

Quantisation of the Speed of Light in Black Holes: A Non-Perturbative Origin of the Universe and a Perturbative Resolution of the Information Paradox

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We develop a quantum-gravity framework in which, within extreme regimes such as black-hole interiors, the fundamental speed of light is promoted to a quantum operator that does not commute with spacetime, thereby inducing a quantised metric whose nonlinear back-reaction prevents the geometry from approaching physical singularities. Within this structure, an internal Hawking-like excitation—the *anti-eigen flux*—can tunnel outward through quantum-corrected horizons. We apply this framework to cosmology by tracing initial conditions to a primordial black hole whose fields saturate the lower bound of non-perturbative quantisation, enforcing a fundamental asymmetry between the tunnelling probabilities of the anti-eigen flux and standard Hawking radiation. This asymmetry drives the non-perturbative primordial black hole along a negentropic trajectory toward an intrinsically unstable Planck-scale black hole, which inevitably generates a quantum-gravitational jerk (recoil) of 2.175×10^{61} hawkings—an event we term *the Jerk Tremendous*—triggering a transition to classical behaviour, decoupling gravitational and quantum dynamics, and yielding an asymptotically free gravitational universe without a cosmological constant problem. In the perturbative regime of black hole quantisation, tunnelling probabilities become symmetric, restoring unitarity and ensuring the conservation of information.

Keywords: Primordial black holes; Noncommutative geometry; Hawking radiation; Information paradox; Singularities; Matter-antimatter asymmetry

1 Fundamental Framework

Status of Extreme Field Theory Extreme Field Theory (EFT), as developed in this work, is formulated as an effective operator-based framework designed to capture quantum-geometric corrections in extreme gravitational regimes—particularly within black hole interiors where semiclassical gravity proves inadequate. The formalism promotes selected geometric quantities to quantum operators, encoding back-reaction effects and correlations across perturbative and non-perturbative regimes, most notably allowing the speed of light to acquire an operator-valued character. Consequently, in the extreme-field regime the fundamental object of the Einstein field equations—the metric tensor—acquires an operator-valued structure, reflecting the quantisation of geometry at extreme curvature. EFT thus complements established quantum gravity programmes, offering a physically motivated intermediate description of extreme-field dynamics with implications for black hole interiors, the information paradox, and early-universe cosmology.

Motivation 1. *Recent observations by the James Webb Space Telescope [148, 149] suggest that our Universe may have originated within a primordial black hole [41, 131], challenging our understanding of black hole interiors and their role in cosmic evolution. Classical frameworks such as Newtonian mechanics, general relativity, and quantum field theory—while successful in many domains—remain incomplete when addressing singularities and the information paradox in extreme gravitational regimes [36, 53, 74, 79]. These limitations motivate the development of an extreme field theory, in which quantum-gravitational effects allow fundamental quantities, such as the speed of light, to acquire an operator-valued character [3, 13, 27, 37, 66, 164].*

Definition 1 (Extreme Quantum-Gravitational Electrodynamics). *In extreme quantum-gravitational regimes, electromagnetic field dynamics obey:*

$$\nabla_{\mu} F^{\mu\nu} = J_{\text{quantum}}^{\nu}, \quad (1)$$

where $F^{\mu\nu} = \nabla^{\mu} A^{\nu} - \nabla^{\nu} A^{\mu}$ is the Maxwell tensor, and J_{quantum}^{ν} represents a non-vanishing quantum source.

Proposition 1. *The field equations (1) yield a modified wave equation with non-trivial source terms that cannot be gauged away.*

Proposition 2 (Quantum Speed of Light). *In extreme quantum-gravitational regimes, the speed of light emerges as a quantum operator $\hat{C}^2 = c^2 + \hat{\chi}$.*

Proof. We begin with the modified Maxwell equation, Eq. (1), expressed in terms of the potential A^μ :

$$F^{\mu\nu} = \nabla^\mu A^\nu - \nabla^\nu A^\mu. \quad (2)$$

Computing the divergence:

$$\nabla_\mu F^{\mu\nu} = \nabla_\mu (\nabla^\mu A^\nu - \nabla^\nu A^\mu) = \nabla_\mu \nabla^\mu A^\nu - \nabla_\mu \nabla^\nu A^\mu. \quad (3)$$

Using the commutator identity [47, 99]:

$$[\nabla_\mu, \nabla^\nu] A^\mu = R^\nu{}_\lambda A^\lambda, \quad (4)$$

we obtain:

$$\nabla_\mu \nabla^\nu A^\mu = \nabla^\nu \nabla_\mu A^\mu + R^\nu{}_\lambda A^\lambda. \quad (5)$$

Generalising the Lorenz gauge condition [85] to include quantum-gravitational back-reaction:

$$\nabla_\mu A^\mu \neq 0. \quad (6)$$

Employing condition (6) in Eq. (5) and substituting into Eq. (3) yields:

$$\nabla_\mu F^{\mu\nu} = \square A^\nu - R^\nu{}_\lambda A^\lambda - \nabla^\nu \nabla_\mu A^\mu, \quad (7)$$

where $\square = \nabla_\mu \nabla^\mu$ is the covariant d'Alembertian operator [61]. Accordingly, substituting Eq. (7) into Eq. (1) yields:

$$\square A^\nu = J^\nu_{\text{quantum}} + R^\nu{}_\lambda A^\lambda + \nabla^\nu \nabla_\mu A^\mu. \quad (8)$$

In classical spacetime coordinates, the covariant d'Alembertian acting on the potential A^ν yields a wave-like term with curvature corrections:

$$\square A^\nu = g^{\alpha\beta} \partial_\alpha \partial_\beta A^\nu - f^\nu(g^{\alpha\beta}, A^\mu, \partial_\mu A^\lambda, \Gamma^\lambda{}_{\mu\rho}, \partial_\mu \Gamma^\alpha{}_{\beta\lambda}). \quad (9)$$

We define the extreme-field source as the sum of a curvature source and a quantum source—the latter derived from the Euler–Lagrange equations established in Remark 3 and defined in Eq. (54):

$$J^\nu_{\text{extreme}} \equiv J^\nu_{\text{quantum}} + R^\nu{}_\lambda A^\lambda + \nabla^\nu \nabla_\mu A^\mu + f^\nu(g^{\alpha\beta}, A^\mu, \partial_\mu A^\lambda, \Gamma^\lambda{}_{\mu\rho}, \partial_\mu \Gamma^\alpha{}_{\beta\lambda}), \quad (10)$$

and interpret this as an extreme-field operator $\hat{\mathcal{E}}$ acting on the potential A^ν :

$$J^\nu_{\text{extreme}} \equiv \hat{\mathcal{E}} A^\nu. \quad (11)$$

Hence, from Eqs. (9), (10), and (11), we obtain:

$$g^{\alpha\beta}\partial_\alpha\partial_\beta A^\nu = \hat{\mathcal{E}}A^\nu. \quad (12)$$

Assuming a plane-wave form for the vector potential,

$$A^\nu(x) = \epsilon^\nu e^{-i(\omega t - \mathbf{k}\cdot\mathbf{x})}, \quad (13)$$

where ϵ^ν is the polarisation vector, it follows that

$$\partial_\alpha\partial_\beta A^\nu = -k_\alpha k_\beta A^\nu. \quad (14)$$

Substituting Eq. (14) into Eq. (12) yields the modified dispersion relation in terms of the operator-valued metric components, ensuring consistency with the extreme-field operator from Eq. (11):

$$\hat{g}^{\alpha\beta}k_\alpha k_\beta A^\nu = \hat{\mathcal{E}}A^\nu. \quad (15)$$

In weak fields, quantum effects vanish and curvature diminishes, recovering the standard Maxwell dispersion relation [109, 147]:

$$\eta^{\alpha\beta}\partial_\alpha\partial_\beta A^\nu = \left(-\frac{\omega^2}{c^2} + |\mathbf{k}|^2\right)A^\nu = 0, \quad (16)$$

which implies

$$\hat{\mathcal{E}}A^\nu|_{\text{weak field}} = 0. \quad (17)$$

By contrast, extreme fields demand

$$\hat{\mathcal{E}}A^\nu|_{\text{extreme field}} \neq 0. \quad (18)$$

In the extreme regime, the generalised Lorenz condition (Eq. (6)) directly imposes the universality condition:

$$\nabla_\nu \hat{\mathcal{E}}A^\nu|_{\text{extreme field}} = [\nabla_\nu, \hat{\mathcal{E}}]A^\nu + \hat{\mathcal{E}}(\nabla_\nu A^\nu) \neq 0. \quad (19)$$

As will be established in Remark 3, promoting the metric tensor in Eq. (15) to the operator $\hat{g}^{\alpha\beta}$ is required to satisfy the extreme-field non-vanishing condition (18) through quantum corrections, and more generally for a self-consistent formulation of quantum gravity. We therefore obtain:

$$\hat{g}^{\alpha\beta}\partial_\alpha\partial_\beta A^\nu = (\hat{g}^{00}\partial_0\partial_0 + \hat{g}^{ii}\partial_i\partial_i)A^\nu. \quad (20)$$

Explicitly, accounting for the Lorentzian signature, we express the temporal term as

$$\hat{g}^{00}\partial_0\partial_0 = -\hat{g}^{tt}\partial_0\partial_0, \quad (21)$$

and substituting this into Eq. (20) yields:

$$\hat{g}^{\alpha\beta}\partial_\alpha\partial_\beta A^\nu = (-\hat{g}^{tt}\partial_0\partial_0 + \hat{g}^{ii}\partial_i\partial_i) A^\nu = \left(-\frac{\omega^2}{c^2}\hat{g}^{tt} + |\mathbf{k}|^2\hat{g}^{ii} + \delta_{\text{ext}}(k_\mu k^\mu)\right) A^\nu. \quad (22)$$

