Recasting Schrödinger's Cloud-Like Entity within Relativistic Geometry

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Abstract

This paper revisits Schrödinger's 1927 concept of a particle as a spatially diffuse "cloud," redefining it as a true spacetime entity equipped with its own intrinsic metric. Each cloud-like entity is structured by constant proper time slices, across which mass density is continuously distributed, while a single world line threading these slices carries all gauge charges. The intrinsic metric is determined by the mass density via a Poisson-like equation, with a Green kernel that exhibits two distinct phases. In the free phase, the kernel has a non-local Coulomb-like form that links every point within the cloud-like entity. During measurement, however, the detector imposes a time-dependent boundary condition, smoothly deforming the kernel into a sharply localized monopole well. This defines a finite "collapse window," during which the wave packet's width contracts, avoiding singularity and allowing for potential re-expansion if the detector pulse ceases prematurely. The formalism provides specific, testable predictions: mass-dependent localization times (of the order of 10 ps for electrons), reversible loss and revival of interference patterns, gradual decay of Bell correlations, and an ultra-weak, transient metric force on nearby probes. Together, these results offer a deterministic yet non-separable framework that bridges quantum non-locality with relativistic causality.

Keywords: spacetime; extended-entity; proper-time; non-locality; configuration space.

1. Introduction

Einstein's theory of relativity recast physics around two foundational principles: the invariance of the speed of light and the dynamic interplay between mass-energy and the geometry of spacetime. Relativity eliminated the concept of a preferred reference frame and unified space and time into a four-dimensional structure, constrained in its causal order by the speed of light [1] [2]. General relativity extended this framework by demonstrating that mass-energy curves spacetime, with free motion following the geodesics of this curvature. In both theories, locality remains a guiding principle, as influences propagate continuously through fields or spacetime curvature, without crossing spacelike separations.

In contrast, quantum mechanics challenges these principles, particularly the concept of locality, with phenomena that defy classical intuition. Through the principle of superposition, a system can exist in a spectrum of mutually exclusive states, while entanglement creates correlations between distant measurements that cannot be accounted for by any local hidden variables. This tension became clear through the EPR argument [3] and Bell's inequality [4], and experiments such as Aspect's photon tests [5] have consistently confirmed the nonlocal predictions of quantum theory, without violating the relativistic no-signalling condition.

To address these quantum challenges, various interpretations have emerged [6]. The Copenhagen interpretation [7], treats quantum mechanics as a pragmatic tool for predicting observable outcomes rather than describing an underlying reality. Quantum Bayesianism QBism [8] reframes wavefunction collapse as a Bayesian update of subjective knowledge rather than an objective event [9]. Everett's Many-Worlds Theory [10], endorsed by contemporary advocates like Wallace [11] [12], posits that every quantum event spawns a branching of worlds, providing a deterministic but highly counterintuitive perspective. The de Broglie-Bohm pilot-wave theory introduces hidden variables, asserting that particles follow deterministic trajectories guided by a wavefunction [13]. Alternatively, GRW's wavefunction-collapse theory [14] incorporates stochastic elements into Schrödinger's equation, offering simplicity but facing issues with relativistic simultaneity.

A central and unresolved issue remains the reconciliation [15] of quantum mechanics with the relativistic doctrine of locality. The dilemma is intensified by the recognition that multi-particle wavefunctions reside within a high-dimensional configuration space. This abstraction raises significant questions about the connection between wavefunctions and the familiar three-dimensional world. The ontological status of the wavefunction continues to spark debate [16], dividing views between those who see it as a tool for knowledge (epistemological) and those who regard it as a real, physical entity (ontological).

In his 1927 Solvay lecture, Schrödinger proposed a shift away from the classical notion of point particles, suggesting instead a continuous entity that permeates space [17]. While this idea is compelling, modelling a particle as an extended mass or charge density introduces significant challenges in the realms of general relativity (GR), quantum field theory (QFT), and continuum mechanics.

For example, treating an electron as a smeared charge cloud faces the problem of divergent self-interaction energies. Experimental evidence, ranging from the photoelectric effect to high-energy scattering, restricts the electron's charge radius to scales no larger than approximately 10^{-22} m. This strongly supports point-like behaviour during energy–momentum exchanges.

In GR, an extended mass distribution generates a gravitational field that curves spacetime beyond the object's boundaries. A smoothly decaying mass density complicates the particle's internal dynamics, as gravitational time dilation causes different regions to experience varying proper times. These variations introduce internal stresses—such as pressure and tension—that contribute to the particle's total mass-energy.

In quantum mechanics, a particle's position is described by a wavefunction that spreads across space until measurement collapses it into a sharply localized state. Extending this idea to a particle's mass density implies a nonlocal collapse of its entire distribution, raising challenges for compatibility with relativity since instantaneous collapse could violate causality.

Despite these difficulties, Schrödinger's vision of "cloud-like" quantum entities has inspired various research directions. Wavefunction realism holds that the wavefunction, defined in configuration space, represents the fundamental ontology from which three-dimensional structures arise [18]**.** Objective-collapse theories, like GRW and CSL, modify quantum dynamics with stochastic elements to ensure spatial localization [19]**.** Schrödinger-Newton models, developed by Diosi and Penrose, incorporate gravitational self-interactions to limit wavefunction spreading [20]**.** Stochastic electrodynamics proposes a Lorentz-invariant random vacuum field—analogous to quantum vacuum fluctuations—to stabilize spread-out charge distributions [21]**.** The Diosi-Penrose model further explores gravity-induced mechanisms for spontaneous wavefunction collapse based on gravitational energy differences [22]**.** Each of these approaches aims to reconcile the spatial extension of quantum systems with empirical localisation and finite self-energies.

Building on this intuition, our approach takes a geometric perspective. Each quantum entity is attributed a non-local spacetime geometry, determined by its diffused mass density. This geometry includes a proper-time foliation, marking a single world-line on which the particle's charge is confined. Within this intrinsic geometry, the Schrödinger–Newton potential emerges naturally—not as an external semiclassical correction but as a property intrinsic to the quantum entity itself. In this framework, quantum mechanics becomes a tool for probing the fine-grained, relational structure of spacetime, while relativity continues to describe its global architecture. Together, these insights create a unified framework that accommodates quantum non-locality and relativistic causality.

2. Fermi–Walker Coordinates around an Arbitrary World-line

We work on a smooth, connected, four-dimensional differentiable manifold M, equipped with a Lorentzian metric $g_{\mu\nu}(X)$ where we adopt the $(+,-,-,-)$ signature. This metric is a position-dependent, symmetric tensor field that defines the invariant line element:

$$ds^2 = g_{\mu\nu}(X)X^{\mu}X^{\nu} \tag{1}$$

This line element allows for the measurement of proper time along time-like curves and spatial distances within spacelike hypersurfaces. Since the metric $g_{\mu\nu}$ is completely general—without assuming any specific coordinate system or symmetry—it can describe any local curvature and any distribution of matter and energy consistent with Einstein's field equations. Throughout what follows, this metric is treated as the underlying geometric framework for constructing specific coordinate systems, such as the Fermi–Walker system, which will be introduced next.

The Fermi–Walker coordinates are a local coordinate system defined in the neighbourhood of a smooth time-like world-line in an arbitrary spacetime with metric $g_{\mu\nu}$. They provide a natural framework for describing the geometry in the vicinity of a reference observer following the world-line, regardless of whether this observer is inertial or undergoing arbitrary acceleration.

Let $\gamma\colon \tau \to X^{\mu}(\tau)$ be the time-like world-line—inertial or accelerated—in the arbitrary spacetime with metric $g_{\mu\nu}$.

An Orthonormal tetrad $\{e_0, e_i\}$ can be associated with the world-line γ, where: $e^{\mu}_{(0)} = u^{\mu} = \dot{X}^{\mu} = dX^{\mu}/d\tau$ is the 4-velocity of the reference observer, satisfying $u^{\mu}u_{\mu} = -1$. The spatial triad $\{e_i\}$ (where e_i are the basis vectors of the spatial triad) satisfies $e^{\mu}_i e_{j\mu} = \delta_{ij}$ and is transported along γ according to Fermi–Walker transport:

$$\frac{De^{\mu}_i}{d\tau} = (a_{\nu}e^{\nu}_i)u^{\mu} - (a_{\nu}u^{\nu})e^{\mu}_i \tag{2}$$

where $a^{\nu} = Du^{\nu}/d\tau$ is the proper acceleration of the world-line γ.

Fermi–Walker coordinates are defined as follows: For any point P sufficiently close to the reference world-line γ, there exists a unique spacelike geodesic σ connecting P to $\gamma(\tau)$ that is orthogonal to $u^{\mu}(\tau)$ at its initial point. Let ξ^{μ} be the unit tangent to σ at $\gamma(\tau)$ and let σ also denote the geodesic distance from $\gamma(\tau)$ to P. Then the Fermi–Walker coordinates (x^0, x^i) of P are $x^0 \equiv \tau,\ x^i \equiv \sigma(\xi^{\mu}e_{i\mu})$.

Thus, $x^0 \equiv \tau$ is the proper time along the reference world-line γ, while $\vec{x} = (x^1, x^2, x^3)$ are the proper distances measured in the transported triad, where the spacelike geodesic connecting P to $\gamma(\tau)$ is orthogonal to $u^{\mu}(\tau)$.

$x^i \equiv \sigma(\xi^{\mu}e_{i\mu})$ are the components of the signed spatial displacement of P along the geodesic with tangent ξ^{μ}, where σ is the geodesic distance and e_i are the basis vectors of the spatial triad.

In these coordinates, the metric admits the standard canonical Fermi expansion to quadratic order in x^i (components of the Riemann tensor evaluated on γ and projected on the tetrad):

$$\begin{aligned} g_{00} &= -(1 + a_k x^k)^2 - R_{0i0j}x^i x^j + O(x^3) \\ g_{0i} &= -\frac{2}{3}R_{0jik}x^j x^k + O(x^3) \\ g_{ij} &= \delta_{ij} - \frac{1}{3}R_{ikjl}x^k x^l + O(x^3) \end{aligned} \tag{3}$$

where $a_k = a_{\mu}e^{\mu}_k$ represents the components of the proper acceleration of the world-line γ and R_{ikjl} are the Riemann curvature tensor components evaluated on γ. In the special case of a geodesic world-line γ ($a_k = 0$) in a flat spacetime, the $O(x^2)$ curvature corrections vanish, and the metric $g_{\mu\nu}$ reduces to the Minkowski form.

