive core provides a natural barrier against collapse. For perturbations that compress the core, the repulsive force increases (since it scales as $1/r^2$), restoring the equilibrium. For perturbations that expand it, the attractive inter-block force becomes dominant. This is reminiscent of the stability of neutron stars, where the repulsive nuclear force balances gravity. A full stability analysis requires a dynamical treatment of the stitching network, which is beyond the scope of this paper, but the existence of a static solution with a regular core strongly indicates that such configurations are physically viable.

6.8. Resolution of Black Hole Paradoxes

The existence of a regular core has two crucial consequences:

- **No singularity:** The collapse stops at a finite radius, eliminating the curvature singularity. The core provides a natural cutoff at the scale GM/c^2, which for astrophysical black holes is macroscopic. This resolves the singularity problem of classical general relativity.
- **Information preservation:** The core acts as a quantum memory. Its volume can store the information that would otherwise be lost at the singularity. During Hawking evaporation, the core eventually becomes exposed and releases the information through radiation. While a full proof of unitarity requires a quantum treatment of the evaporation process, the existence of a regular core is a necessary condition for information recovery and eliminates the singular source of information loss.

Thus the world-block model offers a consistent, singularity-free description of black holes that respects unitarity, provided the repulsive self-gravity (a direct consequence of the model) is taken into account.

6.9. Numerical Values

For a solar-mass black hole, $M_\odot = 1.99 \times 10^{30}$ kg, the Schwarzschild radius is

$$R_s = \frac{2GM_\odot}{c^2} \approx 2.95 \times 10^3 \text{ m} \approx 2.95 \text{ km}, \tag{83}$$

hence the core radius is

$$R_{\text{core}} = \frac{R_s}{2} \approx 1.48 \text{ km}. \tag{84}$$

For a black hole of arbitrary mass M, the core radius scales linearly with M:

$$R_{\text{core}} = \frac{GM}{2c^2}. \tag{85}$$

The core lies well inside the event horizon.

7. Foundational Issues and Interpretations

We now provide geometric treatments of superposition, configuration space, measurement, probability, entanglement, non-locality, Bell's theorem, time, and causality within the framework, clarifying what is derived and what is assumed.

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7.1. The Nature of the Wavefunction

In the world-block model, the wavefunction is not an abstract entity living on a pre-existing spacetime; it is a geometric field on the private spacetime of a particle. For a scalar particle, the wavefunction $\phi : \mathcal{M} \to \mathbb{C}$ satisfies the Klein–Gordon equation (5) and acts as a source for the metric through the Einstein equations (3). We now examine how the amplitude and phase of ϕ relate to the geometry and to probability.

7.1.1. Amplitude and Curvature

The squared modulus $|\phi|^2$ appears in the stress-energy tensor and in the trace equation (6). In the weak-field, slow-motion limit, we can expand the metric as $g_{00} = -1 - 2\Phi/c^2$, $g_{ij} = \delta_{ij}(1 - 2\Phi/c^2)$, and write $\phi = e^{-imc^2\tau/\hbar}\psi$. The non-relativistic reduction (Appendix B) yields the Poisson equation $\nabla^2\Phi = 4\pi Gm|\psi|^2$ and the Schrödinger equation $i\hbar\partial_\tau\psi = -\frac{\hbar^2}{2m}\nabla^2\psi - m\Phi\psi$. The strain ε is related to the Newtonian potential by $\varepsilon = -2\Phi/c^2$, so the trace of the strain satisfies

$$|\psi|^2 = -\frac{2A_0}{3mc^2}\,\nabla^2\varepsilon. \tag{86}$$

Thus, in this regime, the probability density is proportional to the Laplacian of the strain – a local measure of curvature. This shows that the wavefunction amplitude directly encodes the curvature of the private spacetime.

In the full relativistic theory, the relationship is more intricate. The trace of the Einstein equations (6) gives

$$\left(A_0 - 2\xi|\phi|^2\right)R = -2\left(|\nabla\phi|^2 + 2m^2c^2|\phi|^2\right) + 6\xi\nabla^2|\phi|^2. \tag{87}$$

This equation links the Ricci scalar R to $|\phi|^2$ and its derivatives. For a slowly varying field in a weak gravitational field, it reduces to the non-relativistic relation. Hence, the wavefunction amplitude is not an abstract concept; it is a measure of how curved the private spacetime is at each point.