The extreme-field quantum-gravitational correction in Eq. (22), denoted by $\delta_{\text{ext}}(k_\mu k^\mu)$, can be absorbed into the right-hand side of Eq. (12), yielding the following condition for non-vanishing solutions:

$$\hat{g}^{tt}A^\nu \neq \hat{g}^{ii}A^\nu. \quad (23)$$

The invertibility condition of the quantum operator [50, 55, 135, 157] states:

$$\hat{g}^{\mu\nu}\hat{g}_{\mu\nu} = \hat{g}_{\mu\nu}\hat{g}^{\mu\nu} = \mathbb{I}. \quad (24)$$

By rewriting Eq. (22), we obtain:

$$\left(-\frac{\omega^2}{c^2}\hat{g}^{tt}\mathbb{I} + |\mathbf{k}|^2\hat{g}^{ii}\right) A^\nu = \left(-\frac{\omega^2}{c^2}\hat{g}^{tt}\hat{g}_{ii}\hat{g}^{ii} + |\mathbf{k}|^2\hat{g}^{ii}\right) A^\nu = \left(-\frac{\hat{g}^{tt}\hat{g}_{ii}}{c^2}\omega^2 + |\mathbf{k}|^2\right) (\hat{g}^{ii}A^\nu). \quad (25)$$

From Eq. (24) and the definition in Eq. (30), the extreme-field potential is defined as

$$A_\mathcal{E}^\nu \equiv \hat{g}^{ii}A^\nu, \quad \text{with} \quad \hat{g}^{ii}A^\nu = A^\nu + \hat{\mathcal{O}}_2 A^\nu. \quad (26)$$

Inserting the definition from Eq. (26) into Eq. (25) yields the extreme-field wave equation:

$$\hat{g}^{\alpha\beta}\partial_\alpha\partial_\beta A^\nu = \left(-\frac{\hat{g}^{tt}\hat{g}_{ii}}{c^2}\omega^2 + |\mathbf{k}|^2\right) A_\mathcal{E}^\nu. \quad (27)$$

By interpreting the speed of light as a quantum operator $\hat{\mathcal{C}}$ with inverse $\hat{\mathcal{C}}^{-1}$, we obtain:

$$\left(-\frac{\hat{g}^{tt}\hat{g}_{ii}}{c^2}\omega^2 + |\mathbf{k}|^2\right) A_\mathcal{E}^\nu = \left(-\hat{\mathcal{C}}^{-2}\omega^2 + |\mathbf{k}|^2\right) A_\mathcal{E}^\nu. \quad (28)$$

Thus, non-vanishing solutions require $\hat{\mathcal{C}}^2 \neq \omega^2/|\mathbf{k}|^2$, satisfying Eq. (18). To additionally satisfy the universality condition expressed in Eq. (19), we define:

$$\hat{g}^{tt} \equiv (\mathbb{I} + \hat{\mathcal{O}}_1)^{-1}, \quad (29)$$

$$\hat{g}_{ii} \equiv (\mathbb{I} + \hat{\mathcal{O}}_2)^{-1}. \quad (30)$$

Computing the product

$$(\mathbb{I} + \hat{\mathcal{O}}_1)^{-1}(\mathbb{I} + \hat{\mathcal{O}}_2)^{-1} = [(\mathbb{I} + \hat{\mathcal{O}}_2)(\mathbb{I} + \hat{\mathcal{O}}_1)]^{-1} = (\mathbb{I} + \hat{\mathcal{O}}_1 + \hat{\mathcal{O}}_2 + \hat{\mathcal{O}}_2\hat{\mathcal{O}}_1)^{-1}, \quad (31)$$

we define the extreme-field identity as

$$\hat{\mathcal{O}}_1 + \hat{\mathcal{O}}_2 + \hat{\mathcal{O}}_2\hat{\mathcal{O}}_1 \equiv \frac{\hat{\chi}}{c^2}. \quad (32)$$

This yields the inverse speed-of-light operator

$$\hat{c}^{-2} = \frac{1}{c^2} \left(\mathbb{I} + \frac{\hat{\chi}}{c^2} \right)^{-1}, \quad (33)$$

hence the operator reads:

$$\hat{\mathcal{C}}^2 = c^2 + \hat{\chi}. \quad (34)$$

Remark 1 (Positivity of the extreme-field correction). *Physical consistency requires that the metric operators defined in Eqs. (29) and (30) preserve the Lorentzian signature for all physical states. This imposes the positivity conditions:*

$$\langle \mathbb{I} + \hat{\mathcal{O}}_1 \rangle > 0 \quad \text{and} \quad \langle \mathbb{I} + \hat{\mathcal{O}}_2 \rangle > 0. \quad (35)$$

Assuming $\hat{\mathcal{O}}_1, \hat{\mathcal{O}}_2$ are Hermitian with spectra bounded below by -1 and positive semi-definite covariance, Eq. (35) implies

$$\langle (\mathbb{I} + \hat{\mathcal{O}}_2)(\mathbb{I} + \hat{\mathcal{O}}_1) \rangle = \langle \mathbb{I} + \hat{\mathcal{O}}_2 \rangle \langle \mathbb{I} + \hat{\mathcal{O}}_1 \rangle + \text{Cov}(\hat{\mathcal{O}}_2, \hat{\mathcal{O}}_1) > 0. \quad (36)$$

Using the expansion in Eq. (31) and taking the expectation—with $\langle \mathbb{I} \rangle = 1$ and the definition of $\hat{\chi}/c^2$ from Eq. (32)—Eq. (36) yields the bound

$$\left\langle \frac{\hat{\chi}}{c^2} \right\rangle = \langle \hat{\mathcal{O}}_1 + \hat{\mathcal{O}}_2 + \hat{\mathcal{O}}_2\hat{\mathcal{O}}_1 \rangle > -1. \quad (37)$$

Multiplying by c^2 yields the universal bound $\langle \hat{\chi} \rangle > -c^2$, which holds for every physical state and guarantees a real, positive extreme-field speed of light:

$$\sqrt{\langle \hat{\mathcal{C}}^2 \rangle} = \sqrt{c^2 + \langle \hat{\chi} \rangle} > 0. \quad (38)$$

To continue the proof, we multiply Eq. (28) by $\hat{\mathcal{C}}^2$, obtaining

$$\left(-\hat{\mathcal{C}}^2 \hat{c}^{-2} \omega^2 + |\mathbf{k}|^2 \hat{\mathcal{C}}^2 \right) A_{\mathcal{E}}^{\nu} = \left(-\omega^2 + |\mathbf{k}|^2 (c^2 + \hat{\chi}) \right) A_{\mathcal{E}}^{\nu}. \quad (39)$$

Using the classical dispersion relation $(-\omega^2 + |\mathbf{k}|^2 c^2 = 0)$, the non-trivial solution requires:

$$\hat{\chi} A_{\mathcal{E}}^{\nu} = [\hat{\chi}, \hat{g}^{ii}] A^{\nu} + \hat{g}^{ii} (\hat{\chi} A^{\nu}) \neq 0, \quad (40)$$

which is equivalent to the extreme-field condition established in Eq. (18):

$$\hat{\mathcal{C}}^2 \hat{\mathcal{E}} A^{\nu}|_{\text{extreme field}} = (c^2 + \hat{\chi}) \hat{\mathcal{E}} A^{\nu}|_{\text{extreme field}} \neq 0. \quad (41)$$

Satisfying the universality condition in Eq. (19) necessitates that:

$$\hat{\chi} = Z^{\nu} \nabla_{\nu}, \quad (42)$$

for a contravariant vector field Z^{ν} , distinct from the covariant derivative ∇^{ν} , thereby avoiding a contradiction between Eqs. (9) and (19).

To support Remark 1, and following from the quantum–gravitational non-commutative algebra (Subsection 4.3), the contravariant vector field Z^{ν} is interpreted as a quantum operator (Eq. (245)), taking the form

$$\hat{Z}^{\nu} = i\hbar \alpha \partial^{\nu}. \quad (43)$$

Conditionally, $\hat{\chi} = i\hbar \alpha \partial^{\nu} \nabla_{\nu}$ is Hermitian and possesses real eigenvalues, provided $\hbar \alpha \partial^{\nu} \nabla_{\nu}$ acts as a single-action conjugation at the extreme locations of Hermiticity—thereby rendering it anti-Hermitian on the relevant Hilbert space under the boundary conditions established in Subsection 4.2. The factor i subsequently ensures Hermiticity and thus real expectation values.

Remark 2 (Generalised Hermiticity conditions for $\hat{\chi}$). *If $\hat{\chi}$ deviates from conventional Hermiticity, alternative consistency conditions apply.*

1. **\mathcal{PT} -symmetric generalisation.** *Physical admissibility may persist without strict Hermiticity provided Parity-Time (\mathcal{PT}) symmetry holds [21, 22, 57]. Such operators exhibit a real spectrum and generate pseudo-unitary evolution, preserving essential requirements while accommodating broader quantum-gravitational dynamics.*
2. **Anti-symmetric construction.** *Should \mathcal{PT} symmetry also fail, one may adopt the manifestly Hermitian combination*

$$\hat{\chi}_{\text{anti-sym}} = \frac{i\hbar \alpha}{2} (\partial^{\nu} \nabla_{\nu} - \nabla_{\nu} \partial^{\nu}). \quad (44)$$

This operator is globally Hermitian because i times the commutator of two anti-Hermitian operators yields Hermiticity [145]. It eliminates ordering ambiguity, guarantees real expectation values, and fully satisfies the non-vanishing extreme-field condition in Eq. (41).