At each constant value of $x^0 \equiv \tau$, the Fermi–Walker coordinates define a spacelike hypersurface S_τ (Fermi slice), given by:

$$S_\tau = \{(\tau, \vec{x}) \ \mid x \epsilon R^3\} \tag{4}$$

which satisfies the following conditions:

$u^\mu \nabla_\mu = 1$, implying that the time coordinate is equal to proper time;

$u^\mu \nabla_\mu x^i = 0$, meaning that the spatial coordinates remain constant along γ; and

$u^\mu \mathrm{n}_\mu = 0$, where $\mathrm{n}_\mu = -\nabla_\mu x^0$ represents the normal to S_τ.

These surfaces are intrinsically simultaneous for the observer following the world-line γ, and the induced spatial metric on the slice S_τ is:

$$ds^2|_{S_\tau} = h_{ij}(\tau, \vec{x}) dx^i dx^j , \qquad h_{ij} = \delta_{ij} - \frac{1}{3} R_{ikjl} x^k x^l + O(x^3) \tag{5}$$

where h_{ij} remains positive-definite.

All curvature components in Equations (3)-(5) are defined in the tetrad basis associated with γ. The domain of validity is limited to the tubular neighbourhood where the orthogonal geodesic from γ to point P is unique, ensuring the absence of conjugate points.

We propose that spacetime emerges from the integration of intrinsic spacetimes linked to elementary objects. For each world-line γ, the intrinsic spacetime is represented by the Fermi–Walker (FW) tubular neighbourhood, equipped with an FW coordinate chart. Interactions ensure overlaps between such neighbourhoods. The global GR manifold (M, g) is then constructed as the union of these neighbourhoods, ensuring consistency through standard atlas constraints. These include metric agreement through pullbacks, frame matching up to local Lorentz transformations, and gauge consistency up to gauge transformations, as described further in the next Section.

Given the Fermi–Walker–transported spatial triad $\{e_i^\mu(\tau)\}_{i=1}^3$, associated with the reference world-line $\gamma : \tau \to Z^\mu(\tau)$, the FW intrinsic coordinates $(\tau, \vec{x}) \ \epsilon \ R \times R^3$, referred to as 'intrinsic coordinates', can be related to arbitrary spacetime coordinates $(t, \vec{X})$, referred to as 'extrinsic coordinates', by a smooth, locally bijective (diffeomorphic) coordinate transformation:

$$\chi: (\tau, \vec{x}) \to X^\mu(t, \vec{X}) = \left(t(\tau, \vec{x}), \vec{X}(\tau, \vec{x})\right) = exp_{Z(\tau)}\left[x^i e_i^\mu(\tau)\right] \tag{6}$$

Here X^μ represents the extrinsic coordinates of an event in the laboratory frame, while $Z^\mu(\tau)$ represents the specific world-line chosen as the Fermi origin. Specifically, $Z^\mu(\tau)$ (where $Z^\mu(\tau) = X^\mu(\tau, \vec{x} = 0) = \chi(\tau, 0)$) indicates to an external observer the position of the particle's world-line γ at each proper time τ. From the intrinsic perspective, the particle's world-line γ is simply represented as $(\tau, \vec{x}) = (\tau, 0)$.

The mapping described above indicates that, starting from the event $Z(\tau)$ on the world-line γ, we follow the unique spacelike geodesic with an initial tangent vector $x^i e_i^\mu(\tau)$. The endpoint of this geodesic is given in standard extrinsic coordinates as $X^\mu(t, \vec{X})$.

3. Intrinsic hypersurface basis and four-dimensional 'world-block'

Each intrinsic proper-time slice S_τ in the intrinsic FW chart $(\tau, \vec{x})$, may be conceptualised as providing a basis $\{|\vec{x}\rangle\}$ associated with a rigged Hilbert space (Gel'fand triple) $S_\tau \subset \mathcal{H}_\tau \subset S'_\tau$. Here, S_τ represents the space of smooth, rapidly decaying test functions, $\mathcal{H}_\tau$ the usual Hilbert space, and S'_τ its distributional dual, which includes Dirac δ-functions and plane-wave eigenstates.

The elements of this space, denoted as $|\vec{x}\rangle$, are indexed by a continuous variable $\vec{x}$. These elements are normalized using the Dirac δ-function, ensuring orthonormality within the space. The Dirac-normalised basis on every hyperplane slice S_τ, is expressed as follows:

$$\{|\vec{x}\rangle\}_{\vec{x}\epsilon R^3} \quad ; \quad \langle\vec{x}|\vec{x}'\rangle = \delta^{(3)}(\vec{x}-\vec{x}')/\sqrt{h(\tau,\vec{x})} \tag{7}$$

where $h(\tau,\vec{x}) \equiv deth_{ij}(\tau,\vec{x})$ represents the determinant of the induced spatial metric on the slice. Thus a complete basis—and its distributional extension—resides on every slice S_τ. The completeness relation for any given slice S_τ is expressed as:

$$\int d\,^3x\sqrt{h(\tau,\vec{x})}\,|x\rangle\langle x| = 1 \tag{8}$$

Since the evolution from one proper-time slice to the next is unitary, all Hilbert spaces associated with these slices can be considered equivalent and identified as a single, abstract Hilbert space $\boldsymbol{\mathcal{H}_\tau}$.

A single particle at a given proper time τ is represented by its state vector in this Hilbert space:

$$|\varphi(\tau)\rangle \;=\int d\,^3x\sqrt{h(\tau,\vec{x})}\,\varphi(\tau,\vec{x})\,|\vec{x}\rangle; \quad \varphi(\tau,\vec{x}) = \langle\vec{x}|\varphi(\tau)\rangle \tag{9}$$

The normalisation condition $\langle\vec{x}|\varphi(\tau)\rangle = 1$ is expressed as:

$$\int d\,^3x\sqrt{h(\tau,\vec{x})}\,|\varphi(\tau,\vec{x})|^2 = 1 \tag{10}$$

All points $\vec{x}$ within the same proper-time slice S_τ are intrinsically simultaneous, meaning the entire function $\varphi(\tau,\vec{x})$ corresponds to a single "instant" of proper time τ.

Proper-time evolution is governed by the slice-dependent Hamiltonian density $H(\tau)$, through the equation:

$$i\hbar\partial_\tau|\varphi(\tau)\rangle \;= H(\tau)|\varphi(\tau)\rangle \tag{11}$$

and the evolution operator, or propagator, is given by:

$$U(\tau_2,\tau_1) = exp\left[-\frac{i}{\hbar}\int_{\tau_1}^{\tau_2} d\tau' H(\tau')\right] \tag{12}$$

Proper time τ thus provides an intrinsic, observer-invariant ordering of the slices without requiring any external clock variable.

Following Schrödinger's "cloud" model of a free particle, the mass density on each proper-time slice S_τ is expressed as:

$$\mu(\vec{x},\tau) = m|\varphi(\tau,\vec{x})|^2 \tag{13}$$

The total mass m on each proper-time slice S_τ is conserved and is expressed as:

$$\int d\,^3x\sqrt{h(\tau,\vec{x})}\,\mu(\tau,\vec{x}) = m \tag{14}$$

When all slices are combined, they form a four-dimensional world-block $b \equiv \bigcup_{\tau\epsilon R} S_\tau$, which constitutes a smooth manifold. Each point in this manifold is characterized by the scalar field $\mu(\vec{x},\tau)$. In the Fermi–Walker coordinate system, every point of the world-block is uniquely identified by its proper time τ and spatial Fermi coordinates $\vec{x}$.

For simplicity, we absorb the constant m into the scalar field and work instead with the dimensionless density:

$$\rho(\tau,\vec{x}) = |\varphi(\tau,\vec{x})|^2 \ ; \ \int d^{\,3}x\,\sqrt{h(\tau,\vec{x})}\,\rho(\tau,\vec{x}) = 1 \tag{15}$$

For a free Gaussian wave packet, the width $\sigma(\tau)$ increases with proper time τ, causing $\rho(\tau,\vec{x})$ to spread and dilute, precisely as envisioned by Schrödinger's equation.

On each proper-time slice S_τ, there exists a unique centroid $\vec{x}_c(\tau) \in R^3$ such that the first mass dipole moment on the slice S_τ, vanishes. This condition is expressed as:

$$\int \big(\vec{x} - \vec{x}_c(\tau)\big)\rho(\tau,\vec{x})d^3x = 0 \tag{16}$$

Here, the centroid $\vec{x}_c(\tau)$ represents the intrinsic centre of mass-energy for the world-block on the slice S_τ. The reference world-line anchoring the world-block is therefore, the curve γ, which passes through the centroids of the slices $(S_\tau)_{\tau\in R}$ sequentially. Thus, the world-line γ is redefined as:

$$\gamma : \tau \to X^\mu(\tau) = \big(\tau, \vec{x}_c(\tau)\big)\ \epsilon\ b \tag{17}$$

The reference world-line γ represents the intrinsic centre-of-mass line of the world-block and serves as its anchor, carrying all associated gauge charges. Notably, for a particle with charge e, the charge-density $\rho_e(\vec{x},\tau)$, unlike the mass-density $\rho(\vec{x},\tau)$, is entirely confined to the reference world-line γ. This confinement arises as a consequence of local gauge symmetry, renormalizability, and scattering bounds. Consequently, the intrinsic charge density is formulated as:

$$\rho_e(\vec{x},\tau) = e\delta^{(3)}\big(\vec{x} - \vec{x}_c(\tau)\big) \tag{18}$$

More generally, the reference world-line γ serves as the carrier of all strictly conserved point-like quantum numbers, such as electric charge, baryon number, lepton number. In contrast, the mass density is free to form an extended cloud distributed over each proper-time slice S_τ, surrounding the world-line. This distinction arises naturally from symmetry and effective field theory considerations.

Gravity couples to the stress–energy tensor $T_{\mu\nu}$, which can reflect the spatial spread of a quantum state (e.g., $\langle\psi|\hat{T}_{\mu\nu}|\psi\rangle$), provided $\nabla_\mu T_{\mu\nu} = 0$. In contrast, gauge charges couple through local conserved currents j^μ, which transform covariantly under the gauge group and are confined to a single world-line.