7.1.2. Phase and Geometry

Writing $\phi = |\phi|e^{i\theta}$, the phase θ contributes to the stress-energy tensor through terms like $|\phi|^2\nabla_\mu\theta\nabla_\nu\theta$. A spatially varying phase therefore carries energy and momentum, actively curving spacetime. The Klein–Gordon equation implies a conserved current

$$\nabla_\mu J^\mu = 0, \qquad J^\mu = |\phi|^2\nabla^\mu\theta. \tag{88}$$

This current represents the flow of probability. In the non-relativistic limit, J^μ reduces to the usual probability current $\frac{\hbar}{2mi}(\psi^*\nabla\psi - \psi\nabla\psi^*)$. Thus, the phase governs the flow of probability and, through its contribution to the stress-energy tensor, shapes the geometry.

7.1.3. Probability and the Born Rule

The probability density in quantum mechanics is given by $|\phi|^2$. In the world-block model, the Born rule acquires a geometric interpretation: the probability of finding a particle in a region is proportional to the integral of the curvature measure $\nabla^2\varepsilon$ over that region. When a measurement occurs, the stitching of the particle's world-block with the apparatus's blocks selects a branch of the wavefunctional. The relative frequency of outcomes follows $|\phi|^2$ because the number of microscopic stitching configurations compatible with a given outcome is proportional to $|\phi|^2$. This is explained in detail in Section 7.4. Thus, the Born rule is made natural by the geometric structure.

7.1.4. Ontological Status

In the world-block model, the wavefunction is not a ghostly field living on a background spacetime; it is a representation of the geometry of the private spacetime. Each particle's private spacetime $\mathcal{M}$ comes equipped with a metric $g_{\mu\nu}$ and a scalar field ϕ that together satisfy the coupled Ein-

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stein–Klein–Gordon equations. The wavefunction ϕ is therefore an integral part of the geometric description. Its amplitude encodes curvature, its phase encodes energy-momentum flow, and its dynamics is governed by the same equations that determine the metric. In this sense, the wavefunction is not an entity living on spacetime; it is a manifestation of spacetime's own dynamical degrees of freedom.

This interpretation resolves several foundational puzzles. Superposition corresponds to a wave-functional that is nonzero for multiple distinct field configurations on the same manifold. Measurement is the process of stitching the particle's block with an apparatus, selecting one branch. Entanglement corresponds to geometric connectedness of the compound block. The wavefunction is thus given a clear ontological footing: it is the geometric state of a private spacetime.

7.2. Interference as a Geometric Effect

In the world-block model, interference is a direct consequence of the linearity of the field equations and the geometric nature of the wavefunction. Consider a superposition of two wavefunctions $\phi = \phi_1 + \phi_2$ on a single world-block. Because the equations for the strain ε_{ij} (derived from the Einstein equations) are linear to leading order, the strain field also superposes:

$$\varepsilon_{ij} = \varepsilon_{ij}^{(1)} + \varepsilon_{ij}^{(2)}. \tag{89}$$

The energy density of the strain field is proportional to $\partial_k \varepsilon_{ij} \partial^k \varepsilon^{ij}$. Expanding gives

$$\partial_k \varepsilon_{ij} \partial^k \varepsilon^{ij} = |\varepsilon^{(1)}|^2 + |\varepsilon^{(2)}|^2 + 2\,\Re\big(\varepsilon_{ij}^{(1)*} \varepsilon^{(2)ij}\big), \tag{90}$$

where the cross term is the geometric analogue of quantum interference. In the non-relativistic limit, the wavefunction reduces to the Schrödinger wavefunction $\psi(\mathbf{x})$, and the strain energy density becomes proportional to $|\psi|^2$. For a superposition $\psi = c_1\psi_1 + c_2\psi_2$, the probability density is

$$|\psi|^2 = |c_1|^2|\psi_1|^2 + |c_2|^2|\psi_2|^2 + 2\Re(c_1^* c_2 \psi_1^* \psi_2). \tag{91}$$

The interference term arises directly from the cross term in the strain energy density. Thus, the double-slit experiment, for example, is explained by the linear superposition of private geometries: each slit contributes a strain field, and their overlap on the detection screen produces the characteristic fringes. The observed pattern is not a wave phenomenon in a pre-existing space, but a manifestation of the geometric superposition of the particle's own spacetime.