In this work, we conditionally adopt $\hat{\chi} = i\hbar\alpha\partial^\nu\nabla_\nu$ as the Hermitian operator at the extreme locations of Hermiticity. The antisymmetric construction—a natural generalisation—is reserved for a forthcoming paper.

Having established $\hat{\chi} = i\hbar\alpha\partial^\nu\nabla_\nu$ as Hermitian, consistency requires that $\hbar\alpha\partial^\nu\nabla_\nu$ be anti-Hermitian, so that the factor i renders it Hermitian. However, the operator $\hat{\chi}$ contains only two derivatives (via $\partial^\nu\nabla_\nu$), suggesting two possibilities. The first is an anticommutative geometry that conditionally endows $\partial^\nu\nabla_\nu$ with anti-Hermiticity. The second interprets α as a unified running coupling that acquires an operator character—acting as a differential generator in a generalised coordinate—so that the combination $\alpha\partial^\nu\nabla_\nu$ introduces an odd number of derivatives, thereby ensuring the odd count of integrations by parts required for anti-Hermiticity. This allows us to define the scalar running-coupling operator, with dimension $[G/c]$ (Eq. (281)) up to a constant factor α_0 , as

$$\alpha \equiv \alpha_u \partial_u, \quad \alpha_u = \alpha_0. \quad (45)$$

It is evident that the second possibility requires more assumptions and complexity than the first; both are investigated in Subsection 4.2 under the Extreme Hermitian Geometry, which ensures the conditional Hermiticity of $\hat{\chi}$.

This construction of the conditional Hermitian operator $\hat{\chi}$ further suggests a natural geometric interpretation in Planck units: the product $\hbar\alpha$ yields the Planck length squared, $l_{\text{Pl}}^2 = G\hbar/c^3$, allowing the leading-order form of $\hat{\chi}$ in Eq. (42) to be written as

$$\hat{\chi} \sim i c^2 l_{\text{Pl}}^2 \partial^\mu \nabla_\mu, \quad (46)$$

where both derivatives carry dimensions of $1/L$ —equivalently, energy E in natural units [6]—and their action on a test function Φ , taken to be a generic non-eigenstate, yields

$$(\partial^\nu \nabla_\nu \Phi)^\sigma = \partial^\nu \partial_\nu \Phi^\sigma + (\partial^\nu \Gamma_{\nu\lambda}^\sigma) \Phi^\lambda + \Gamma_{\nu\lambda}^\sigma \partial^\nu \Phi^\lambda. \quad (47)$$

This generates a non-trivial curvature contribution. By Theorem 3 (Subsection 4.3), this contribution is bounded in the extreme limit; for a conservative structure, this bound must satisfy the universality condition established in Remark 1. For the pure inverse-length-squared contributions, since lengths are by definition positive, the dimensionless quantity $\langle \hat{\chi} \rangle / c^2$ is therefore positive-definite:

$$\frac{\langle \hat{\chi} \rangle}{c^2} \sim \frac{l_{\text{Pl}}^2}{L^2} > 0. \quad (48)$$

Consequently, the expectation value of the squared operator-valued speed of light in the black hole extreme field is

$$\langle \hat{\mathcal{C}}^2 \rangle \simeq c^2 \left(1 + \frac{l_{\text{Pl}}^2}{L^2} \right), \quad (49)$$

which naturally defines two distinct physical regimes:

- **Perturbative regime:** For $L \gg l_{\text{Pl}}$, the quantum correction is small, $\langle \hat{\chi} \rangle / c^2 \ll 1$, and propagation remains effectively classical, corresponding to the perturbative description of the extreme field.
- **Non-perturbative regime:** For $L \sim l_{\text{Pl}}$, the correction becomes of order unity, $\langle \hat{\chi} \rangle / c^2 \sim 1$, signalling the breakdown of perturbation theory and the onset of fully non-perturbative quantum-gravitational effects.

□

Remark 3 (Lagrangian Foundation for the Operator-Valued Speed of Light). *The operator-valued speed of light in Eq. (34) can be motivated from first principles by considering the electromagnetic Lagrangian extended to the extreme-field regime. We begin with the standard Lagrangian density for electrodynamics in curved spacetime, including the corresponding extreme-field corrections:*

$$\mathcal{L} = -\frac{1}{4\mu_0} F^{\mu\nu} F_{\mu\nu} - J^\mu A_\mu + \delta\mathcal{L}_{ext}. \quad (50)$$

The generally covariant Euler–Lagrange equations [39] for the field A_ν , derived from the action $S = \int d^4x \sqrt{-g} \mathcal{L}$, read

$$\nabla_\mu \left(\frac{\partial \mathcal{L}}{\partial(\nabla_\mu A_\nu)} \right) - \frac{\partial \mathcal{L}}{\partial A_\nu} = 0. \quad (51)$$

Solving Eq. (51) using Eq. (2) yields the modified Maxwell equations (1), which incorporate extreme-field quantum-gravitational corrections as follows:

$$\nabla_\mu \left(\frac{\partial \mathcal{L}}{\partial(\nabla_\mu A_\nu)} \right) = \nabla_\mu \left(-\frac{1}{\mu_0} F^{\mu\nu} + \frac{\partial(\delta\mathcal{L}_{ext})}{\partial(\nabla_\mu A_\nu)} \right), \quad (52)$$

$$\frac{\partial \mathcal{L}}{\partial A_\nu} = -J^\nu + \frac{\partial(\delta\mathcal{L}_{ext})}{\partial A_\nu}, \quad (53)$$

$$J_{quantum}^\nu \equiv \mu_0 \left[J^\nu - \frac{\partial(\delta\mathcal{L}_{ext})}{\partial A_\nu} + \nabla_\mu \left(\frac{\partial(\delta\mathcal{L}_{ext})}{\partial(\nabla_\mu A_\nu)} \right) \right]. \quad (54)$$

In extreme-field regimes, the stress–energy tensor can be derived from the variational derivative of the extreme-field effective action with respect to the metric as

$$T_{\mu\nu}^{ext} = -\frac{2}{\sqrt{-g}} \frac{\delta S_{matter}^{eff}}{\delta g^{\mu\nu}}. \quad (55)$$

In conventional physical models—such as perfect fluids, scalar fields, and electromagnetism—the solutions are functionals of the metric tensor. Consequently, their quantisation inherits an explicit functional dependence on the metric operator:

$$T_{\mu\nu}^{ext} \equiv T_{\mu\nu}[\hat{g}_{\mu\nu}]. \quad (56)$$

When considering the Einstein field equations [55],

$$G_{\mu\nu} = \kappa T_{\mu\nu}, \quad (57)$$

the left-hand side (the Einstein tensor) is a functional of the metric $g_{\mu\nu}$, while the right-hand side depends on the metric through the stress–energy tensor, as in Eq. (56). This classical framework becomes inadequate in regimes of extreme gravity, due to the incompatibility between the classical geometry of spacetime and the quantum nature of matter and fields. While the semiclassical approximation—replacing $T_{\mu\nu}$ with its expectation value—suffices in most astrophysical settings, a fully consistent quantum-gravitational treatment requires promoting the metric tensor to a quantum operator [44]. The quantised Einstein field equations then act on a quantum-gravitational state $|\Phi\rangle$:

$$\hat{G}_{\mu\nu}|\Phi\rangle = \kappa \hat{T}_{\mu\nu}|\Phi\rangle. \quad (58)$$

Consistency of Eq. (58) thus requires promoting the metric itself to a quantum operator. This quantisation of geometry naturally implies that both sides of Einstein’s equations are operator-valued.

By induction, this argument justifies the promotion of the inverse metric components to quantum operators in the extreme-field regime, as introduced in Eqs. (29) and (30). Consequently, the mathematical identities (32) acquire a clear physical meaning: they encode the back-reaction of quantised geometry on the propagation of light. The derivation of the inverse speed-of-light operator in Eq. (33) and the final operator expression $\hat{C}^2 = c^2 + \hat{\chi}$ in Eq. (34) then follow directly, establishing that the operator-valued speed of light emerges naturally from the quantisation of the metric in extreme gravitational regimes.