Renormalisation clarifies this distinction. The UV consistency of Quantum Electrodynamics (QED) requires point-like charge sources, while allowing mass-energy to be spatially distributed in gravity aligns with low-energy effective field theory. Experimentally, electromagnetic scattering limits the electron's charge radius to less than $10^{-22}\ m$, consistent with a world-line current. At the same time, effects like interferometric phase shifts are sensitive to the spatial distribution of mass-energy.

A consistent framework emerges: Einstein's equations are sourced by the (possibly smeared) mass-energy $T_{\mu\nu}$ of the quantum state, while Maxwell or Yang–Mills equations are sourced by a conserved current confined to the world-line. Standard conservation laws $\nabla \cdot T = 0$ and $\partial \cdot j = 0$ ensure consistency across overlapping intrinsic spacetimes. Electromagnetic phenomenology remains unchanged, while gravity, in principle, can probe wave-packet widths through Schrödinger–Newton-type effects. This approach leads to testable predictions without requiring ad hoc assumptions.

Thus the world-block b represents the complete ontic entity. its slices carry the extended mass density, its centre-of-mass line γ carries all gauge charges, and the intrinsic proper time τ sequentially orders the slices independently of any external clock. It is important to note that variations in initial conditions, external fields, or experimental set-ups yield different coordinate transformations χ, but none of them alters the intrinsic foliation $\{S_\tau\}$: all points of a given slice always share the same proper time τ.

We note that the Fermi-Walker coordinates are applicable within the normal-convex tube surrounding the reference world-line γ, with the radius of this tube determined by the smaller of two scales: the tidal-curvature scale l_{curv} and the proper acceleration scale l_{acc}. These coordinates are well-suited to describe the mass-density distribution $\rho(\tau, \vec{x}) = |\varphi(\tau, \vec{x})|^2$ of a particle with a wavefunction, such as a three-dimensional Gaussian:

$$\rho(\tau, \vec{x}) = \left(\frac{1}{2\pi\sigma^2}\right)^{3/2} e^{-r^2/2\sigma^2}; \qquad r = \sqrt{\delta_{ij} x^i x^j} \tag{19}$$

provided that the effective support $r_{eff} \equiv N\sigma$ (with $N = 3$ for a Gaussian) satisfies the following condition:

$$r_{eff} \equiv N\sigma = min\left\{l_{curv} \equiv \left|R_{ikjl}\right|^{-1/2}, l_{acc} \equiv c^2/|a|\right\} \tag{20}$$

where $l_{curv} \equiv \left|R_{ikjl}\right|^{-1/2}$ is the local curvature radius, $l_{acc} \equiv c^2/|a|$ is the inverse-acceleration length, and N is chosen to define the "edge" of the wave packet (for a Gaussian, $N = 3$ already encloses 99.7 % of the density ρ).

In practice, this condition is easily satisfied. For a free proton or electron in the lab, with typical wave packet width $\sigma \approx 1 - 100\ \mu m$ and $l_{curv} > 10^4\ m$, the inequality $N\sigma \ll l_{curv}$ holds by at least eight orders of magnitude. For a proton in the Large Hadron Collider (LHC), where strong bending occurs, the typical width is $\sigma \approx 10 - 50\ \mu m$ and the acceleration magnitude $|a| \approx 1.8 \times 10^{21} ms^{-2}$ leads to a proper acceleration scale $l_{acc} \approx 50\ \mu m$. In this case, $3\sigma \approx l_{acc}$, and the inequality still holds. For an electron, the proper acceleration scale $l_{acc} \approx 20\ m$ and thus the inequality is largely satisfied.

Thus, for modern laboratory wave packets, the effective support $r_{eff} \equiv N\sigma$ lies comfortably within the region where the exponential map along spacelike geodesics orthogonal to the world-line γ remains injective and the canonical Fermi-normal metric of Equation (3) remains a good approximation.

Although a Gaussian never truly vanishes, the coordinate expansion only breaks down beyond a finite $N\sigma$, where the tails of the wavefunction extend. This presents no physical issues, as the density contribution from these tails is exponentially small. If necessary, spacetime can be patched smoothly with another normal chart without losing accuracy. All meaningful observables, where the density ρ is significant, are fully captured within the Fermi tube.

Ultimately, for any realistic Gaussian-like density distribution of an elementary particle, present-day accelerators and laboratory fields satisfy $N\sigma \ll min\{l_{curv}, l_{acc}\}$. As a result, Fermi coordinates provide a valid and efficient framework for describing both the metric and the density profile throughout the entire region where the particle's wavefunction has a significant amplitude.

When two intrinsic Fermi–Walker neighbourhoods U_α and U_β overlap, their stitching is governed by the physics of interactions. Let $g^{(\alpha)}$ and $g^{(\beta)}$ represent the metrics in the α and β charts respectively. On the overlap $U_\alpha \cap U_\beta$, the intrinsic coordinates are related by a C^2 transition map $\chi_{\beta\alpha} = \chi_\beta \circ \chi_\alpha^{-1}$, which transforms α-coordinates into β-coordinates.

To ensure metric compatibility—required for all local charts to describe the same global spacetime—the β-metric, when pulled back along $\chi_{\beta\alpha}$, must equal the α-metric: $g^{(\alpha)} = \chi_{\alpha\beta}^{*} g^{(\beta)}$. Specifically, let $\chi_\alpha : U_\alpha \to R^4$ and $\chi_\beta : U_\beta \to R^4$ be the FW charts, and g the single spacetime metric on M. The coordinate metrics are then given by pullbacks of g using the inverse charts:

$$g^{(\alpha)} = (\chi_\alpha^{-1})^{*} g, \quad g^{(\beta)} = \left(\chi_\beta^{-1}\right)^{*} g \tag{21}$$

It follows that:

$$\chi_{\beta\alpha}^{*} g^{(\beta)} = \left(\chi_{\beta} \circ \chi_\alpha^{-1}\right)^{*} \left(\left(\chi_\beta^{-1}\right)^{*} g\right) = (\chi_\alpha^{-1})^{*} g = g^{(\alpha)} \tag{22}$$

Thus both charts describe the same geometry on the overlap. Equivalently, $g^{(\beta)} = \chi_{\alpha\beta}^{*} g^{(\alpha)}$, where $\chi_{\alpha\beta} = \chi_{\alpha} \circ \chi_\beta^{-1}$. Additionally, orthonormal frames match up to a local Lorentz transformation, determined by parallel transport along the unique orthogonal spacelike geodesic.

On common slices, conserved currents are consistent under pullback and the probability measure remains invariant:

$$\rho^{(\alpha)} \sqrt{h^{(\alpha)}} d^3 x_\alpha = \rho^{(\beta)} \sqrt{h^{(\beta)}} d^3 x_\beta \tag{23}$$

With a single stress–energy tensor $T_{\mu\nu}$ defined by the matter and fields of all objects, the Einstein–matter equations $G_{\mu\nu}[g] = 8\pi T_{\mu\nu}$ along with $\nabla_\mu T_{\mu\nu} = 0$ enforce these identifications across overlaps, ensuring a unique global metric g.

Therefore, the atlas of intrinsic Fermi–Walker charts generated by interacting world-blocks is mathematically sufficient to reconstruct a single, smooth general-relativistic spacetime: the stitched global manifold.

4. Intrinsic non-separable metric

Within the Fermi chart, every point in the world-block is identified by its proper-time coordinate τ and spatial Fermi coordinates $\vec{x}$. On each proper-time slice S_τ, we consider the induced metric $h_{ij}(\tau, \vec{x})$ with its determinant given by $h = det h_{ij}$. At a given proper time τ, the one–particle state $|\varphi(\tau)\rangle$ is defined as:

$$|\varphi(\tau)\rangle = \int d^{\,3}x \sqrt{h(\tau, \vec{x})}\, \varphi(\tau, \vec{x})\, |\vec{x}\rangle; \qquad \int d^{\,3}x \sqrt{h}\, |\varphi|^2 = 1 \tag{24}$$

The mass density is given by $\mu(\tau, \vec{x}) = m|\varphi(\tau, \vec{x})|^2 = m\rho(\tau, \vec{x})$. Under the hypothesis that the world-block represents a distribution of mass density, it is natural to relate the intrinsic metric to the density $\rho(\tau, \vec{x})$. Since proper time τ has already been identified as the temporal parameter, the action S can be reformulated as an integral over proper time, involving a purely spatial Lagrangian density L, which may encode the intrinsic geometry and dynamics of the world-block. This decomposition can be expressed as:

$$S = \int d\tau \int d^3x \sqrt{h} L \left(h_{ij}(\tau, \vec{x}), \rho(\tau, \vec{x}), \varphi(\tau, \vec{x})\right) \tag{25}$$

where $h(\tau,\vec{x})$ is the determinant of the induced metric $h_{ij}(\tau,\vec{x})$ on the proper-time slice S_τ.This decomposition incorporates the intrinsic geometry and dynamics of the world-block in the context of proper-time evolution.

A minimal form for the Lagrangian density L, which couples the intrinsic geometry to the cloud density and recovers ordinary Schrödinger dynamics, is described as:

$$L = A\mathrm{D}_k h_{ij} D^k h^{ij} + B h_{ij} h^{ij} \rho - \frac{\hbar^2}{2m} h^{ij} \partial_i \varphi^* \partial_j \varphi + i\hbar \varphi^* \partial_\tau \varphi \tag{26}$$

where D_i denotes the covariant derivative associated with the intrinsic spatial metric h_{ij}. The first term $A\mathrm{D}_k h_{ij} D^k h^{ij}$, represents the "metric stiffness," which characterises the 'elastic energy' of the intrinsic spatial metric h_{ij}, quantifying how the geometry resists deformation. The second term $Bh_{ij}h^{ij}\rho$, introduces a 'geometry-matter', linking the intrinsic geometry of the slice to the mass/energy density. The third and fourth terms represent the standard Schrödinger kinetic energy term and the proper time-derivative term, respectively.

This formulation establishes a unified and consistent framework that connects the intrinsic geometry, the mass/energy distribution, and the quantum matter dynamics within a single world-block.

To derive the Euler-Lagrange equations, we calculate variations of the Lagrangian (26) with respect to h^{ij} and φ^*.