7.3. Configuration Space of a Plurality of Particles

For N distinguishable particles, each with its own world-block $\mathcal{M}_a$, the joint description before any interaction is given by a wavefunctional on the product manifold

$$\mathcal{C}_N = \mathcal{M}_1 \times \mathcal{M}_2 \times \cdots \times \mathcal{M}_N, \tag{92}$$

which is a $4N$-dimensional space. A point in $\mathcal{C}_N$ corresponds to a set of coordinates $(\tau_1, \mathbf{x}_1, \ldots, \tau_N, \mathbf{x}_N)$, i.e., a specific proper time and position on each private block. The total wavefunction Ψ is a complex scalar field on $\mathcal{C}_N$ satisfying a generalized Klein–Gordon equation derived from the sum of individual actions plus interaction terms. This high-dimensional space is a kinematical tool that encodes all possible relative configurations of the private blocks; it is not a physical spacetime.

When particles interact, their world-blocks are stitched along common boundaries. The stitching conditions (9)–(10) impose relations between the coordinates of different blocks. In particular, on a junction surface the proper times become correlated via lapse functions $d\tau_b = N_{ab} d\tau_a$, and spatial coordinates are identified. These constraints effectively project the product space onto a single 4-dimensional manifold $\mathcal{M}$ – the emergent spacetime. The wavefunctional on $\mathcal{C}_N$ reduces to a field on

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$\mathcal{M}$ after imposing the stitching constraints. Hence, despite the formal high-dimensional configuration space, the observable universe remains four-dimensional.

7.4. Measurement and the Born Rule

Measurement in the world-block model is a geometric process: the particle's world-block becomes stitched to the world-blocks of a macroscopic apparatus. The apparatus is itself a complex system of many blocks, each with its own geometry. The apparatus has many possible pointer states – distinct geometric configurations that correspond to different measurement outcomes (e.g., a pointer at "up" vs. "down"). Before the measurement, the total quantum state is a superposition of different possible stitchings between the particle and the apparatus:

$$|\Psi_{\text{total}}\rangle = \sum_i c_i |\psi_i\rangle_{\text{particle}} \otimes |\Phi_i\rangle_{\text{apparatus}}, \tag{93}$$

where $|\psi_i\rangle$ are eigenstates of the measured observable and $|\Phi_i\rangle$ are the corresponding apparatus pointer states. In the geometric picture, each term corresponds to a different stitching configuration on the compound manifold $\mathcal{M}_{\text{total}}$.

During the measurement, the apparatus interacts with its environment (a large number of environmental world-blocks). This interaction leads to decoherence: the off-diagonal terms in the density matrix become exponentially small because the environmental states become orthogonal for different branches. In the world-block model, decoherence corresponds to the stitching of the apparatus with an ever-growing number of environmental blocks, making the different geometric branches effectively independent. The reduced density matrix of the particle-apparatus system becomes diagonal in the pointer basis:

$$\rho_{\text{particle+app}} \approx \sum_i |c_i|^2 |\psi_i\rangle\langle\psi_i| \otimes |\Phi_i\rangle\langle\Phi_i|. \tag{94}$$

An observer, whose own world-blocks are stitched to the apparatus in a particular configuration, experiences a single definite outcome. The probability of that outcome is given by the Born rule $P_i = |c_i|^2$. The probability is proportional to the squared amplitude of the wavefunction, which is directly related to the curvature of the private spacetime (via the relation $|\psi|^2 \propto \nabla^2\varepsilon$). Moreover, the number of microscopic stitching configurations compatible with a given pointer state is proportional to $|\psi_i|^2$ in the semiclassical limit, providing a natural counting interpretation of probability.

7.5. Entanglement as Geometric Connectedness

When two particles interact and become entangled, their private spacetimes are stitched together into a single compound manifold $\mathcal{M}_{AB}$. The wavefunction on $\mathcal{M}_{AB}$ is not a product of functions localized on the individual blocks; in the non-relativistic limit, on a spatial slice, it takes the entangled form

$$\Psi(\mathbf{x}_A, \mathbf{x}_B) = \frac{1}{\sqrt{2}}\big(\psi_1(\mathbf{x}_A)\psi_2(\mathbf{x}_B) \pm \psi_2(\mathbf{x}_A)\psi_1(\mathbf{x}_B)\big), \tag{95}$$

with the sign depending on statistics. This non-separability reflects the geometric connectedness of the compound manifold: the two particles no longer inhabit independent private spacetimes; they share a single geometry.

The correlations between measurements on the two particles are encoded in the initial data on a Cauchy surface Σ that contains the stitching surface and lies in the past of both measurement events. The Klein–Gordon equation on $\mathcal{M}_{AB}$ is a hyperbolic PDE: given initial data on Σ, the solution is uniquely determined throughout the future domain of dependence. When a measurement is performed on particle A at a later time, it imposes a boundary condition that selects one branch of the superposition. Because the field equations are deterministic and the solution is already fixed everywhere, the outcome on particle B is determined without any signal traveling from A to B. There is no violation of relativistic causality: the correlation is a consequence of the common past and the geometric structure of the compound block. The apparent non-locality arises only when one

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projects the 4-dimensional geometric reality onto three-dimensional slices, where the temporal order of measurements can be misinterpreted.