Motivation 2. *Quantising the speed of light yields a quantised noncommutative framework that simultaneously regularises the physical singularity problem [124] and resolves the black hole information paradox [76]. This is achieved via an interior version of Hawking radiation, where outward-propagating negative-energy modes—the anti-eigen flux—provide the quantum-geometric mechanism for unitarity and information preservation [70].*

Proposition 3 (Non-singular extreme-field quantum-gravitational regularisation). *The quantisation of the speed of light in extreme-field regimes is expressed by the operator*

$$\hat{C}^2 = c^2 + \hat{\chi}, \quad (59)$$

where $\hat{\chi}$ is defined in Eq. (42). Promoting the radial coordinate to an operator \hat{r} , as permitted by Principle 1, consequently induces a fundamental noncommutativity [116],

$$[\hat{r}, \hat{\chi}] \neq 0, \quad (60)$$

which is required to be finite over the entire physical domain, including at the classical singularity $r = 0$. This noncommutativity implies a quantum-gravitational uncertainty relation for r and χ ,

$$\Delta r \Delta \chi \geq \frac{1}{2} |\langle [\hat{r}, \hat{\chi}] \rangle|, \quad (61)$$

thereby providing a quantum-geometric regularisation (see Subsection 4.3) that ensures the expectation value of the operator-valued metric tensor remains finite:

$$\lim_{r \rightarrow 0} \langle \hat{g}_{\mu\nu} \rangle < \infty, \quad \forall \mu, \nu. \quad (62)$$

Principle 1 (Extreme-field quantisation of spacetime geometry). *In extreme-field regimes, the metric tensor $\hat{g}_{\mu\nu}$ is promoted to a quantum operator, making the line element $d\hat{s}^2 = \hat{g}_{\mu\nu} dx^\mu dx^\nu$ operator-valued. Coordinates remain classical parameters that label spacetime points, while spatial coordinates may be treated as operators and time strictly remains a classical external parameter [7, 82]. Thus, quantum nature resides entirely in geometry, not coordinate labels. Quantising spacetime thereby means that distances, intervals, and causal structure fluctuate quantum mechanically, while the underlying coordinate manifold remains classical.*

Proposition 4 (Metric Quantisation). *Replacing c by \hat{C} in the spacetime line element promotes the metric to an operator-valued quantity, with corrections interpretable as renormalisation terms.*

Proposition 5 (Quantum-Corrected Metric). *For quantum-corrected metrics consistent with Einstein field equations in presence of non-vanishing quantum energy-momentum tensor:*

$$\hat{G}_{\mu\nu} = \kappa \hat{T}_{\mu\nu}^{(\text{quantum})}, \quad \hat{T}_{\mu\nu}^{(\text{quantum})} \neq 0, \quad (63)$$

the general static, spherically symmetric quantum-corrected Schwarzschild metric ansatz is [28]:

$$ds^2 = -c^2 \left(1 - \frac{2GM}{rc^2} + f(r) \right) dt^2 + \left(1 - \frac{2GM}{rc^2} + g(r) \right)^{-1} dr^2 + r^2 d\Omega^2. \quad (64)$$

Corollary 1. *The functions $f(r)$ and $g(r)$ represent independent quantum-gravitational corrections; the choice $f(r) = g(r)$ restricts the equation of state, while generically $f(r) \neq g(r)$.*

Proposition 6 (Quantised Schwarzschild Metric). *Quantising the classical Schwarzschild metric by mapping $c \mapsto \hat{C}$ yields an operator-valued line element; its expectation value reduces to the quantum-corrected form in Eq. (64).*

Proof. Starting from classical Schwarzschild metric:

$$ds^2 = -c^2 \left(1 - \frac{2GM}{rc^2} \right) dt^2 + \left(1 - \frac{2GM}{rc^2} \right)^{-1} dr^2 + r^2 d\Omega^2, \quad (65)$$

promote the classical quantities to their operator counterparts as $1 \mapsto \mathbb{I}$, $c \mapsto \hat{C}$ and $r \mapsto \hat{r}$:

$$d\hat{s}^2 = -\hat{C}^2 \left(\mathbb{I} - 2GM\hat{C}^{-2}\hat{r}^{-1} \right) dt^2 + \left(\mathbb{I} - 2GM\hat{C}^{-2}\hat{r}^{-1} \right)^{-1} dr^2 + \hat{r}^2 d\Omega^2. \quad (66)$$

Substituting Eq. (59) and handling noncommutativity:

$$\begin{aligned} d\hat{s}^2 &= -(c^2 + \hat{\chi} - 2GM\hat{r}^{-1}) dt^2 + \left(\mathbb{I} - 2GM\hat{r}^{-1}\hat{C}^{-2} - 2GM[\hat{r}^{-1}, \hat{C}^{-2}] \right)^{-1} dr^2 + r^2 d\Omega^2 \\ &= -c^2 \left(\mathbb{I} - \frac{2GM}{c^2}\hat{r}^{-1} + \frac{\hat{\chi}}{c^2} \right) dt^2 + \left(\mathbb{I} - \frac{2GM}{c^2}\hat{r}^{-1} \left(\mathbb{I} + \frac{\hat{\chi}}{c^2} \right)^{-1} + \Lambda \right)^{-1} dr^2 + r^2 d\Omega^2, \end{aligned} \quad (67)$$

where the non-trivial commutator:

$$\Lambda \equiv -\frac{2GM}{c^2}\hat{r}^{-1}\hat{r}[\hat{r}^{-1}, c^2\hat{C}^{-2}] = -\frac{2GM}{c^2}\hat{r}^{-1} + \mathcal{O}(\hat{r}, \hat{r}^{-1}, \hat{\chi}). \quad (68)$$

Rewriting Eq. (67):

$$d\hat{s}^2 = -c^2 \left(\mathbb{I} - \frac{2GM}{c^2} \hat{r}^{-1} + \hat{Q}_t \right) dt^2 + \left(\mathbb{I} - \frac{2GM}{c^2} \hat{r}^{-1} + \hat{Q}_r \right)^{-1} dr^2 + r^2 d\Omega^2, \quad (69)$$

with quantum corrections:

$$\hat{Q}_t \equiv \frac{\hat{\chi}}{c^2}, \quad \hat{Q}_r \equiv -\frac{2GM}{c^2} \hat{r}^{-1} \left(\mathbb{I} + \frac{\hat{\chi}}{c^2} \right)^{-1} + \mathcal{O}(\hat{r}, \hat{r}^{-1}, \hat{\chi}). \quad (70)$$

Expectation values yield Eq. (64) with:

$$f(r) = \langle \hat{Q}_t \rangle, \quad g(r) = \langle \hat{Q}_r \rangle. \quad (71)$$

Consistency of Eqs. (29) and (30) from Proposition 2, when employed in Eq. (67) for purely radial motion (negligible angular components, Section 3, Eq. (98)), demands the following definitions to verify extreme-field identity Eq. (32):

$$\hat{\mathcal{O}}_1 + \hat{\mathcal{O}}_2 + \hat{\mathcal{O}}_2 \hat{\mathcal{O}}_1 = \frac{\hat{\chi}}{c^2} : \quad (72)$$

$$\mathbb{I} \equiv 1, \quad B \equiv \frac{2GM}{c^2} \hat{r}^{-1}, \quad x \equiv \frac{\hat{\chi}}{c^2}, \quad \hat{\mathcal{O}}_1 = x - B, \quad \hat{\mathcal{O}}_2 = -B(1+x)^{-1} + \Lambda. \quad (73)$$

Inserting these definitions yields:

$$\begin{aligned} \hat{\mathcal{O}}_1 + \hat{\mathcal{O}}_2 + \hat{\mathcal{O}}_2 \hat{\mathcal{O}}_1 &= (x - B) + [-B(1+x)^{-1} + \Lambda] + [-B(1+x)^{-1} + \Lambda](x - B) \\ &= x - B - B(1+x)^{-1} + \Lambda - B(1+x)^{-1}(x - B) + \Lambda(x - B). \end{aligned} \quad (74)$$

For $x \ll 1$, the series expansion of $(1+x)^{-1}$ fails; the exact form applies:

$$\hat{\mathcal{O}}_1 + \hat{\mathcal{O}}_2 + \hat{\mathcal{O}}_2 \hat{\mathcal{O}}_1 = x + [-B - B(1+x)^{-1} - B(1+x)^{-1}(x - B) + \Lambda(1+x - B)]. \quad (75)$$

For $x \ll 1$, the perturbative expansion is valid:

$$(1+x)^{-1} = 1 - x + x^2 - x^3 + x^4 - x^5 + x^6 + \mathcal{O}(x^7). \quad (76)$$

Substituting this expansion into the definition of $\hat{\mathcal{O}}_2$ yields

$$\begin{aligned} \hat{\mathcal{O}}_2 &= -B(1 - x + x^2 - x^3 + x^4 - x^5 + x^6) + \Lambda + \mathcal{O}(x^7) \\ &= -B + Bx - Bx^2 + Bx^3 - Bx^4 + Bx^5 - Bx^6 + \Lambda + \mathcal{O}(x^7). \end{aligned} \quad (77)$$

Computing the product to sixth order:

$$\begin{aligned}\hat{\mathcal{O}}_2\hat{\mathcal{O}}_1 &= (-B + Bx - Bx^2 + Bx^3 - Bx^4 + Bx^5 - Bx^6 + \Lambda)(x - B) \\ &= B^2 - Bx + Bx^2 - Bx^3 + Bx^4 - Bx^5 + Bx^6 - BxB + Bx^2B - Bx^3B \\ &\quad + Bx^4B - Bx^5B + Bx^6B + \Lambda x - \Lambda B + \mathcal{O}(x^7).\end{aligned}\tag{78}$$

Combining all contributions yields the perturbative expansion:

$$\hat{\mathcal{O}}_1 + \hat{\mathcal{O}}_2 + \hat{\mathcal{O}}_2\hat{\mathcal{O}}_1 = x + \text{corrections}.\tag{79}$$

Both the non-perturbative result in Eq. (75) and the perturbative expression in Eq. (79) satisfy Eq. (32) up to residual corrections. To preserve full consistency, counter-corrections naturally arise within the metric components, absorbing higher-order effects in analogy with renormalisation in quantum field theory. As a result, the metric tensor emerges as a quantum-renormalised operator that is intrinsically free of singularities. \square

Remark 4. *An initial resolution of the physical singularity will be discussed in Subsection 4.3, in particular within Theorem 3, while a complete analysis of the metric tensor quantisation and singularity resolution will be presented in a forthcoming paper. In this work, we focus on demonstrating how the non-perturbative quantisation framework establishes the initial conditions for the emergence of the universe, and how the perturbative regime provides a mechanism for resolving the information paradox.*