The variation with respect to h^{ij} yields a Poisson-type constraint, linking the intrinsic spatial geometry to the mass/energy density. This constraint takes the general form:

$$\mathrm{D}^2 h_{ij}(\tau,\vec{x}) = \frac{B}{2A} h_{ij} \rho(\tau,\vec{x}) \tag{27}$$

The left-hand side of the equation represents the spatial Laplacian of the intrinsic metric h_{ij}, which captures the spatial variation of the geometry on a constant proper time slice S_τ. The right-hand side provides the 'self-source term' responsible for modifying the intrinsic geometry. This constraint can be interpreted as an equation governing the evolution of the intrinsic geometry h_{ij} in response to the density $\rho(\tau,\vec{x})$.

When $h_{ij} \approx \delta_{ij}$ the Laplacian simplifies to the flat ∇^2 and Equation (27) is solved by:

$$h_{ij}(\tau,\vec{x}) = -\frac{B}{2A} \delta_{ij} [\varphi_C(\tau,\vec{x}) + h_H(\tau,\vec{x})] \tag{28}$$

where $\varphi_C(\tau,\vec{x})$ is a Coulomb-type integral, and $h_H(\tau,\vec{x})$ is a harmonic solution, satisfying the following:

$$\varphi_C(\tau,\vec{x}) = \int d^3y \frac{\rho(\tau,\vec{x})}{4\pi|\vec{x}-\vec{y}|}; \qquad \nabla^2 h_H(\tau,\vec{x}) = 0 \tag{29}$$

The intrinsic metric is structurally non-local because each value of $h_{ij}(\tau,\vec{x})$ is determined by convolving the mass-density $\rho(\tau,\vec{x})$ with the Green's kernel of the slice Laplacian. The integral over $\vec{y}$ indicates that the metric h_{ij} at $\vec{x}$ depends on the density distribution $\rho(\tau,\vec{x})$ across the entire slice. In other words, it is influenced by contributions from all other points $\vec{y}$ within the world-block. This does not involve direct action-at-a-distance but rather reflects an elliptic, instantaneous relationship.

Furthermore, the intrinsic metric is non-separable, as it cannot be decomposed into a product or sum of independent sub-metrics corresponding to disjoint spatial regions within the world-block. This non-separability implies that any disturbance in the intrinsic metric in one region cannot be described independently of its effects on other regions. Consequently, this introduces a form of acausal dependence, where changes in one part of the metric affect the entire system without a causal propagation.

On the other hand, the variation of the Lagrangian with respect to $\varphi^*(\tau,\vec{x})$ results in the standard Schrödinger-Newton equation [22] [23] which can be expressed as:

$$i\hbar\partial_\tau\varphi(\tau,\vec{x}) = -\frac{\hbar^2}{2m}\nabla^2\varphi(\tau,\vec{x}) + \frac{3B^2}{2A}[\varphi_C + h_H]\varphi(\tau,\vec{x}) \tag{30}$$

In this formalism, the Schrödinger-Newton equation provides a self-consistent description of the quantum dynamics of the wavefunction $\varphi(\tau,\vec{x})$, while the intrinsic metric h_{ij} evolves under the constraints imposed by the Poisson-type Equation (27) derived earlier. This coupled system ties the quantum dynamics to the geometry of a single world-block.

It is to be noted that Equation (30) is not a Newtonian approximation. In this equation the wavefunction $\varphi(\tau,\vec{x})$ evolves in the particle's own proper time τ like a Schrödinger equation. The first term represents the standard quantum kinetic energy term, while the second acts as an effective geometric potential $(3B^2/2A)[\varphi_C + h_H]$, which arises from the way gravity shapes space on each slice.

The spacetime foliation is expressed as $ds^2 = d\tau^2 - h_{ij}dx^i\, dx^j$, written in proper-time gauge with unit lapse and vanishing shift, assuming time-symmetric slices ($K_{ij} = 0$). In this configuration, the exact Hamiltonian constraint of general relativity reduces to a curved-space Poisson equation $\nabla_h^2\phi_G = 4\pi G\mu$, where ∇_h^2 is the Laplacian constructed from the spatial metric h_{ij}. The scalar ϕ_G determines the conformal factor of the spatial metric h_{ij}, and solving this constraint provides the geometric potential that influences the evolution of $\varphi(\tau,\vec{x})$.

The world-block operates in an extremely weak, non-radiative regime, with $|\phi_G|/c^2 < 10^{-20}$ for $\sigma > 10\ nm$. In this regime, higher-order GR terms are entirely negligible for all proposed tests. The coupling between matter and geometry is fully consistent with GR. The proper-time Schrödinger dynamics govern the evolution of the matter state, while the Hamiltonian constraint simultaneously updates the spatial metric on each proper-time slice. Therefore, Equation (30) represents the exact Hamiltonian constraint for the slices, not a severe approximation.

To derive the Madelung equations (the continuity and the Hamilton-Jacobi equation) from the given Schrödinger-Newton-like Equation (30) we substitute the polar representation $\varphi = \sqrt{\rho}e^{iS/\hbar}$ of the wavefunction, into the Schrödinger-Newton Equation (30) and then separate it into its real and imaginary parts.

The imaginary part results in the continuity equation:

$$\partial_\tau\rho + \vec{\nabla}\cdot\left(\rho\,\vec{\nabla}S/m\right) = 0 \tag{31}$$

On the other hand, the real part of Equation (30) leads to the following Hamilton-Jacobi equation:

$$\partial_\tau S + \frac{1}{2m}\left(\vec{\nabla}S\right)^2 + Q + \frac{3B^2}{2A}[\varphi_C + h_H] = 0 \tag{32}$$

where the quantum potential is given by:

$$Q(\vec{x},\tau) = -(\hbar^2/2m)\,\nabla^2\sqrt{\rho}/\sqrt{\rho} \tag{33}$$

Nonlocality is inbuilt in the intrinsic metric, as manifested by the equations of continuity and Hamilton-Jacobi. In fact, when the quantum potential term Q is expressed in terms of h_{ij}, it involves a nonlocal convolution. Consequently, the potential Q, and therefore the dynamics of the phase S, depend nonlocally on the distribution of the density ρ, and is non-separable across the slice, reflecting the intrinsic, geometry-induced non-locality of the world-block.

To determine the coefficients A and B, we first observe that the dimension of the metric gradient term in the Lagrangian Equation (26) is $\left[D_k h_{ij} D^k h^{ij}\right] = L^{-2}$. Consequently, the coefficient A must have the dimension of energy per unit length. Furthermore, A determines the stiffness of the metric, making it reasonable to choose A as the particle's rest energy mc^2 per Compton wavelength λ_c:

$$A = mc^2/\lambda_c = m^2c^3/\hbar \tag{34}$$

To replicate the standard Newton-Schrodinger kernel, and given the above choice for A, we pose the following expression for the coefficient B:

$$B = m\sqrt{8\pi AG} = m^2\sqrt{8\pi Gc^3/\hbar} \tag{35}$$

For simplicity, we retain the symbols A and B to track the algebraic factors.

5. Two-phase intrinsic metric

Quantum dispersion on each slice S_τ is driven by the gradient term in the Lagrangian. In the Madelung representation, this produces the quantum-pressure potential $Q(\tau,\vec{x})$, whose force $-\vec{\nabla}Q(\tau,\vec{x})$, drives the outward spreading of the density distribution.

When left undisturbed, a free world-block continues to expand or dilate from slice to slice. However, the introduction of a detector reverses this trend by imposing a harmonic contribution h_H to the intrinsic metric in the general solution:

$$\begin{gathered} h_{ij}(\tau,\vec{x}) = -\frac{B}{2A}\delta_{ij}[\varphi_C(\tau,\vec{x}) + h_H(\tau,\vec{x})] \\ \nabla_x^2\varphi_C = 4\pi\rho \;\; ,\nabla^2 h_H(\tau,\vec{x}) = 0 \quad \textit{outside the detector} \end{gathered} \tag{36}$$

This harmonic term $h_H(\tau,\vec{x})$ is directly affected by the presence of a detector. However, its Laplacian must remain zero outside the boundary of the detector's domain. This constraint means that the detector domain must be specifically defined as a small spherical region centred at $\vec{x}_0$, with a radius r_d beyond which the Poisson equation holds. This requirement can be enforced by introducing a Dirichlet boundary ∂D condition:

$$h_H(\tau,\vec{x}\in\partial D) = -I(\tau)C_0; \qquad 0 \le I(\tau) \le 1 \tag{37}$$

where $I(\tau)$ is a smooth transition function (or switch function), and C_0 is a constant.

The aforementioned condition should ensure that h_H remains harmonic everywhere except at the fixed detector surface. Specifically, the condition $\nabla_{\vec{x}}^2 h_H = 0$ should be satisfied within the region where the Poisson equation is applicable. This requirement can be accomplished by introducing a harmonic 'monopole well' located outside the boundary ∂D of the detector, defined as follows:

$$h_H(\tau,\vec{x}) = -I(\tau)\frac{C_0 r_d}{|\vec{x}-\vec{x}_0|}\delta_{ij}; \qquad |\vec{x}-\vec{x}_0| > r_d \tag{38}$$

The detector domain radius r_d defines the surface ∂D of the sphere around a detector focus $\vec{x}_0$. The constant $C_0 > 0$ is dimensionless and represents the strength of the harmonic monopole, chosen such that the post-collapse size is on the order of the radius r_d. Specifically, $C_0 r_d$ corresponds to the depth of the monopole well, which is fixed by the detector's active potential: r_d is the detector's effective capture radius and C_0 (where $C_0{\sim}1$) is the bias energy required to register one charge. The switch $I(\tau)\in[0,1]$ modulates the amplitude of the harmonic monopole.

Equation (38) demonstrates that the presence of the detector focus $\vec{x}_0$ imposes a Dirichlet-type boundary condition on the intrinsic metric h_{ij} by constraining it on a small surface. When the detector is idle, the switch $I(\tau)$ is at its lowest level ($I(\tau) = 0$), and no constraints are applied. Conversely, when the detection ramp generated by the detector reaches a high level $I(\tau) \approx 1$, the boundary ∂D enforces the intrinsic metric h_{ij} to conform to the monopole form.

It is important to note that the switch function $I(\tau)$ cannot take the form of a simple step function. A sudden jump (δ-spike) would result in an infinite time derivative, which would correspond to infinite energy. Instead, it must exhibit a smooth profile, transitioning gradually from 0 to 1 over a finite proper-time interval $\Delta\tau$.