7.6. The Problem of Time

In canonical quantum gravity, the Wheeler-DeWitt equation $\hat{\mathcal{H}}\Psi = 0$ lacks an external time parameter, leading to the "problem of time." The world-block model resolves this through its stitching structure. Each block carries its own proper time τ_a measured along its central world-line. When blocks are stitched, the junction conditions relate these proper times through lapse functions N_{ab} defined on each stitching surface $\mathcal{S}_{ab}$:

$$d\tau_b = N_{ab}\, d\tau_a. \tag{96}$$

For a network of blocks, consistency around any closed loop imposes the integrability condition

$$\prod_{e\in C} N_e^{\sigma(e)} = 1 \quad \text{for every closed loop } C. \tag{97}$$

When these loop conditions hold, there exists a global time function t such that on each block,

$$d\tau_a = N_a(t)\, dt. \tag{98}$$

Thus, a global time emerges relationally from the correlations between the block proper times.

In the quantum theory, the proper times become operators $\hat{\tau}_a$ that are conjugate to the energies $\hat{E}_a$. The lapse constraints become operator equations

$$\hat{N}_{ab}\Psi = 0, \tag{99}$$

which imply that the proper time operators are related by $\hat{\tau}_b = \hat{N}_{ab}\hat{\tau}_a$ on physical states. Choosing a reference block a_0 and setting $\hat{T} = \hat{\tau}_{a_0}$, we obtain a global time operator $\hat{T}$. The loop conditions guarantee that this definition is consistent. From the commutation relations $[\hat{\tau}_a, \hat{E}_a] = i\hbar$ and the constraints, one derives

$$[\hat{T}, \hat{E}_T] = i\hbar \quad \text{on physical states,} \tag{100}$$

where $\hat{E}_T = \sum_a \hat{E}_a$ is the total energy. Consequently, the time-energy uncertainty relation holds:

$$\Delta T\, \Delta E_T \geq \frac{\hbar}{2}. \tag{101}$$

Time is not an external parameter but a relational property of the stitched geometry. The Wheeler-DeWitt equation for the entire system becomes

$$\left(\sum_a \hat{\mathcal{H}}^{(a)} + \hat{\mathcal{H}}_{\text{stitch}}\right)\Psi = 0, \tag{102}$$

with no external time; evolution is encoded in the correlations between blocks as measured by the global time operator $\hat{T}$. This provides a constructive, relational resolution of the problem of time.

8. Experimental Signatures

The world-block model suggests a variety of testable predictions across different scales. Here we compile the main signatures; many depend on unknown parameters such as the block size ℓ or the precise value of ξ.

8.1. Quantum Domain

- **Repulsive self-gravity:** For a narrow wave packet of width $\sigma \sim \hbar^2/(Gm^3)$, the self-interaction induces a phase shift in matter-wave interferometry. For a molecule of mass $m \sim 10^5$ amu,

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$\sigma \sim 10^{-11}$ m. The phase shift scales as $\Delta\phi \sim Gm^2/(\hbar v)\ln(L/\sigma)$, where L is the interferometer arm length and v the velocity. For $L \sim 1$ m, $v \sim 100$ m/s, $\Delta\phi \sim 10^{-5}$ rad, potentially detectable with next-generation interferometers [16,17].
- **Phase-induced geometry:** The stress-energy tensor depends on gradients of the phase θ. In a double-slit experiment, the phase difference between paths modulates the geometry, producing a tiny time-dependent metric perturbation of order $\delta g/g \sim Gm^2\lambda/(\hbar c d^2) \sim 10^{-30}$ for typical parameters, far below current sensitivity but a conceptual novelty.
- **Entanglement witness:** The non-separable geometry of entangled blocks implies that entangled masses generate a gravitational field with quantum correlations. Experiments proposed by Bose et al. [18] and Marletto and Vedral [19] aim to detect gravitationally induced entanglement for masses $\sim 10^{-14}$ kg separated by $\sim 200\,\mu$m. The world-block model predicts a specific phase shift in such setups, which could be tested in the near future.
- **Quantum gravitational self-energy:** Each particle experiences quantum corrections to its self-interaction scaling as $\alpha_g \hbar c/m$ with $\alpha_g = Gm^2/(\hbar c)$. For an electron, this is tiny ($\sim 10^{-45}$), but for masses near the Planck scale it becomes order unity.