Definition 2. *A quantum probabilistic metric is a metric that deviates from classical behaviour through extreme operator-valued corrections, transforming the deterministic classical metric into a quantum-mechanical probabilistic object. Superposition between space and time components is allowed, conferring upon the black hole an effective double-slit behaviour.*

Theorem 1 (Preservation of black-hole interior geometry). *Extreme-field corrections, established in Proposition 6, confer their quantum probabilistic nature to the metric, thereby allowing a non-vanishing probability that the interior black-hole geometry is preserved. This keeps the temporal and radial metric components stable and prevents the classical collapse or reversal of the interior causal structure.*

Proof. For quantum-corrected metric operator, Eq. (69), with components:

$$\hat{g}_{00} = -\left(\hat{\mathbb{I}} - \frac{2GM}{c^2}\hat{r}^{-1} + \hat{Q}_t\right), \quad \hat{g}_{rr} = \left(\hat{\mathbb{I}} - \frac{2GM}{c^2}\hat{r}^{-1} + \hat{Q}_r\right)^{-1}, \tag{80}$$

there exist non-trivial quantum corrections to the Lorentzian signature, i.e.,

$$\langle \hat{g}_{00} \rangle < 0 \quad (\text{temporal character preserved}), \quad (81)$$

$$\langle \hat{g}_{rr} \rangle > 0 \quad (\text{spatial character preserved}). \quad (82)$$

provided \hat{Q}_t and \hat{Q}_r evolve positively, as discussed in Remark 1. \square

Proposition 7 (Quantum Escapable Horizons). *Quantum probabilistic metrics admit quantum escapable horizons, where extreme acausal excitations allow information to escape as anti-eigen flux.*

Proof. The semiclassical Hawking radiation tunnelling probability is determined by the metric components [8, 90, 121, 154]:

$$\Gamma \sim \exp \left[f(\langle \hat{g}_{00} \rangle, \langle \hat{g}_{rr} \rangle) \right]. \quad (83)$$

By Theorem 1, the quantum-probabilistic metric preserves the interior causal structure, enabling extreme acausal excitations of the *anti-eigen flux* to tunnel through a quantum-escapable horizon. This mechanism yields a consistent quantum-gravitational channel for information escape with bounded probability, constituting a *necessary* condition for unitary evolution. A full derivation is given in Section 3, see Subsections 3.2 and 3.5. \square

Motivation 3. *A primordial black hole whose acausal field saturates the lower boundary of non-perturbative quantisation naturally establishes the cosmological initial conditions. This process is driven by an intrinsic asymmetry: the creation rate of eigen-Hawking radiation substantially exceeds that of standard Hawking emission. The resulting negentropic flux possesses a contracting area that approaches the Planck scale—forming what we term the primordial Planckian black hole.*

Theorem 2 (Primordial Planck Black Hole as a Jerk-Driven Cosmological Seed). *The primordial cosmological state is a primordial Planck black hole occupying an extreme negentropic phase ($\Delta S \ll 0$). This state undergoes an extremisation phase transition mediated by a quantum-gravitational jerk (see Subsection 4.5). This transition enforces the classical limit $\mathbf{S} \gg \hbar$, thereby marking the onset of classical behaviour [60], the decoupling of gravity, and the beginning of inflation—the birth of the Universe. As a direct consequence, the gravitational and quantum couplings separate, producing an emergent universe with asymptotically free gravitational dynamics.*

Proof. A complete derivation is presented in subsections 3.3, 3.4, and 3.7 of Section 3, as well as in Section 4. \square

2 Introduction

Standard theoretical frameworks describe early-universe dynamics as governed by quantum vacuum fluctuations. When these fluctuations become nonlinear—modelled in perturbative extreme field theory as corrections to the classical metric arising from higher-order quantum nonlinear perturbations, as established in Proposition 6—they can seed primordial black hole (PBH) formation [10, 32, 72, 92, 112, 130, 150]. In conventional models, such PBHs subsequently evaporate via Hawking radiation, a process that adheres to the second law of thermodynamics ($\Delta S > 0$) [38, 43, 70].

Extreme field theory departs from the standard picture by positing a three-stage ‘Pre-Being’ epoch, characterised by a *reversed* causal flow.

Event 1 involves the initial Pre-Being Fluctuations (PBFs)—causal excitations obeying special relativity [54](see Subsection 3.1).

Event 2 is the subsequent formation of primordial black holes (PBHs), whose fields exhibit *acausal, perturbative* excitations, described by a perturbed speed of light (see Subsection 3.2):

$$c_p = \sqrt{c^2 + v}. \quad (84)$$

Conventionally, the sequence of the Pre-Being events is assumed to be

$$\text{Event 1 (PBFs)} \xrightarrow{\text{led to}} \text{Event 2 (PBHs)}. \quad (85)$$

However, this ordering gives rise to a *Pre-Being Paradox*: PBFs require spacetime, yet spacetime had not yet emerged. Our resolution—that acausal excitations inside PBHs arise due to the quantisation of the speed of light—addresses this paradox through non-causal ordering. Specifically, if a causal event generates a superluminal one, the latter may occur first. Consequently, the cosmological sequence is reversed: primordial black holes emerge first, notably as strongly entangled pairs [104], thereby establishing the spacetime arena necessary for the precursor propagation of pre-being fluctuations [153].

$$\text{Event 1 (PBFs)} \xleftarrow{\text{preceded}} \text{Event 2 (PBHs)}. \quad (86)$$

Event 3 is a maximally entangled primordial black hole whose acausal field configuration saturates the lower non-perturbative quantisation boundary (see Subsections 3.3 and 3.7), departing from the perturbative physics of Event 2 and governed by the non-perturbative speed-of-light operator:

$$\hat{C} = \sqrt{c^2 + \hat{\chi}}. \quad (87)$$

This intrinsically non-perturbative primordial black hole—embedded within the nonlinear (non-Gaussian) primordial fluctuation model—represents the extreme limit of such nonlinearity. Its core must substantially deviate from translational symmetry, thereby intrinsically exhibiting dynamics that generate a *negentropic flux* accompanied by area contraction, violating the Hawking area theorem [1,75]. These dynamics culminate at the Planck scale in the formation of the *Primordial Planck Black Hole* (PPBH) [30], which represents the upper limit of non-perturbative quantisation. At this extremisation threshold, the Planckian black hole undergoes an irreversible jerk (recoil), constituting the proposed seed of the universe (see Section 4).

The remainder of this paper develops the theoretical machinery underlying this three-stage Pre-Being scenario. Section 3 develops the theoretical framework for deriving Hawking radiation tunnelling probabilities, distinguishing three regimes: classical pre-being fluctuations (Event 1, Subsection 3.1), the perturbative framework for primordial black holes that addresses the information paradox via extreme-field symmetry (Event 2, Subsection 3.2), and the non-perturbative framework that departs from this symmetry (Event 3, Subsection 3.3), alongside a general interpretation of time in Subsection 3.4. Within this structure, we interpret the anti-eigen flux as a negative-energy solution (Subsection 3.5), introduce an extreme-field eigen-intrinsic asymmetry parameter to account for matter–antimatter asymmetry—favouring the tunnelling of the anti-eigen flux as antimatter while confining eigen-Hawking radiation as matter (Subsection 3.6)—and examine the negentropic evolution of non-perturbative primordial black holes toward the Planckian extremisation threshold (Subsection 3.7).

Section 4 focuses entirely on the non-perturbative framework describing the origin of the universe. It begins by examining the cosmological initial asymmetries (Subsection 4.1), the conditional Hermiticity of the quantum-gravitational conjugate and its associated boundary conditions (Subsection 4.2), and the extreme-field algebra together with its implications for singularity resolution (Subsection 4.3). After presenting the extreme-field equation (Subsection 4.4), we examine the quantum-gravitational jerking of the primordial Planck black hole as the natural mechanism for the birth of the universe, where quantum-gravitational effects decouple (Subsection 4.5). We then present numerical values for the constants of creation (Subsection 4.6) and propose an asymptotically free gravitational universe, in which these constants are enhanced inversely by the attenuated unified coupling—offering a potential resolution to the cosmological constant problem (Subsection 4.7).

3 Theory

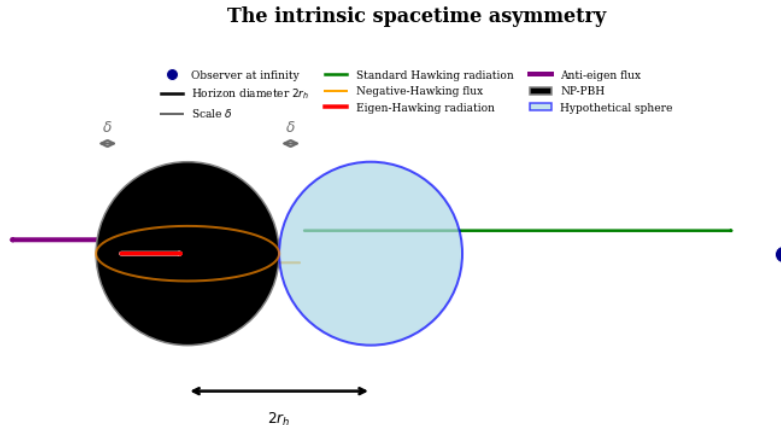


Figure 1: Intrinsic spacetime asymmetry. Two spheres of equal volume share a common tangent orthogonal to their aligned horizons: a black sphere representing the non-perturbative primordial black hole (NP-PBH), governed by extreme field theory; and a blue sphere representing the hypothetical sphere of the pre-being epoch, governed by general relativity. At a distance δ outside the right horizon (on the blue sphere), standard Hawking radiation (green) escapes outward, while negative-Hawking flux (orange) fall toward the NP-PBH centre. Translating this configuration by the horizon diameter $2r_h$ places the same distance δ inside the left horizon of the NP-PBH. There, eigen-Hawking radiation (red) flows inward toward the centre, while anti-eigen flux (thick purple) tunnels outward to the left. The thicker purple beam symbolises the enhanced tunnelling probability of anti-eigen flux in the non-perturbative regime compared to classical Hawking radiation, revealing the fundamental asymmetry.