A convenient smooth profile is the logistic form:

$$I(\tau) = 1/\left(1 + e^{-(\xi(\tau)-1)/\varepsilon}\right), \;\; \xi \ll 1 \tag{39}$$

where $\xi(\tau)$ represents the driving ratio between the pulse or ramp applied by the detector and the quantum pressure:

$$\xi(\tau) = \frac{V_{det}(\tau)}{\hbar^2/(mL^2)} \tag{40}$$

The numerator $V_{det}(\tau)$ represents the detector energy scale, which increases as the ramp progresses. The denominator $\hbar^2/(mL^2)$ corresponds to the quantum pressure of a free wave packet with a size $L(\tau)$ prior to any interaction with the detector.

For $I(\tau) \approx 0$, the quantum-pressure dominates, reaching its minimum at infinity. Consequently, as long as the detector remains weak, with $\xi(\tau) < 1$, the wave packet continues to expand freely, following Schrödinger-Newton-like dynamics.

A critical threshold time τ_c is reached when the logistic switch $I(\tau)$ crosses its midpoint $I(\tau_c) = 1/2$, which corresponds to $\xi(\tau_c) = 1$. At this moment, the detector energy $V_{det}(\tau_c)$ becomes equal to the quantum-pressure $\hbar^2/(mL^2)$. This relationship provides a means to estimate the collapse window $\Delta\tau$, as detailed in the next Section.

Once the driving ratio $\xi(\tau)$ surpasses the critical threshold ($\xi(\tau) = 1$), a second global minimum emerges within the monopole well. When $I(\tau) \approx 1$, the quantum-pressure is significantly suppressed, allowing the detector energy to dominate. As a result, the wave packet collapses into the monopole well, whose depth is $C_0 r_d$.

Putting together the free and monopole contributions results in an interpolating kernel that effectively transitions between the two regimes:

$$K(\tau; \vec{x}, \vec{y}) = \left(1 - I(\tau)\right)\frac{1}{4\pi|\vec{x} - \vec{y}|} - I(\tau)\frac{C_0 r_d}{|\vec{x} - \vec{x}_0|} \tag{41}$$

The interpolating kernel $K(\vec{x}, \vec{y}; \tau)$ serves as an intrinsic solver, redistributing the 'geometric weight' through the weighting factor $I(\tau)$. This kernel transitions smoothly between the non-local Coulomb-type free kernel K_{free} (free spreading when $I(\tau) \approx 0$) and the sharply localised kernel K_{loc} (monopole well when $I(\tau) \approx 1$). Importantly, the inclusion of $I(\tau)$ in the interpolating kernel is not an additional postulate but a natural result of enforcing the detector boundary condition while maintaining the solution to the same Poisson-like equation.

Both kernels K_{free} and K_{loc} are Green functions of the same Laplace–Beltrami operator $\mathrm{D}^2 h_{ij}$ that appears in the Hamiltonian constraint of Equation (27). The Coulomb free kernel K_{free} is the free-space solution, acting as the inverse of $\mathrm{D}^2 h_{ij}$ on an unbounded slice. The localised kernel K_{loc}, on the other hand, is the solution in the presence of a point-source image placed just outside the detector boundary.

Both kernels are inherently spatial and act within the intrinsic metric, derived from the same variational action (mass term + curvature term) that leads to Equation (27). They are not ad hoc constructions but natural outcomes arising consistently from the intrinsic geometric framework.

The intrinsic metric throughout the collapse window for the world block can thus be expressed as follows:

$$h_{ij}(\tau,\vec{x}) = -\frac{B}{2A}\delta_{ij}\int d^3y\rho(\tau,\vec{y})K(\tau;\vec{x},\vec{y}) \tag{42}$$

The integration of the density $\rho(\tau,\vec{y})$ over the entire spatial domain of the world-block signifies the inherent non-separability between all its points. The two-phase behaviour of the intrinsic metric is controlled by the smooth measuring switch function $I(\tau)$, which dynamically transitions the metric between the free phase and the localising phase within the collapse window $\Delta\tau$.

For $I \approx 0$, the kernel K reduces to the non-local Coulomb form, and the metric remains in the free phase. Conversely, for $I \approx 1$, only the localising monopole term persists, and the metric enters the collapsed phase. Since the interpolating kernel K incorporates both the density ρ and the switch function $I(\tau)$, this transition represents an intrinsic geometric response—not an external rule—and remains compatible with the Poisson constraint, the continuity equation, and the Hamilton–Jacobi dynamics on every slice.

6. Measurement-induced localisation

Assume that the initial intrinsic wave packet is described by a free Gaussian with an initial width σ_0. At $\tau = 0$, let:

$$\varphi_0(\vec{x}) = \left(\frac{1}{\pi\sigma_0^2}\right)^{3/4} e^{-|\vec{x}|^2/2\sigma_0^2}; \qquad \rho = |\varphi_0|^2 \tag{43}$$

Assuming the wave packet remains Gaussian, with $\sigma(\tau)$ evolving due to the variational energy contributions from quantum pressure and the detector-induced monopole well potential. The variational energy $E(\sigma, I)$ of the Gaussian packet, as a function of its width σ and the switch $I(\tau)$, is expressed as:

$$E(\sigma, I) = \frac{3\hbar^2}{4m\sigma^2} - I(\tau)\frac{2V_0}{\sqrt{\pi}\sigma}; \quad V_0 = \frac{3B^2C_0r_d}{2A} \tag{44}$$

The first term of Equation (44) represents the gradient energy of the amplitude, commonly referred to as quantum pressure. The second term corresponds to the monopole well potential. The first term is positive and promotes spreading, as the energy decreases for larger σ. Conversely, the second term is negative and promotes contraction, as the energy decreases for smaller σ, however, this term is only present when the detector switch $I(\tau)$ is not zero.

When the detector is off, $I(\tau) = 0$, the variational energy $E(\sigma, I)$ has its global first minimum at $\sigma \to \infty$, causing the world-block to expand (free phase). As $I(\tau)$ increases from 0 to 1, the variational energy $E(\sigma, I)$ rises monotonically. Initially, when $I(\tau)$ is small, $E(\sigma, I)$ develops a shallow local minimum at a finite σ, indicating that the monopole attraction starts to compete with the gradient energy but does not yet dominate, maintaining the global minimum at infinity. Once $I(\tau)$ becomes sufficiently large, the finite σ well sinks below the $\sigma \to \infty$ plateau. This results in the emergence of a new global minimum at $\sigma \sim r_d$. Consequently, the world-block transitions from the free minimum to the confined minimum, leading to collapse.

More specifically, by substituting the Gaussian profile into the bulk Lagrangian Equation (26) and integrating over space, we derive the effective single-variable Lagrangian:

$$L_{eff}(\sigma,\dot{\sigma};\tau)=\frac{m}{2}\dot{\sigma}^2-\frac{3\hbar^2}{4m\sigma^2}+I(\tau)\frac{2V_0}{\sqrt{\pi}\sigma^2} \tag{45}$$

By computing the Euler Lagrangian derivative, $\partial L_{eff}/\partial\tau=\dot{I}2V_0/\sqrt{\pi}\sigma$, we derive the following dynamical equation:

$$m\ddot{\sigma}=\frac{3\hbar^2}{2m\sigma^3}-I(\tau)\frac{2V_0}{\sqrt{\pi}\sigma^2}+\dot{I}(\tau)\frac{2V_0}{\sqrt{\pi}\sigma} \tag{46}$$

The first two terms on the right-hand-side of Equation (46) represent the gradient of the effective potential associated with the variational energy $E(\sigma,I)$ of Equation (44). The final term accounts for the time-dependent driving force, which injects energy into the system during the ramping of the detector ($\dot{I}(\tau)\neq 0$).

In the adiabatic regime, where $\dot{I}(\tau)\approx 0$, the equilibrium width σ_{eq} can be determined by setting $\ddot{\sigma}=0$ in the Equation (46). This leads to the following expression for the equilibrium width, σ_{eq}:

$$\sigma_{eq}(\tau)=\frac{3\hbar^2\sqrt{\pi}}{4mV_0I(\tau)} \tag{47}$$

For $I(\tau)\to 0$, the equilibrium width $\sigma_{eq}(\tau)\to\infty$, indicating a free spreading regime.

The onset of collapse occurs when $\sigma_{eq}(\tau)$ becomes comparable to the free spreading width $L(\tau)$. Based on the definition of the logistic switch (39), the threshold instant τ_c is determined when the driving ratio $\xi(\tau_c)=1$. Using the relationship of Equation (40), this implies:

$$V_{det}(\tau_c)=V_c=\hbar^2/\big(mL^2(\tau_c)\big) \tag{48}$$

To estimate the collapse window $\Delta\tau$, we linearise $V_{det}(\tau)$ around the threshold instant τ_c as:

$$V_{det}(\tau)\approx V_c+s(\tau-\tau_c) \tag{49}$$

where $s=[V_{det}/d\tau]_{\tau_c}$ is the slope of $V_{det}(\tau)$ at the threshold instant τ_c.

Substituting $V_{det}(\tau)$ into the expression for $\xi(\tau)$, and then differentiating with respect to τ, one finds:

$$d\tau=[\hbar^2/(mL^2s)]d\xi \tag{50}$$

By selecting a practical criterion where the collapse window is defined as the time interval during which the driving ratio $\xi(\tau)$ transitions from 10% to 90%, one can estimate the collapse window $\Delta\tau$:

$$\Delta\tau=4.4\,\varepsilon\hbar^2/(mL^2s) \tag{51}$$

This provides an estimate for the duration of the collapse process based on the rate of change of the driving potential. For example, considering an electron with a free width $L=100\ nm$, a ramp slope of $s=10^{-4}\ eV/ps$, and a small dimensionless parameter $\varepsilon=0.01$, the collapse duration is estimated as $\Delta\tau\sim 7ps$.