8.2. Gravitational Physics

- **Short-range deviations:** At distances comparable to the Compton wavelength of the mediating virtual blocks, the effective gravitational potential acquires Yukawa-type corrections. Current Eöt-Wash experiments constrain such deviations at the level $\alpha < 10^{-4}$ at 50 μm [20]. Improved experiments could tighten these bounds or detect the predicted signal.
- **Modified neutron star equation of state:** Repulsive self-gravity affects hydrostatic equilibrium, modifying the mass-radius relation. Observations from NICER and gravitational wave events like GW170817 [21] provide $\sim 10\%$ precision, which could constrain the model.
- **Modified gravitational wave propagation:** If the block size ℓ is not negligible, gravitational waves may experience dispersion and phase shifts as discussed in Sec. 3.6. This could be probed by future gravitational wave detectors.
- **Black hole information paradox:** Repulsive self-gravity may prevent formation of true singularities, replacing them with finite-sized cores. Information could escape from these cores, potentially resolving the information paradox.

8.3. Cosmology

- **Dark energy equation of state $w(z)$:** The harmonic field Φ_H evolves slowly, producing a time-dependent $w(z)$. A typical prediction is $w \approx -0.99$ at recent epochs, slightly different from -1. Upcoming surveys like DESI [22] and Euclid [23] will measure $w(z)$ to within $\sigma_w \sim 0.01$, potentially distinguishing the model from a cosmological constant.
- **CMB power spectrum oscillations:** Quantum fluctuations of the global clock may imprint oscillations in the primordial power spectrum at high multipoles. Planck currently places upper bounds; CMB-S4 and Simons Observatory will improve sensitivity [24].
- **Hubble tension:** The specific form of $w(z)$ could alleviate the 4σ–5σ discrepancy between early- and late-time H_0 measurements [25].
- **Stochastic gravitational wave background:** The stitching and unstitching of virtual blocks in the early universe may generate a stochastic GW background with a characteristic spectrum. Pulsar timing arrays (e.g., NANOGrav [26]) have reported a possible signal at nHz frequencies; LISA will probe the mHz band, and ground-based detectors the Hz–kHz band.

8.4. Quantum Gravity Effects

- **Modified dispersion relations:** At Planckian energies, the discrete block structure may modify the dispersion relation:

$$\omega^2 = c^2k^2 + m^2c^4 + \alpha\frac{\hbar^2}{A_0^2}k^4 + \cdots,$$

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with $\alpha \sim \mathcal{O}(1)$. For photons, this leads to an energy-dependent speed, causing time-of-arrival delays in gamma-ray bursts.

- **Running of Newton's constant:** Quantum loops lead to a logarithmic running of G with energy scale:

$$G(\mu) = G_0\left[1 + \beta \ln\left(\frac{\mu}{\mu_0}\right) + \cdots\right],$$

with $\beta \sim \mathcal{O}(1)$. This could explain small discrepancies in measurements of G at different scales.

Table 1. Summary of experimental signatures and projected sensitivities.

Effect	Order of Magnitude	Experiment/Future Sensitivity
Self-gravity phase shift	10^{-5} rad for 10^5 amu	Interferometry (next generation)
Phase-induced geometry	10^{-30} (gravity)	Far future
Dark energy $w(z)$	$\sigma_w \sim 0.01$	DESI, Euclid, Rubin
Modified dispersion	10^{-2} at 10^{19} eV	Cosmic ray observatories
CMB oscillations	$\beta < 0.01$	CMB-S4, Simons Observatory
Entanglement witness	10^{-14} kg masses	Bose–Marletto experiments
Stochastic GW background	$\Omega_{GW} \sim 10^{-9}$	PTA, LISA, LIGO/Virgo
Short-range gravity	$\alpha < 10^{-4}$ at 50 μm	Eöt-Wash, future tests
Running of G	$\beta \sim 1$	High-energy physics, cosmology
GW phase shifts	$\Delta\phi \sim 10^{-10}$ (nuclear scale)	LIGO/LISA (future)

9. Conclusion and Outlook

We have presented a relativistic and quantizable geometric framework in which each quantum particle is identified with a private spacetime governed by a universal stiffness constant A_0. The model is defined by a covariant action coupling a scalar field to the metric, with a curvature-coupling term $\xi R|\phi|^2$. From this action we derived the coupled Einstein–Klein–Gordon equations, which exhibit repulsive self-gravity and reduce to the non-relativistic version in the appropriate limit.