Extreme field theory posits that, within the extreme gravitational fields of a black hole, the speed of light emerges as a quantum-gravitational operator whose expectation value can exceed its classical value—thereby challenging the second postulate of special relativity regarding the universal constancy of the speed of light [54]. Excitations within these fields become *acausal*, and extreme-field corrections emerge from Proposition 6. This structure gives rise to a quantum-gravitational, probabilistic field that, with non-vanishing

probability, preserves the causal structure of the metric tensor: the time coordinate remains timelike while the radial coordinate remains spacelike, thus preventing the conventional sign flip inside the horizon [62, 110].

This causal invariance enables eigen-Hawking radiation to seed the black hole core while the anti-eigen flux tunnels outward. To a distant observer, the flux appears as backward-in-time particles due to its inwardly increasing radial flow, whereas forward-in-time Hawking radiation flows toward the singularity in a radially decreasing flux.

A necessary condition for unitarity preservation is satisfied if the tunnelling probabilities of the outer standard Hawking radiation (SHR) and the inner anti-eigen flux are symmetric. Within this framework, information propagates quantum-gravitationally as an acausal anti-eigen flux through a quantum-escapable horizon, which need not coincide with the classical event horizon. This model applies to both contemporary and primordial black holes under the perturbative quantisation of the speed of light.

The formation of a Planck-scale black hole (a Primordial Planck Black Hole) necessitates that the core field of its progenitor saturate the lower boundary of the non-perturbative quantisation of the speed of light. This saturation enforces an intrinsic spacetime asymmetry, profoundly differentiating the creation rates of eigen-Hawking radiation (EHR) and standard Hawking radiation (SHR), with EHR overwhelmingly exceeding SHR, as formalised in Theorem 2. This *intrinsic spacetime asymmetry* is rooted in the unique features of non-perturbative quantisation. The process drives the primordial black hole towards a highly ordered, negentropic Planckian state, triggering cosmic genesis from an inherently asymmetric origin [9, 45, 64, 71].

The intrinsic spacetime asymmetry is further elucidated by the thought experiment depicted in Fig. 1. Consider a hypothetical sphere, equal in volume to the non-perturbative primordial black hole (NP-PBH), positioned just outside its horizon. In this region, spacetime is a continuous manifold well described by general relativity [55, 56], where the sphere and the NP-PBH are connected at two points by a common tangent orthogonal to their aligned horizons. At an infinitesimal distance δ outside the horizon, Pre-Being fluctuations allow Hawking radiation to escape outward toward the sphere's centre, while negative-energy particles fall into the NP-PBH's core. Translating this sphere to coincide with the NP-PBH reveals a different regime: inside, spacetime is quantised and governed by extreme-field physics. At the same infinitesimal distance δ *inside* the horizon, extreme fluctuations enable positive-energy eigen-Hawking radiation to flow—in the same direc-

tional sense as standard Hawking radiation—toward the singularity, while negative-energy particles move away from it. Here, the negative-energy anti-eigen flux represents tunnelled particles, while the positive-energy EHR seeds the NP–PBH. Spacetime asymmetry emerges if extreme-field physics permits a significantly higher creation rate for inner EHR compared to outer SHR. This disparity produces a negentropic singular state, where all primordial microstates diminish via the outward anti-eigen flux, leaving the NP–PBH in a state of extreme negentropy. Such order surpasses even that of zero entropy, representing a highly destabilised Planckian threshold that necessitates an extreme phase transition to macroscopic scales—the hypothesised mechanism of cosmic genesis.

To elucidate the intrinsic spacetime asymmetry, it is instructive to examine the classical mechanism underlying Hawking radiation. In the vicinity of a black-hole event horizon, particle–antiparticle pairs spontaneously materialise from the vacuum. Under the intense gravitational field near the horizon, one particle may be captured while its partner escapes to infinity [151]. Employing the semiclassical WKB approximation [31, 95, 121, 161], the tunnelling probability for the escaping particle is given by

$$\Gamma \sim \exp\left(-\frac{2\Im S}{\hbar}\right). \quad (88)$$

The expression highlights the tunnelling rate’s exponential sensitivity to the imaginary part of the particle’s classical action. Within the WKB framework, this imaginary component corresponds to the imaginary contribution arising from the radial integral of the particle’s momentum, p_r , as it propagates in the curved black hole spacetime [154]:

$$\Im S = \Im \int_{r_{\text{in}}}^{r_{\text{out}}} p_r dr. \quad (89)$$

Here, the integration limits span from a point just inside the horizon, $r_{\text{in}} = r_s - \epsilon$ (where $r_{\text{in}} < r_s$), to a point just outside, $r_{\text{out}} = r_s + \epsilon$ (where $r_{\text{out}} > r_s$), with $r_s = 2GM/c^2$ denoting the Schwarzschild radius.

In general relativity [55, 56], the motion of a particle in curved spacetime is governed by the geodesic equation. For a spherically symmetric spacetime, the line element may be expressed as

$$ds^2 = g_{tt}c^2 dt^2 + g_{rr}dr^2 + r^2 d\Omega^2, \quad (90)$$

where g_{tt} and g_{rr} depend only on the radial coordinate r , and $d\Omega^2$ denotes the metric on the unit 2-sphere. In the Schwarzschild solution [146], which represents the unique static, spherically symmetric vacuum solution to Einstein's field equations [55, 56], both the energy-momentum tensor and the Einstein tensor vanish identically, $T_{\mu\nu} = 0$ and $G_{\mu\nu} = 0$. The non-trivial metric components are therefore

$$g_{tt} = - \left(1 - \frac{2GM}{rc^2} \right), \quad (91)$$

$$g_{rr} = \left(1 - \frac{2GM}{rc^2} \right)^{-1}. \quad (92)$$

The four-momentum of a particle, fundamental in both special and general relativity, is defined by [110]

$$p_\mu = mu_\mu = m \frac{dx_\mu}{d\tau}, \quad (93)$$

where m is the rest mass, u_μ the covariant four-velocity, and τ the proper time. The temporal component, associated with the conserved energy in a stationary spacetime, is given by

$$p_t = -mu_t. \quad (94)$$

Since

$$u^\mu = \frac{dx^\mu}{d\tau}, \quad (95)$$

and the metric is time-independent, the conserved energy E is

$$E = -p_t c = -g_{tt} m c^2 \frac{dt}{d\tau}. \quad (96)$$

Similarly, the radial momentum is

$$p_r = mu_r = m g_{rr} \frac{dr}{d\tau}. \quad (97)$$

Assuming the angular components of the four-velocity, $d\phi/d\tau$ and $d\theta/d\tau$, vanish or are negligible, the motion can be treated as purely radial. Then the normalisation condition for a timelike worldline, $u^\mu u_\mu = -c^2$ [110], gives

$$g_{tt} c^2 \left(\frac{dt}{d\tau} \right)^2 + g_{rr} \left(\frac{dr}{d\tau} \right)^2 = -c^2. \quad (98)$$

Using Eq. (96), we can express

$$\frac{dt}{d\tau} = \frac{E}{-g_{tt}mc^2}. \quad (99)$$

Substituting (99) into (98) yields

$$g_{tt}c^2 \left(\frac{E}{-g_{tt}mc^2} \right)^2 + g_{rr} \left(\frac{dr}{d\tau} \right)^2 = -c^2, \quad (100)$$

which simplifies to

$$\frac{-1}{g_{tt}} \left(\frac{E}{mc} \right)^2 - g_{rr} \left(\frac{dr}{d\tau} \right)^2 = c^2. \quad (101)$$

Multiplying Eq. (101) by m^2g_{rr} and using the definition of p_r from Eq. (97), we obtain

$$p_r^2 = g_{rr} \left(\frac{E^2}{-g_{tt}c^2} - m^2c^2 \right). \quad (102)$$

We now impose two physically motivated conditions. First, we consider the high-energy regime,

$$E \gg mc^2, \quad (103)$$

and second, we assume a spherically symmetric, non-rotating black hole described by the Schwarzschild solution. In this case,

$$g^{rr} = \frac{1}{g_{rr}} = -g_{tt}. \quad (104)$$

Applying (103) and (104) to Eq. (102) and taking the positive root, we obtain

$$p_r \approx \frac{E}{-g_{tt}c}. \quad (105)$$

Substituting (105) into Eq. (89) and inserting the explicit form of g_{tt} from (91), we obtain

$$\Im S = \frac{E}{c} \Im \int_{r_{\text{in}}}^{r_{\text{out}}} \frac{1}{\left(1 - \frac{2GM}{rc^2}\right)} dr. \quad (106)$$

The integrand exhibits a simple pole at $r = r_s = \frac{2GM}{c^2}$. To evaluate the integral, the contour is deformed into the complex r -plane according to the standard tunnelling prescription (equivalent to the Feynman $i\epsilon$ rule), whereby

the integration path is deformed only minimally to avoid the pole rather than encircling it completely. Consequently, a semicircular $i\epsilon$ deformation of the real-axis pole yields a half-residue contribution, giving π instead of 2π . Expanding near the pole gives

$$\frac{1}{1 - \frac{2GM}{rc^2}} = \frac{r}{r - r_s} \approx \lim_{r \rightarrow r_s} \frac{r_s}{r - r_s}, \quad (107)$$

and thus the residue at the pole is $r_s = 2GM/c^2$. Hence,

$$\Im \int_{r_{\text{in}}}^{r_{\text{out}}} \frac{dr}{1 - \frac{2GM}{rc^2}} = \pi \frac{2GM}{c^2}. \quad (108)$$

Substituting (108) into (106), we obtain

$$\Im S = \frac{E}{c} \cdot \pi r_s = \frac{2\pi GME}{c^3}. \quad (109)$$

Each member of the Hawking pair contributes a semicircular (π) imaginary part upon crossing the horizon pole; taken together, the pair yields an overall factor 2π .

$$\Im S_{\text{total}} = \frac{4\pi GME}{c^3}. \quad (110)$$

Inserting Eq.(110) into the WKB formula (88) yields the tunnelling probability for the standard Hawking radiation (SHR)

$$\Gamma_{\text{SHR}} \sim \exp\left(-\frac{8\pi GME}{\hbar c^3}\right). \quad (111)$$

Eq. (111) represents the standard semi-classical tunnelling probability in the fixed-background approximation, i.e. when the black-hole mass is treated as constant during the emission process.