To analyse the behaviour of the system within the collapse window, i.e. for $|\tau-\tau_c|<\Delta\tau$, we simplify the governing Equation (46) by neglecting the $1/\sigma^3$ term once σ has shrunk to a value smaller than $1/L^2$. Under this approximation, Equation (46) can then be integrated to yield:

$$\sigma(\tau)\approx\frac{\sigma_0}{\sqrt{1+(\tau-\tau_c)/\Delta\tau}}\ ;\ \tau_c\le\tau\le\tau_c+3\Delta\tau \tag{52}$$

So the width decreases approximately as the inverse square-root of the elapsed proper time, scaled by the collapse window $\Delta\tau$. After three logistic e-folds, the collapse process is nearly complete, and the width reaches:

$$\sigma_\infty \approx \frac{\sigma_0}{\sqrt{1+3}} = \frac{\sigma_0}{2} \tag{53}$$

The with σ_∞ asymptotically approaches r_d as $I \rightarrow 1$.

A notable feature of the finite collapse window $\Delta\tau$ is the fact that if the detector is abruptly switched off (i.e. $I \rightarrow 0$) while the system is still within or near the collapse window, (i.e. at $\tau_{off} < \tau_c + \Delta\tau$), then Equation (46) reduces to describing the dynamics driven solely by quantum pressure, leading to:

$$m\ddot{\sigma} = \frac{3\hbar^2}{2m\sigma^3} \tag{54}$$

Integrating Equation (54) twice yields the following expression for the width σ:

$$\sigma^2(\tau) = \sigma_\infty^2 + \frac{3\hbar^2}{m^2\sigma_\infty^2}\left(\tau - \tau_{off}\right)^2 \tag{55}$$

This result illustrates the free re-expansion of the wave packet's width driven by quantum pressure after the detector is abruptly turned off, leading to the possible reappearance of interference patterns. This behaviour is a distinctive signature of the finite-window collapse process and stands in contrast to instantaneous projection scenarios, where such re-expansion and the revival of interference patterns do not occur.

The revival takes place only if the detector bias is turned off before the logistic switch reaches a value of approximately $I(\tau) \approx 0.9$. In the case of an electron, as described in the above example, the revival happens within a proper-time interval that is shorter than the collapse window, $\Delta\tau \approx 7\ ps$.

7. Example of the double-slit experiment

In the double-slit experiment a single quantum particle—for example, an electron—approaches an opaque plate with two narrow slits. The particle's wave packet, upon reaching the plate, is constrained to pass through the slits. The following analysis uses proper time to explore how the particle's intrinsic wave packet splits, evolves, and ultimately generates the well-known interference pattern, all while the particle's localized charge follows one of the resulting lobes.

Before interacting with the plate, the particle is solely characterised by its intrinsic wave packet $\varphi(\tau, \vec{x}) = \sqrt{\rho}e^{iS/\hbar}$ on each slice S_τ, valid for for $\tau < \tau_P$, where τ_P denotes the proper time at which the particle encounters the plate. Prior to reaching the plate, the intrinsic wave packet takes the form of a single Gaussian centred at $\vec{x}_C(\tau)$ and evolves according to free dynamics. The squared magnitude ρ of the intrinsic wave packet represents the mass-density distribution, while the centroid world-line $\vec{x}_C(\tau)$ carries the particle's charge and remains within the support of the intrinsic wave packet.

At the proper time τ_P, the particle interacts with the plate, positioned on the intrinsic 3-surface S_{τ_P}. At $\tau = \tau_P$, the plate imposes a boundary condition on the wave packet with a potential V_M. Specifically, V_M is set to zero within two narrow windows W_A and W_B, centred at $\vec{x}_A$ and $\vec{x}_B$, respectively, while it is very large everywhere else. This interaction causes the particle's wave packet to split as it propagates through the windows defined by the slits.

Multiplying the intrinsic wave packet φ by the respective window indicator functions Θ_A and Θ_B gives, immediately after the plate:

$$\varphi(\tau_P^+,\vec{x}) = \varphi(\tau_P^-,\vec{x})[\theta_A + \theta_B] = \varphi_A + \varphi_B \tag{56}$$

Thus, immediately after interacting with the mask, the single intrinsic wave packet $\varphi(\tau_P^+,\vec{x})$ develops two amplitude peaks (lobes) corresponding to the two windows. Despite this, the intrinsic wave-packet $\varphi(\tau_P^+,\vec{x})$ remains a single, complex function, although its support is now split into two distinct lobes. Importantly, phase coherence is preserved across the two lobes since both components inherit the same global phase of the incoming wave-packet. The reference charge world-line $\vec{x}_C(\tau)$ is entirely confined within one of the lobes, either φ_A or φ_B.

For $\tau > \tau_P$, the intrinsic wave-packet resumes free evolution. At the detector slice, defined by τ_D, the wave-packet is described as:

$$\varphi(\tau_D,\vec{x}) = \varphi_A(\tau_D,\vec{x}) + \varphi_B(\tau_D,\vec{x}) \tag{57}$$

where each component evolves independently under the same free propagation dynamics. However, due to differences in optical path lengths, the two components acquire a relative phase, such that:

$$\varphi_B(\tau_D,\vec{x}) = e^{i\Delta\alpha(\vec{x})}\varphi_A(\tau_D,\vec{x}) \tag{58}$$

The world-line $\vec{x}_C(\tau)$ carrying the particle's charge, remains confined within the same initial lobe that was determined after the interaction with the mask.

At the detector, the mass-density distribution of the intrinsic wave-packet is given by:

$$|\varphi(\tau_D,\vec{x})|^2 = |\varphi_A|^2 + |\varphi_B|^2 + 2Re[\varphi_A^*\varphi_B] \tag{59}$$

The final term represents the intrinsic cross-term responsible for interference, which remains non-zero because the off-diagonal element of the following density matrix $\rho(\tau_D;\vec{x},\vec{x}')$ maintains coherence between the two lobes:

$$\rho(\tau_D;\vec{x},\vec{x}') = \varphi^*(\tau_D,\vec{x}')\varphi(\tau_D,\vec{x}) \tag{60}$$

Therefore, the interference pattern, manifested as regions of brightness and darkness, is inherently encoded within the space of the intrinsic wave-packet itself.

When the particle is detected during a single experimental run, its trajectory ends at a specific absorber located at $\vec{x}_{hit}\epsilon\{\varphi_A\ or\ \varphi_B\}$, corresponding to the lobe within which it resides. This results in the collapse of the intrinsic wave-packet φ, as described in Section 7.

Across many identical experimental runs, the initial placement of $\vec{x}_C(\tau_P^+)$ is statistically sampled from the distribution $|\varphi(\tau_P^+)|^2$. Since this distribution already encompasses the interference pattern, the histogram of hit positions converges to a profile proportional to $|\varphi_A + \varphi_B|^2$, perfectly matching the observed interference fringes.

8. Configuration Space

A single, free quantum world-block b, described by the intrinsic coordinates $(\vec{x},\tau)$, is defined within a Hilbert space $\mathcal{H} = span\{|\vec{x}\rangle\}$. Its state $|\varphi\rangle$ is expressed as follows:

$$|\varphi\rangle = \int d^3x\varphi(\vec{x},\tau)\,|\vec{x}\rangle \tag{61}$$

More generally, a system consisting of N particles is conceptualized as a collection of N world-blocks. To account for the individuality of each world-block, the system of N blocks is represented in an abstract space $\mathcal{S}$ (or manifold), which is defined as the tensor product of the N individual blocks within the Hilbert space $\mathcal{H}^{(N)}$:

$$\mathcal{S} = b_1 \otimes b_2 \otimes \cdots \otimes b_N \tag{62}$$

where the Hilbert space $\mathcal{H}^{(N)}$ is itself the tensor product of the N single-block Hilbert spaces:

$$\mathcal{H}^{(N)} = \mathcal{H}_1 \otimes \mathcal{H}_2 \otimes \cdots \otimes \mathcal{H}_N \tag{63}$$

The generic case involves one proper time per block. A single world-block is charted by $(\vec{x}_i, \tau_i)$, where τ_i represents the proper time measured along the reference world-line anchoring block b_i, and $\vec{x}_i$ lies on the space-like slice $S^{(i)}_{\tau_i}$. Proper time along a world-line is uniquely defined up to an additive constant and a scale factor, both determined by the block's motion relative to its environment. For two blocks b_1 and b_2, which exhibit different modes of motion, their proper times τ_1 and τ_2 accumulate at different rates and, consequently, cannot generally be aligned slice-by-slice.

In the case of an N-block system, the configuration space is characterized by coordinates in $R^{3N} \times R^N$. Thus, the system involves N proper-time parameters:

$$(\vec{x}_1, \tau_1; \vec{x}_2, \tau_2; \cdots; \vec{x}_N, \tau_N) \,\epsilon\, R^{3N} \times R^N \tag{64}$$

This system can also be described by a wavefunction that encapsulates the entire ensemble. Consequently, the wavefunction of the N-particle system is defined in the 4N-dimensional configuration space $R^{3N} \times R^N$. The general state of the system is thus represented as a multi-time wavefunction Φ:

$$\Phi(\vec{x}_1, \tau_1; \vec{x}_2, \tau_2; \cdots; \vec{x}_N, \tau_N) \tag{65}$$

We can work with the N proper-time parameters. However, it is also possible to adopt the standard formalism of "Tomonaga–Schwinger gauge'' [24] to introduce a global time parameter for the N-block system, replacing the N individual proper times. This involves selecting a monotonic real parameter λ with $\dot{\lambda} > 0$. For instance, λ can be chosen as the proper time of one reference world-block. For each individual world-block b_i, we then define a smooth, strictly increasing clock map:

$$\tau_i = \tau_i(\lambda)\,; \;\; \dot{\tau}_i(\lambda) = d\tau_i/d\lambda > 0 \tag{66}$$

The global parameter λ is used solely to label composite slices $\left(S^{(1)}_{\tau_1(\lambda)}, \cdots, S^{(N)}_{\tau_N(\lambda)}\right)$ corresponding to the N world-blocks. Importantly, this formalism does not assume a common tick rate for the proper times of the individual blocks.