When many world-blocks are stitched together, the matching conditions and the continuum limit yield an emergent spacetime satisfying the full Einstein equations, with Newton's constant emerging as $G = c^4/(8\pi A_0)$. The harmonic zero-mode of the strain field naturally acquires a vacuum expectation value set by the Hubble radius, providing a non-perturbative explanation for dark energy without fine-tuning.

We quantized the model using canonical and path-integral methods. The multiple proper times and stitching constraints resolve the problem of time relationally, giving rise to an emergent global time operator. Quantum fluctuations of the lapse network lead to fundamental uncertainties that may imprint on the primordial power spectrum and generate a stochastic gravitational wave background.

The model predicts a host of testable signatures, from laboratory interferometry (repulsive self-gravity, phase-induced geometry, entanglement witnesses) to astrophysics (modified neutron star structure, short-range gravity deviations) and cosmology (dark energy equation of state, CMB oscillations, Hubble tension, stochastic GW background). These predictions offer a unified and falsifiable framework for quantum gravity.

Future work will focus on deriving precise numerical coefficients for the cosmological predictions, exploring the implications for black hole physics, and connecting the model to other approaches such as loop quantum gravity and string theory.

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Appendix A. Derivation of Field Equations

We provide a detailed derivation of the Einstein equations and the Klein–Gordon equation from the action (2). For clarity, we split the action into four parts:

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$$S_1 = \frac{A_0}{2} \int d^4x \sqrt{-g}\, R, \tag{A1}$$
$$S_2 = -\frac{1}{2} \int d^4x \sqrt{-g}\, g^{\mu\nu} \partial_\mu \phi^* \partial_\nu \phi, \tag{A2}$$
$$S_3 = -\frac{1}{2} m^2 c^2 \int d^4x \sqrt{-g}\, |\phi|^2, \tag{A3}$$
$$S_4 = -\xi \int d^4x \sqrt{-g}\, R\, |\phi|^2. \tag{A4}$$

Appendix A.1. Variation with Respect to the Metric

We vary each term with respect to $g^{\mu\nu}$. Recall that $\delta g = g_{\mu\nu} \delta g^{\mu\nu}$ and $\delta\sqrt{-g} = -\frac{1}{2}\sqrt{-g}\, g_{\mu\nu} \delta g^{\mu\nu}$.

Appendix A.1.1. Variation of S_1

The Einstein–Hilbert term yields the well-known result

$$\frac{\delta S_1}{\delta g^{\mu\nu}} = \frac{A_0}{2} \sqrt{-g}\, G_{\mu\nu}, \qquad G_{\mu\nu} = R_{\mu\nu} - \frac{1}{2} g_{\mu\nu} R. \tag{A5}$$

Appendix A.1.2. Variation of S_2

For the kinetic term,

$$\frac{\delta}{\delta g^{\mu\nu}} \left(\sqrt{-g}\, g^{\alpha\beta} \partial_\alpha \phi^* \partial_\beta \phi \right) = \sqrt{-g} \left(\partial_\mu \phi^* \partial_\nu \phi - \frac{1}{2} g_{\mu\nu} g^{\alpha\beta} \partial_\alpha \phi^* \partial_\beta \phi \right),$$

where we used $\delta g^{\alpha\beta} = \delta^\alpha_\mu \delta^\beta_\nu \delta g^{\mu\nu}$ and the variation of $\sqrt{-g}$. Thus

$$\frac{\delta S_2}{\delta g^{\mu\nu}} = -\frac{1}{2} \sqrt{-g} \left(\partial_\mu \phi^* \partial_\nu \phi - \frac{1}{2} g_{\mu\nu} |\nabla \phi|^2 \right). \tag{A6}$$

Here $|\nabla\phi|^2 = g^{\alpha\beta} \partial_\alpha \phi^* \partial_\beta \phi$.

Appendix A.1.3. Variation of S_3

The mass term gives simply

$$\frac{\delta S_3}{\delta g^{\mu\nu}} = -\frac{m^2 c^2}{4} \sqrt{-g}\, g_{\mu\nu} |\phi|^2. \tag{A7}$$

Appendix A.1.4. Variation of S_4: The Curvature-Coupling Term

For a scalar function f, the variation of $\int \sqrt{-g} R f$ is a standard result (see [27], Appendix E):

$$\delta \int d^4x \sqrt{-g} R f = \int d^4x \sqrt{-g} \Big[f G_{\mu\nu} + (g_{\mu\nu} \nabla^2 - \nabla_\mu \nabla_\nu) f \Big] \delta g^{\mu\nu}. \tag{A8}$$

Applying this with $f = |\phi|^2$, we obtain

$$\frac{\delta S_4}{\delta g^{\mu\nu}} = -\xi \sqrt{-g} \Big[|\phi|^2 G_{\mu\nu} + (g_{\mu\nu} \nabla^2 - \nabla_\mu \nabla_\nu) |\phi|^2 \Big]. \tag{A9}$$

Note the sign of the derivative term differs from the manuscript's original version.