A complete treatment must strictly enforce energy conservation: emission of a quantum of energy E reduces the black-hole mass from M to $M - E/c^2$. This back-reaction modifies the horizon radius entering the radial momentum p_r , consequently altering the imaginary part of the action. Evaluating the contour integral with the corrected horizon, $r_s^- = 2G(M - E/c^2)/c^2$, yields the corrected tunnelling probability for standard Hawking radiation [121]:

$$\Gamma_{\text{SHR}}^{\text{BR}} \sim \exp\left[-\frac{8\pi GE}{\hbar c^3} \left(M - \frac{E}{2c^2}\right)\right] = \exp\left[\frac{-8\pi GE\Delta^-}{\hbar c^3}\right]. \quad (112)$$

In Eq.(112) we introduced the quantity $\Delta^- \equiv M - E/(2c^2)$, where the term $E/(2c^2)$ is the leading back-reaction correction. In the limit $E \ll Mc^2$ this term is negligible, and Eq.(111) is recovered; when E becomes comparable to Mc^2 the correction makes the spectrum slightly non-thermal, reflecting the fact that the black hole shrinks as it radiates.

We now derive the analogous tunnelling probability for the acausal excitation emerging from the black-hole interior, which we term the anti-eigen flux. A physically meaningful tunnelling description for this flux from within the black hole is possible only if the excitation preserves the causal structure of the metric tensor: i.e., the temporal component remains negative (the time coordinate stays timelike) and the radial component remains positive (the radial coordinate stays spacelike), as established in Theorem 1. Assuming the Schwarzschild form of the metric remains valid, with quantum-corrected components still satisfying Eq. (105), it suffices to examine the conservative structure of the temporal metric component, denoted j_{tt} .

In the classical Schwarzschild geometry, the temporal metric component outside the event horizon is negative (timelike) and can be written as

$$g_{tt} = -\left(1 - \frac{f(r)}{c^2}\right), \quad (113)$$

where $f(r)$ is a function of the radial coordinate. The horizon r_h is defined by the condition $f(r_h) = c^2$. Consequently,

$$\frac{f(r)}{c^2} \begin{cases} = 1, & r = r_h, \\ < 1, & r > r_h, \\ > 1, & r < r_h. \end{cases} \quad (114)$$

Eq. (114) shows that inside the horizon ($r < r_h$) the quantity $1 - f(r)/c^2$ becomes negative, so g_{tt} turns positive. In the standard interpretation, this sign flip renders the coordinate t spacelike and r timelike, forcing all future-directed causal trajectories toward decreasing r (i.e., toward the singularity). In a purely classical picture, the notion of a particle tunnelling outward from inside the horizon is physically meaningless, as no timelike or null worldline connects the interior to the exterior. Consequently, any semiclassical tunnelling analysis must either modify the causal structure near the horizon or invoke a quantum-gravitational resolution of the coordinate singularity.

Extreme field theory alters the causal structure of black hole interiors through its central postulate: the black hole interior can support acausal

excitations characterised by a quantised speed of light. This speed emerges as a quantum-gravitational operator, with both a perturbative representation (Eq. (84)) and a unique non-perturbative form (Eq. (87)). We denote the non-perturbative operator as $\hat{\chi}$, which encodes the full quantum-gravitational corrections and reduces to the perturbative form v in the appropriate limit.

This framework requires the invertibility of the quantised speed-of-light operator as a key property for a well-defined quantum-gravitational operator algebra and for ensuring the mathematical consistency of the underlying field theory. This invertibility condition reads

$$\hat{\mathcal{C}}\hat{\mathcal{C}}^{-1} = \hat{\mathcal{C}}^{-1}\hat{\mathcal{C}} = \mathbb{I}. \quad (115)$$

Without loss of generality, the quantised metric $\hat{j}_{\mu\nu}$ within the extreme field theory can be formulated as a quantum operator acting on the extreme-field Φ :

$$d\hat{s}^2\Phi = [dt^2\hat{\mathcal{C}}^2 \hat{j}_{tt} + d\mathbf{x}^2 \hat{j}_{xx}]\Phi, \quad (116)$$

where the speed of light is promoted into its quantum operator version and its inverse. By utilising Eqs. (87) and (115), we can rewrite Eq. (116) as:

$$d\hat{s}^2\Phi = [-dt^2\hat{\mathcal{C}}^2(1 - \zeta(r)\hat{\mathcal{C}}^{-2}) + d\mathbf{x}^2\hat{j}_{xx}]\Phi \quad (117)$$

$$= [-dt^2(\hat{\mathcal{C}}^2 - \zeta(r)\hat{\mathcal{C}}^2\hat{\mathcal{C}}^{-2}) + d\mathbf{x}^2\hat{j}_{xx}]\Phi \quad (118)$$

$$= [-dt^2(\hat{\mathcal{C}}^2 - \zeta(r)\mathbb{I}) + d\mathbf{x}^2\hat{j}_{xx}]\Phi \quad (119)$$

$$= [-dt^2(c^2 + \hat{\chi} - \zeta(r)\mathbb{I}) + d\mathbf{x}^2\hat{j}_{xx}]\Phi \quad (120)$$

Therefore, in the semiclassical picture, the normalised temporal component of the quantised metric is obtained by taking the expectation value:

$$\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle = -(c^2 + \chi - \zeta(r)), \quad (121)$$

where the expectation value for the quantum operator $\hat{\chi}$ is:

$$\langle \hat{\chi} \rangle = \chi. \quad (122)$$

Consistency with Remark 1 and Theorem 1—which establish a non-vanishing probability for a conservative structure of the quantised metric when corrected with extreme-field conservative corrections, strictly in the non-perturbative

regime—implies that the expectation value of the normalised temporal component of the quantised metric (Eq. (121)) must be negative in the black-hole interior:

$$\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle < 0. \quad (123)$$

If Eq. (123) is satisfied, thereby realising a conservative structure, then Eq. (121) yields the following extreme-field conservative inequality:

$$\chi > \zeta(r) - c^2. \quad (124)$$

Comparing the normalised temporal components of the classical metric—Eq. (113) multiplied by c^2 —and its quantised counterpart in Eq. (121), we first write the classical expression at a point a distance δ outside the horizon, $r = r_h + \delta$, where a virtual pair is created:

$$c^2 g_{tt} = - (c^2 - f(r_h + \delta)). \quad (125)$$

In the thought-experiment, we consider the pair-creation process at a distance δ inside the horizon, $r = r_h - \delta$. The normalised temporal component of the quantised metric at this interior point incorporates the quantum corrections from Propositions 4 and 6, which satisfy the operator identity (32). With these corrections, the expectation value becomes

$$\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle = -c^2 \left(1 + \frac{\chi}{c^2} - \frac{\zeta(r_h - \delta)}{c^2} + \Upsilon \right). \quad (126)$$

Eq. (126) incorporates extreme-field corrections via the speed of light correction χ , the renormalisation term Υ , and the radial function $\zeta(r)$. It yields negativity of the renormalised temporal component when the conservative conditions in Eqs. (123) and (124) hold. This negativity is necessary for a tunnelling description of interior excitations.