The configuration space is thus reduced to $R^{3N} \times R$, where all the world-blocks are evaluated on the same global slice defined by $\lambda = const$. Consequently, the multi-time wavefunction ϕ and the configuration-space density ρ can be redefined as follows:

$$\begin{gathered} \varphi(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda) = \phi\big(\vec{x}_1, \tau_1(\lambda); \vec{x}_2, \tau_2(\lambda); \cdots; \vec{x}_N, \tau_N(\lambda)\big) \\ \rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda) = |\varphi|^2 \end{gathered} \tag{67}$$

For a single world-block, the wavefunction's density $\rho(\vec{x}, \tau) = |\varphi|^2$and the mass density $\mu(\vec{x}, \tau)$ are equivalent up to a proportionality constant:

$$\mu(\vec{x}, \tau) = m\rho(\vec{x}, \tau) \tag{68}$$

However, for a system of N world-blocks, it is necessary to distinguish between the global configuration-space density $\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N)$, defined in the configuration space R^{3N}, and the total mass density $\mu(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N)$, defined in ordinary physical space R^3. While the mass-density of each block and the total mass-density can be derived from the global configuration-space density, the reverse derivation is generally not possible.

The total mass-density operator $\hat{\mu}(\vec{x}) = \hat{\mu}(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N)$, defined in ordinary space R^3, for N world-blocks of masses m_k is:

$$\hat{\mu}(\vec{x}) = \sum_{k=1}^{N} m_k \delta^{(3)}(\vec{x} - \hat{x}_k) \tag{69}$$

where $\hat{x}_k$ represents the position operator acting on the world-block b_k

For any normalised wavefunction $\varphi(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda)$, with configuration-space density $\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda) = |\varphi|^2$, the total mass-density in ordinary space can be expressed in terms of the global configuration-space density as follows:

$$\mu_{tot}(\vec{x}, \lambda) = \langle \varphi | \hat{\mu}(\vec{x}) | \varphi \rangle = \sum_{k=1}^{N} m_k \rho_k^{(mar)}(\vec{x}, \lambda) \tag{70}$$
$$\mu_{tot}(\vec{x}, \lambda) = \sum_{k=1}^{N} m_k \int d^{3N}x \, \rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N) \delta^{(3)}(\vec{x} - x_k)$$

where $\rho_k^{(mar)}(\vec{x}, \lambda)$ is the marginal density associated with world-block b_k. The marginal density is defined in terms of the global configuration-space density ρ, as:

$$\rho_k^{(mar)}(\vec{x}, \lambda) = \int d^{3N}x \, \rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N) \delta^{(3)}(\vec{x} - x_k) \tag{71}$$

In essence, $\rho_k^{(mar)}(\vec{x}, \lambda)$ is obtained by performing the 3N-3 integrals over all coordinates $\vec{x}_l \neq \vec{x}_k$, effectively marginalizing over the degrees of freedom associated with all other world-blocks.

Equation (70) demonstrates that the physical 3-space mass-density is the sum of the N marginal densities, each weighted by their respective masses m_k. No cross-terms appear in this expression because each Dirac delta function $\delta^{(3)}(\vec{x} - x_k)$ acts as a projection operator that isolates the diagonal part of the configuration-space density.

Phase correlations between world-blocks, while critical for two-point or higher-order observables, do not influence the one-body mass density. Even in cases of full quantum entanglement, the one-body marginal densities $\rho_k^{(mar)}(\vec{x}, \lambda)$ contain all of the information required to compute the real-space mass distribution. This is because mass is an additive, one-body observable. Although entanglement impacts multi-point correlations and the geometric structure of the configuration space, it does not hinder the recovery of the physical-space mass-density, which can still be obtained through the summation provided in Equation (70).

9. Two-phase intrinsic metric in configuration space

In configuration space, the two-phase intrinsic metric for the N-block system can be formalised, as follows:

$$h_{ij} = -\frac{B}{2A} \delta_{ij} \int d^{3N}y \rho(\vec{y}_1, \cdots, \vec{y}_N; \lambda) \prod_{k=1}^{N} K_k^{(I_k(\lambda))}(\vec{x}_k, \vec{y}_k; \lambda) \tag{72}$$

where the index k refers to the k^{th} world-block; $I_k(\lambda)$ is the per-block switch; and $K_k^{(I_k(\lambda))}$ represents the following intrinsic metric Kernel:

$$K_k^{(I_k(\lambda))}(\vec{x}_k, \vec{y}_k; \lambda) = [1 - I_k(\lambda)] K_k^{(0)}(\vec{x}_k, \vec{y}_k) + I_k(\lambda) K_k^{(1)}(\vec{x}_k) \tag{73}$$

where $K_k^{(0)}$ and $K_k^{(1)}$ are referred to as the free Poisson kernel and the harmonic monopole kernel, respectively, and are formulated as:

$$K_k^{(0)}(\vec{x}_k,\vec{y}_k) = \frac{1}{4\pi|\vec{x}_k - \vec{y}_k|} \; ; \; K_k^{(1)}(\vec{x}_k) = -\frac{C_{0k}r_{dk}}{4\pi|\vec{x}_k - \vec{x}_{0k}|} \tag{74}$$

where $\vec{x}_{0k}$ and r_{dk} are the detector focus and radius for world-block k; C_{0k} is the dimensionless depth (of order 1).

The smooth logistic switch $I_k(\lambda) \epsilon [0,1]$, undergoes a rise when the driving ratio $\xi(\lambda)$ crosses 1, where $\xi(\lambda) = V_{det,k}(\lambda)/\hbar^2/(m_k L_k^2)$.

The product kernel $\prod_{k=1}^{N} K_k^{(I_k)}(\vec{x}_k,\vec{y}_k)$ acts on $h_{ij}(\vec{x}_1,\cdots,\vec{x}_N;\lambda)$ with the configuration-space Laplacian $\sum_k \nabla_k^2$. For each k the Laplacian acting on $K_k^{(0)}$ produces $\delta^{(3)}(\vec{x}_k - \vec{y}_k)$. Acting on $K_k^{(1)}$ yields 0 because it is harmonic in $\vec{x}_k$. Therefore, we precisely obtain the required multi-block Poisson equation:

$$\left(\sum_{k=1}^{N} \nabla_k^2\right) h_{ij} = \frac{B}{2A}\delta_{ij}\rho(\vec{y}_1,\cdots,\vec{y}_N;\lambda) \tag{75}$$

For the free phase, where $I_k(\lambda) \approx 0$, the k[th] factor corresponds to the Posson kernel $K_k^{(0)}(\vec{x}_k,\vec{y}_k)$, and the metric consequently couples every point of block k to every other point.

In the localized (measurement) phase, where $I_k(\lambda) \approx 1$, the influence of the long-range kernel $K_k^{(0)}(\vec{x}_k,\vec{y}_k)$ vanishes, leaving only the block's local monopole kernel $K_k^{(1)}(\vec{x}_k)$. This results in a confining metric well that is centred on the detector focus $\vec{x}_{0k}$, forcing the marginal density for block k to contract toward r_{dk}.

For an intermediate $I_k(\lambda)$ between the two phases, both kernels $K_k^{(0)}(\vec{x}_k,\vec{y}_k)$ and $K_k^{(1)}(\vec{x}_k)$ contribute continuously. The collapse window is characterized by $\Delta\lambda \approx \varepsilon\hbar^2/(m_k L_k^2 s_k)$, where s_k represents the detector ramp slope.

10. Dependent and independent world-blocks

In the case where the wavefunction of N world-blocks, factorises as:

$$\varphi(\vec{x}_1,\vec{x}_2,\cdots,\vec{x}_N,\lambda) = \prod_{k=1}^{N} \varphi_k(\vec{x}_k,\lambda) \tag{76}$$

Then, the global configuration-space density $\rho(\vec{x}_1,\vec{x}_2,\cdots,\vec{x}_N,\lambda)$ can correspondingly be expressed as the product of the individual densities:

$$\rho(\vec{x}_1,\vec{x}_2,\cdots,\vec{x}_N,\lambda) = \prod_{k=1}^{N} \rho_k(\vec{x}_k,\lambda) \; ; \rho_k(\vec{x}_k,\lambda) = |\varphi_k|^2 \tag{77}$$

In this case, the marginal density corresponds exactly to the individual density for each block. Specifically, integrating over the coordinates of all blocks except k yields exactly $\rho_k(\vec{x}_k,\lambda)$:

$$\begin{aligned} \rho_k^{(mar)}(\vec{x},\lambda) &= \int d^3\vec{x}_1,\cdots,d^3\vec{x}_{k-1},d^3\vec{x}_{k+1},\cdots,d^3\vec{x}_N \; \rho(\vec{x}_1,\vec{x}_2,\cdots,\vec{x}_N) \\ \rho_k^{(mar)}(\vec{x},\lambda) &= \rho_k(\vec{x}_k,\lambda)\prod_{l\neq k}\int d^3\vec{x}_l \; \rho_l(\vec{x}_l,\lambda) = \rho_k(\vec{x}_k,\lambda) \end{aligned} \tag{78}$$

With this product density, the Poisson equation for the metric decouples additively, and consequently, Equation (72) factorises into a direct sum of single-block metrics:

$$h_{ij}(\vec{x}_1, \cdots, \vec{x}_N, \lambda) = \sum_{k=1}^{N} h_{ij}^{(k)}(\vec{x}_k, \lambda) \tag{79}$$
$$\text{where } h_{ij}^{(k)}(\vec{x}_k, \lambda) = -\frac{B}{2A}\delta_{ij} \int d^3y_k \rho_k(\vec{y}_k, \lambda)\, K^{I_k(\lambda)}(\vec{x}_k, \vec{y}_k)$$

with each $h_{ij}^{(k)}$ having its own monopole switch.

Thus, if the individual densities are uncorrelated (i.e. total density $\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda)$ can be expressed as the product of the individual densities), then the intrinsic metrics are separable, and the N blocks are independent.

As a result, each world-block collapses or expands independently based on its own monopole switch $I_k(\lambda)$. An operation on world-block k does not affect world block $l \neq k$.

On the other hand, if the wavefunction of N world-blocks does not factorise, then the global configuration-space density $\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda)$ cannot be expressed as a product of individual densities:

$$\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda) \neq \prod_{k=1}^{N} \rho_k(\vec{x}_k, \lambda) \tag{80}$$

In this case, it is not possible to reconstruct the global configuration-space density, $\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda)$, from the set of marginal densities $\left\{\rho_k^{(mar)}(\vec{x}, \lambda)\right\}_{k=1}^{N}$. This is because the marginal density $\rho_k^{(mar)}(\vec{x}, \lambda)$ associated with block k does not contain any information about its correlation with the other blocks. Consequently, the integral (78) couples all coordinates, and therefore, cannot be factorised into a direct sum of single-block metrics:

$$h_{ij}(\vec{x}_1, \cdots, \vec{x}_N, \lambda) \neq \sum_{k=1}^{N} g_{ij}^{(k)}(\vec{x}_k, \lambda) \tag{81}$$

Correlations in the density $\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda)$ create a geometric interdependence among the N world-blocks. A measurement on any world-block b_i (flipping $I_i(\lambda)$ from 0 to 1), alters the convolution weight, subsequently modifying the global intrinsic metric $h_{ij}(\vec{x}_1, \cdots, \vec{x}_N, \lambda)$, as a function of all the coordinates $\vec{x}_k$. This modification represents the intrinsic signature of entanglement, yet the evolution remains unitary in λ.