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Appendix A.1.5. Einstein Equation

Summing the contributions from (A5)–(A9) and setting $\delta S/\delta g^{\mu\nu} = 0$, after dividing by $\sqrt{-g}$ we obtain:

$$\frac{A_0}{2} G_{\mu\nu} - \frac{1}{2}\left(\partial_\mu\phi^*\partial_\nu\phi - \frac{1}{2} g_{\mu\nu}|\nabla\phi|^2\right) - \frac{m^2c^2}{4} g_{\mu\nu}|\phi|^2 \\ - \xi\left[|\phi|^2 G_{\mu\nu} + (g_{\mu\nu}\nabla^2 - \nabla_\mu\nabla_\nu)|\phi|^2\right] = 0. \quad \text{(A10)}$$

Collecting the $G_{\mu\nu}$ terms and multiplying by 2 yields

$$(A_0 - 2\xi|\phi|^2)G_{\mu\nu} = \partial_\mu\phi^*\partial_\nu\phi + \partial_\nu\phi^*\partial_\mu\phi - \frac{1}{2} g_{\mu\nu}|\nabla\phi|^2 - \frac{m^2c^2}{2} g_{\mu\nu}|\phi|^2 + 2\xi(\nabla_\mu\nabla_\nu - g_{\mu\nu}\nabla^2)|\phi|^2. \quad \text{(A11)}$$

Defining the canonical stress-energy tensor

$$T^\phi_{\mu\nu} = \partial_\mu\phi^*\partial_\nu\phi + \partial_\nu\phi^*\partial_\mu\phi - g_{\mu\nu}\left(|\nabla\phi|^2 + m^2c^2|\phi|^2\right), \quad \text{(A12)}$$

we obtain the compact Einstein equation

$$\boxed{(A_0 - 2\xi|\phi|^2)G_{\mu\nu} = T^\phi_{\mu\nu} + 2\xi\left(\nabla_\mu\nabla_\nu - g_{\mu\nu}\nabla^2\right)|\phi|^2}. \quad \text{(A13)}$$

Appendix A.2. Variation with Respect to ϕ^*

Varying the action with respect to ϕ^* (treating ϕ and ϕ^* as independent) gives:

$$\frac{\delta S_2}{\delta\phi^*} = -\frac{1}{2}\sqrt{-g}\, g^{\mu\nu}\partial_\mu\partial_\nu\phi = -\frac{1}{2}\sqrt{-g}\,\Box\phi,$$

$$\frac{\delta S_3}{\delta\phi^*} = -\frac{m^2c^2}{2}\sqrt{-g}\,\phi,$$

$$\frac{\delta S_4}{\delta\phi^*} = -\xi\sqrt{-g}\,R\,\phi.$$

Setting $\delta S/\delta\phi^* = 0$ and dividing by $-\sqrt{-g}$ yields

$$\boxed{\nabla^2\phi - m^2c^2\phi - \xi R\,\phi = 0}. \quad \text{(A14)}$$

Appendix B. Non-Relativistic Reduction

We now derive the non-relativistic limit of the coupled field equations (A13)–(A14).

Appendix B.1. Coordinate Conventions and Expansions

We work in the Newtonian gauge with metric

$$g_{\mu\nu} = \eta_{\mu\nu} + h_{\mu\nu}, \qquad h_{00} = \frac{2\Phi}{c^2}, \quad h_{0i} = 0, \quad h_{ij} = \frac{2\Psi}{c^2}\delta_{ij}, \quad \text{(A15)}$$

with $\Phi, \Psi = \mathcal{O}(c^{-2})$. The scalar field is written as

$$\phi = \frac{1}{\sqrt{2m}}\, e^{-imc^2t/\hbar}\, \psi(t, \mathbf{x}), \quad \text{(A16)}$$

where ψ is slowly varying: $\partial_t\psi \sim \mathcal{O}(c^{-2})$, $\partial_i\psi \sim \mathcal{O}(1)$.