Thus, if the individual densities are correlated (i.e. the global density $\rho(\vec{x}_1, \vec{x}_2, \cdots, \vec{x}_N, \lambda)$ cannot be expressed as a product of the individual densities), the intrinsic metrics become inseparable, and the N world-blocks are inherently interdependent.

11. Entangled world-blocks in the intrinsic–extrinsic framework

We consider two entangled world-blocks (labelled b_1 and b_2) and model a position-selective measurement performed exclusively on block b_1. The reasoning can be generalized to an arbitrary number of world-blocks.

For two world-blocks, the full quantum state is described by $\varphi(\vec{x}_1, \vec{x}_2; \lambda)$, and the configuration-space density is given by $\rho_{12}(\vec{x}_1, \vec{x}_2; \lambda) = |\varphi|^2$, where λ serves as the global clock.

In the case of entanglement, the state φ cannot be factorised into a product of separate wavefunctions for each block. This, equivalently, implies that $\rho_{12} \neq \rho_1\rho_2$, where ρ_1 and ρ_2 correspond to the marginal densities of blocks b_1 and b_2, respectively.

When both blocks are in the free phase (i.e., $I_1 = I_2 = 0$), the intrinsic configuration-space metric is determined by the product kernel:

$$h_{ij}(\vec{x}_1, \vec{x}_2, \lambda) = -\frac{B}{2A}\delta_{ij}\int d^3y_1 d^3y_2 \frac{\rho_{12}(\vec{y}_1, \vec{y}_2; \lambda)}{4\pi|\vec{x}_1 - \vec{y}_1|4\pi|\vec{x}_2 - \vec{y}_2|} \tag{82}$$

Each point in block b_1 is geometrically linked to all points in block b_2 through the double non-local Coulomb kernel.

When the position-selective detector is activated for block b_1, with a focus at $\vec{x}_{01}$ and radius $\vec{r}_{d,1}$, the corresponding logistic switch $I_1(\lambda)$ is given by:

$$I_1(\lambda) = 1/\left(1 + e^{-(\xi_1(\lambda)-1)/\varepsilon}\right); \quad \xi_1(\lambda) = V_{det,1}(\lambda)/(\hbar^2/(m_1 L_1^2)) \;\; ; \varepsilon \ll 1 \tag{83}$$

Block b_2 remains unaffected, and its logistic switch is fixed at: $I_2(\lambda) = 0$. Thus, the per-block kernels are given by:

$$K_1^{(I_1(\lambda))}(\vec{x}_1, \vec{y}_1; \lambda) = \frac{[1 - I_1(\lambda)]}{4\pi|\vec{x}_1 - \vec{y}_1|} - \frac{I_1(\lambda) C_{01} r_{d,1}}{|\vec{x}_1 - \vec{x}_{01}|}; \quad K_1^{(I_2(\lambda))} = K_2^{(0)}(\vec{x}_2, \vec{y}_2) = \frac{1}{4\pi|\vec{x}_2 - \vec{y}_2|} \tag{84}$$

The kernel $K_1^{(I_1(\lambda))}$ of block b_1 is a combination of the long-range Coulomb-like kernel $K_1^{(0)}(\vec{x}_1, \vec{y}_1) = 1/4\pi|\vec{x}_1, \vec{y}_1|$ and the harmonic monopole kernel $K_1^{(1)}(\vec{x}_1) = -C_{01} r_{d1}/4\pi|\vec{x}_1 - \vec{x}_{01}|$. In contrast, the kernel $K_1^{(I_2(\lambda))}$ of block b_2 is composed only of the long-range Coulomb-like kernel $K_2^{(0)}$.

The intrinsic metric during the collapse window is given by:

$$h_{ij}(\vec{x}_1, \vec{x}_2, \lambda) = -\frac{B}{2A}\delta_{ij}\int d^3y_1 d^3y_2 \rho_{12}(\vec{y}_1, \vec{y}_2; \lambda) K_1^{(I_1)}(\vec{x}_1, \vec{y}_1; \lambda) K_2^{(0)}(\vec{x}_2, \vec{y}_2) \tag{85}$$

As $I_1(\lambda)$ increases during the collapse window $\Delta\lambda = \varepsilon\hbar^2/(m_1 L_1^2 s_1)$, the long-range Coulomb kernel $K_1^{(I_1(\lambda))}$ of block b_1 gradually loses weight, while the harmonic monopole kernel, proportional to $1/|\vec{x}_k - \vec{x}_{01}|$ gains weight.

Block b_2, however, continues to experience a time-dependent Coulomb-like weight $\left(1 - I_1(\lambda)\right) K_1^{(0)}$ of block b_1 in its effective potential until the collapse window $\Delta\lambda$ ends. After the collapse window closes, the product kernel retains only the monopole factor $K_1^{(1)}$from block b_1 and the Coulomb factor $K_2^{(0)}$ from block b_2. The coupling weight between the two blocks decreases from its maximum value to nearly zero over the duration of the collapse window $\Delta\lambda$, signalling the onset of decoherence.

12. Experimental signatures

The proposed interpretation of quantum mechanics introduces several experimental predictions that distinguish it from other interpretations. A key prediction is the existence of a finite collapse window $\Delta\tau \approx \varepsilon\hbar^2/(mL^2 s)$ with $\varepsilon \ll 1$. For electrons, the collapse time is approximately 10 ps, while for heavy atoms, it is around 100 ns. This finite collapse duration is a central feature of the model.

Closely related to this finite collapse time is the prediction of wave-packet re-expansion and the revival of interference patterns if the detector pulse ends before full localization is achieved. This phenomenon is not accounted for by interpretations based on the projection postulate or by spontaneous collapse models with a constant rate. For example, in an electron double-slit interferometer, placing a 50 nm radius electrostatic "tip" just behind one slit and pulsing it for about 20 ps results in a visibility decay that follows the logistic curve. Furthermore, the model predicts that a second, delayed pulse can recover part of the lost contrast, meaning the interference fringes reappear if the detector is turned off before the collapse time $\Delta\tau$ is complete. This behavior is fundamentally inconsistent with the instantaneous projection interpretation, which does not allow for fringe revival.

Another distinct feature is the mass dependence of the collapse time, where localization begins only when the detector potential satisfies $V_{det} \geq \hbar^2/(mL^2)$. As a result, lighter quantum wave-packets require a stronger detector energy for localization compared to heavier ones.

Additionally, the model predicts a gradual decay of Bell-type correlations between two entangled systems over the finite collapse window, as opposed to an abrupt and instantaneous jump.

The model also anticipates a transient, metric-induced force acting on nearby test masses. This effect manifests as a weak, prompt acceleration outside the detector's influence zone, synchronized with the collapse process. The magnitude of this acceleration is extremely small and is estimated to be of the order of $a(r) = 12\pi G m^2 C_0 r_d / r^2$, depending solely on the world-block's mass and geometry. Despite its small magnitude, this effect is theoretically measurable using optomechanical instruments. For example, a world-block represented by a mesoscopic molecule with a mass on the order of $m \sim 10^{-21} kg$ coupled with a micro-cantilever probe located approximately $r = 5\ \mu m$ from the detector's focal point (with a width of $r_d = 50\ nm$), would generate a small, distance-dependent acceleration of approximately $a \sim 10^{-14} ms^{-2}$. This acceleration coincides with the visibility drop occurring during the collapse window. The force acts only during the interval when $I(\tau)$ is rising. Notably, a cryogenic SiN cantilever exhibits a thermal noise-limited acceleration sensitivity of approximately $a_{th} \sim 10^{-15} ms^{-2}$, which is within an order of magnitude of the predicted signal. Furthermore, using optical interferometric readout over the duration of the collapse window ($\Delta\tau \sim 10ps - 10ns$), the transient acceleration signal could, in principle, be resolved.

These experimentally testable predictions provide a clear framework for differentiating this interpretation of quantum mechanics from others and offer a means for its potential falsification.

13. Conclusion

Re-examining Schrödinger's concept of a diffuse "cloud" through geometric tools suggests a new perspective in which a quantum entity is not a point-like particle residing in pre-existing spacetime, but rather an extended four-dimensional "world block" whose intrinsic geometry encodes the wavefunction. In this framework, the mass density spread across each proper time slice generates a metric, via a Poisson-like relation, that non-locally links all parts of the block. Measurement emerges when an external apparatus imposes a smooth, time-dependent boundary condition, distorting the Green kernel of this metric from its free Coulomb-like form into a sharply localized monopole well. Since this deformation occurs continuously over a finite "collapse window," localisation becomes a gradual geometric phase transition rather than a sudden, stochastic event.

The contraction of the mass distribution during the collapse window, remains finite, and can re-expand if the detector's pulse ends prematurely—allowing interference patterns to revive. This approach naturally integrates entanglement: multiple world blocks share a composite configuration space metric that only decouples when the overall density factorizes. Otherwise, a measurement on one block gradually diminishes the non-local Coulomb-like correlation with others, resulting in a smooth decay of Bell-type correlations.

This interpretation offers specific, testable predictions that depend on mass, including picosecond-scale localisation for electrons, sub-microsecond scales for larger molecules, logistic (rather than step-like) visibility loss, revivable interference fringes, and a transient ultra-weak acceleration of nearby probes during the collapse window. These effects are, in principle, measurable with current ultrafast interferometry and optomechanical sensing technologies.

By treating the wavefunction as the relational fabric of a particle's own spacetime, this interpretation provides a deterministic yet non-separable framework that reconciles quantum non-locality with relativistic causality, avoiding the need for hidden variables or many-worlds interpretations. Furthermore, it invites experimental verification through its unique dynamical and gravitational predictions.

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