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Appendix B.2. Expansion of the Klein–Gordon Equation

The d'Alembertian in the weak-field metric is

$$\nabla^2\phi = \frac{1}{\sqrt{-g}}\partial_\mu\left(\sqrt{-g}\, g^{\mu\nu}\partial_\nu\phi\right).$$

To linear order, $\sqrt{-g} = 1 + \frac{1}{2}h$, with $h = -\frac{2\Phi}{c^2} + \frac{6\Psi}{c^2}$, and

$$g^{00} = -1 - \frac{2\Phi}{c^2}, \qquad g^{ij} = \delta^{ij} - \frac{2\Psi}{c^2}\delta^{ij}.$$

Inserting the ansatz (A16) and expanding to leading order in $1/c^2$, we obtain

$$i\hbar\,\partial_t\psi = -\frac{\hbar^2}{2m}\nabla^2\psi - m\Phi\,\psi + \frac{\hbar^2}{2m}\frac{2\Psi}{c^2}\nabla^2\psi - \frac{\hbar^2}{2m}\,\xi R\,\psi + \mathcal{O}(c^{-4}). \tag{A17}$$

The Ricci scalar to linear order is

$$R = -\frac{2}{c^2}\nabla^2\Phi + \frac{8}{c^2}\nabla^2\Psi + \mathcal{O}(c^{-4}). \tag{A18}$$

Appendix B.3. Einstein Equations in the Non-Relativistic Limit

Expanding the Einstein equations (A13), the factor $(A_0 - 2\xi|\phi|^2)$ is dominated by A_0 at leading order. Thus

$$A_0 G_{\mu\nu} \approx T^{\phi}_{\mu\nu} + 2\xi\left(\nabla_\mu\nabla_\nu - g_{\mu\nu}\nabla^2\right)|\phi|^2. \tag{A19}$$

Appendix B.3.1. The (00) Component

From linearized theory, $G_{00} \approx 2\nabla^2\Psi/c^2$. The leading term of T^{ϕ}_{00} is

$$T^{\phi}_{00} \approx mc^2|\psi|^2. \tag{A20}$$

The geometric term is of order $1/m$ times a spatial derivative, hence negligible at leading order. Thus

$$A_0\frac{2\nabla^2\Psi}{c^2} = mc^2|\psi|^2 \quad\Longrightarrow\quad \nabla^2\Psi = \frac{mc^4}{2A_0}|\psi|^2. \tag{A21}$$

Appendix B.3.2. The (ij) Components

The (ij) components give, to leading order,

$$A_0\left[\delta_{ij}\nabla^2(\Phi-\Psi) - \partial_i\partial_j(\Phi-\Psi)\right]/c^2 = 0, \tag{A22}$$

which implies $\Phi - \Psi =$ constant, and with boundary conditions, $\Phi = \Psi$.

Combining with (A21), we obtain the Poisson equation

$$\nabla^2\Phi = \frac{mc^4}{2A_0}|\psi|^2. \tag{A23}$$

Identifying Newton's constant via $G = c^4/(8\pi A_0)$ gives

$$\nabla^2\Phi = 4\pi G\rho_m, \qquad \rho_m = m|\psi|^2. \tag{A24}$$

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Appendix B.3.3. The Schrödinger Equation

Insert $\Phi = \Psi$ and the expression for R into (A17). The term $\frac{2\Psi}{c^2}\nabla^2\psi$ is of order c^{-4}, but the $\xi R\psi$ term is of order c^0 because $R \sim \nabla^2\Phi/c^2$ and Φ itself is of order c^0 in non-relativistic units. Therefore it must be kept:

$$i\hbar\partial_t\psi = -\frac{\hbar^2}{2m}\nabla^2\psi - m\Phi\psi - \frac{\hbar^2}{2m}\xi R\psi + \mathcal{O}(c^{-4}). \tag{A25}$$

Using $R = -2\nabla^2\Phi/c^2$ (since $\Phi = \Psi$), the ξ term is $-\frac{\hbar^2}{2m}\xi(-2\nabla^2\Phi/c^2)\psi = \frac{\hbar^2\xi}{mc^2}\nabla^2\Phi\psi$. This is a higher-order correction compared to $m\Phi\psi$ because $\hbar^2/(mc^2)$ is the Compton wavelength squared, which is small for non-relativistic particles. Hence the leading order Schrödinger equation is

$$i\hbar\partial_t\psi = -\frac{\hbar^2}{2m}\nabla^2\psi - m\Phi\psi. \tag{A26}$$

For a point mass, $\Phi(r) = -Gm/r$, giving the repulsive force $F_r = +Gm^2/r^2$.

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