Building upon Eq. (126), we formulate the acausal counterpart of Eq. (98) within the extreme-field regime. Using the linearity of the expectation, $\mathbb{E}[X + Y] = \mathbb{E}[X] + \mathbb{E}[Y]$, and treating τ as a classical parameter, we obtain the analogous relation:

$$\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle \left(\frac{dt}{d\tau} \right)^2 + \langle \hat{j}_{rr} \rangle \left(\frac{dr}{d\tau} \right)^2 = -\langle \hat{\mathcal{C}}^2 \rangle. \quad (127)$$

Applying Jensen's inequality [87], which asserts that for any positive operator the reciprocal of its expectation value does not exceed the expectation

value of its reciprocal, and adopting the emergent quantum corrections truncated at a physical cut-off as proposed in Propositions 4 and 6, the quantised radial component may be approximately decomposed into a classical and an extreme contribution, $\langle \hat{j}_{rr} \rangle \approx g_{rr} + e_{rr}$. Furthermore, upon employing Eq. (126) to truncate higher-order quantum effects—thereby recovering Eq. (125) up to the leading term χ —Eq. (127) reduces to

$$[g_{tt}c^2 - \chi] \left(\frac{dt}{d\tau} \right)^2 + [g_{rr} + e_{rr}] \left(\frac{dr}{d\tau} \right)^2 = -c^2 - \chi. \quad (128)$$

This yields the conventional invariant causal term as given by Eq. (98):

$$g_{tt}c^2 \left(\frac{dt}{d\tau} \right)^2 + g_{rr} \left(\frac{dr}{d\tau} \right)^2 = -c^2. \quad (129)$$

The analogue of Eq. (129) for the acausal term χ is obtained by transforming the temporal and radial coordinates, t and r , to new coordinates \mathbf{t} and \mathbf{r} . This transformation is mediated by the quantum-gravitational temporal component \mathbf{n}_{tt} and the radial component $\mathbf{n}_{rr} = e_{rr} \left(\frac{\partial r}{\partial \mathbf{r}} \right)^2$, yielding the corresponding expression:

$$\mathbf{n}_{tt} \chi \left(\frac{d\mathbf{t}}{d\tau} \right)^2 + \mathbf{n}_{rr} \left(\frac{d\mathbf{r}}{d\tau} \right)^2 = -\chi. \quad (130)$$

Building on the preceding analysis, the tunnelling probability for the anti-eigen flux (AEF) is derived from Eq. (127), with all operators expressed through their expectation values. In general, the expectation value can be decomposed as

$$-\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle = \langle \hat{\mathcal{C}}^2 \rangle \langle -\hat{j}_{tt} \rangle + \text{Cov}(\hat{\mathcal{C}}^2, -\hat{j}_{tt}), \quad (131)$$

where $\text{Cov}(\hat{\mathcal{C}}^2, -\hat{j}_{tt})$ captures the quantum correlations. Using this decomposition, the expectation value of the radial momentum operator in the extreme-field regime can be expressed as

$$\langle \hat{p}_r \rangle = \sqrt{m^2 \langle \hat{j}_{rr} \rangle^2 \left(\frac{dr}{d\tau} \right)^2}. \quad (132)$$

Accordingly, starting from Eq. (127) and following the procedure developed in Eqs. (98)–(105), we obtain the corresponding expression for the radial momentum expectation value in the extreme-field regime as

$$\langle \hat{p}_r \rangle = \sqrt{\frac{\langle \hat{j}_{rr} \rangle E^2}{\langle \hat{\mathcal{C}}^2 \rangle \langle -\hat{j}_{tt} \rangle + \text{Cov}(\hat{\mathcal{C}}^2, -\hat{j}_{tt})}}. \quad (133)$$

For notational convenience, and building upon Proposition 2 and Remark 1, we define

$$\sqrt{\langle \hat{\mathcal{C}}^2 \rangle} \equiv c_*, \quad c_* > c. \quad (134)$$

The expectation value of the quantised temporal component, derived from Eq. (126), reads

$$\langle \hat{j}_{tt} \rangle = - \left(1 - \frac{2GM}{rc_*^2} + \Upsilon \right), \quad (135)$$

and, by invoking Jensen's inequality [87],

$$\frac{1}{\langle \hat{x} \rangle} \leq \left\langle \frac{1}{\hat{x}} \right\rangle, \quad (136)$$

in order to satisfy the Schwarzschild restriction condition from Eq. (104), it follows directly that the expectation value of the quantised radial component is bounded from below as

$$\langle \hat{j}_{rr} \rangle = \left\langle \frac{1}{-\hat{j}_{tt}} \right\rangle \geq \frac{1}{\langle -\hat{j}_{tt} \rangle}. \quad (137)$$

Utilising Eq. (137) in Eq. (133), with the perturbative definition of the speed of light c_p , and focusing on the weakly coupled regime where the covariance term can be consistently neglected, Eq. (133) simplifies to a form directly analogous to the classical radial momentum in Eq. (105), namely

$$\langle \hat{p}_r \rangle \approx \frac{E}{\langle -\hat{j}_{tt} \rangle \sqrt{\langle \hat{\mathcal{C}}^2 \rangle}} = \frac{E}{\langle -\hat{j}_{tt} \rangle c_p}. \quad (138)$$

In the general, non-perturbative case, however, the covariance term does not vanish and can become significant. Accordingly, the expectation value of the radial momentum operator is more accurately expressed as

$$\langle \hat{p}_r \rangle = \frac{E}{\sqrt{\langle -\hat{j}_{tt} \rangle^2 + \frac{\text{Cov}(\hat{\mathcal{C}}^2, -\hat{j}_{tt})}{\langle \hat{\mathcal{C}}^2 \rangle}} \sqrt{\langle \hat{\mathcal{C}}^2 \rangle}} = \frac{E}{\langle -\hat{j}_{tt}^* \rangle c_*}. \quad (139)$$

Here, Eq. (139) utilises Eq. (135), where, upon squaring, the covariance term is included within the quantum correction Υ , thereby providing a more complete definition of the temporal component operator of the Schwarzschild metric in the extreme-field, non-perturbative regime, denoted as \hat{j}_{tt}^* .

Applying the procedure of Eqs. (88)–(111) and the results of Eqs. (133)–(139), the imaginary part of the expectation value of the action operator for anti-eigen flux reduces, in the semiclassical approximation, to the radial integral of the expectation value of the radial momentum. By linearity of the expectation value under integration over r , this is expressed as

$$\Im \langle \hat{S} \rangle = \Im \left\langle \int_{r_{\text{in}}}^{r_{\text{out}}} \hat{p}_r dr \right\rangle = \Im \int_{r_{\text{in}}}^{r_{\text{out}}} \langle \hat{p}_r \rangle dr = \frac{E}{c_*} \Im \int_{r_{\text{in}}}^{r_{\text{out}}} \frac{1}{\left(1 - \frac{2GM}{rc_*^2} + \Upsilon\right)} dr. \quad (140)$$

The integrand possesses a simple pole at the quantum-corrected horizon [133]

$$r_s^* = \frac{2GM}{c_*^2(1 + \Upsilon)}. \quad (141)$$

Per Proposition 6, extreme-field corrections may be radially dependent, i.e. $\Upsilon = \Upsilon(r)$. The denominator in Eq. (140) is then

$$D(r) = 1 + \Upsilon(r) - \frac{2GM}{rc_*^2}. \quad (142)$$

The pole $r_0 = r_s^*$ follows from $D(r_0) = 0$, which via Eq. (141) reads

$$(1 + \Upsilon(r_0))r_0 = \frac{2GM}{c_*^2}. \quad (143)$$

A first-order Taylor expansion of $D(r)$ around r_0 yields

$$D(r) = D'(r_0)(r - r_0) + \mathcal{O}((r - r_0)^2). \quad (144)$$

Hence, in the vicinity of the pole, the denominator exhibits the linear behaviour

$$D(r) \simeq D'(r_0)(r - r_0). \quad (145)$$

The residue of the integrand $J(r) = 1/D(r)$ at the pole r_0 is therefore

$$\text{Res}_{r_0} J(r) = \frac{1}{D'(r_0)} = \left[\Upsilon'(r_0) + \frac{2GM}{c_*^2 r_0^2} \right]^{-1}. \quad (146)$$

To encode the contribution of the radial derivative of the correction to the residue at the pole $r_0 = r_s^*$, we introduce the compact notation ϱ and define

$$\varrho \equiv 1 + \frac{c_*^2 r_0^2}{2GM} \Upsilon'(r_s^*). \quad (147)$$

Using Eq. (141), the residue in Eq. (146) for the integrand in Eq. (140) yields the semicircular contour contribution

$$\int_{c_{1/2}} J(r) dr = \pi i \operatorname{Res}_{r_s^*} J(r) = \pi i \frac{2GM}{c_*^2(1 + \Upsilon)^2 \varrho}. \quad (148)$$

Including the conjugate contour contribution from the eigen-Hawking counterpart, which yields an identical term, the contour integral evaluates to

$$\oint J(r) dr = 2\pi i \frac{2GM}{c_*^2(1 + \Upsilon)^2 \varrho}. \quad (149)$$

This fully determines the total imaginary contribution to the action in Eq. (140), given by

$$\Im \langle \hat{S}_{\text{total}} \rangle = \frac{4\pi GME}{c_*^3(1 + \Upsilon)^2 \varrho}. \quad (150)$$

Accounting for the back-reaction of the positive-energy eigen-Hawking counterpart, all correction terms arising from the back-reaction integral [121] are fully absorbed into the extreme-field corrections Υ and ϱ , thereby preserving exact consistency between Eqs. (150), (151) and their classical counterparts (111), (112). Hence, the corrected tunnelling probability for the anti-eigen flux becomes

$$\Gamma_{\text{AEF}}^{\text{BR}} \sim \exp\left[-\frac{8\pi GE}{\hbar c_*^3(1 + \Upsilon)^2 \varrho} \left(M + \frac{E}{2c_e^2}\right)\right] = \exp\left[\frac{-8\pi GE \Delta^+}{\hbar c_*^3(1 + \Upsilon)^2 \varrho}\right]. \quad (151)$$

It is worth noting that, in the limit $\chi \rightarrow 0$, the speed of light recovers its constancy, $c_* \rightarrow c$ (Eq. (34)). In this classical limit, the quantum corrections and their derivatives vanish ($\Upsilon, \Upsilon' \rightarrow 0$), so $\varrho \rightarrow 1$ (Proposition 6); moreover, the energy term becomes negative, $E \rightarrow -E$, implying $\Delta^+ \rightarrow \Delta^-$ (Subsection 3.5). Hence, the tunnelling probability of the anti-eigen flux in Eq. (151) reduces to that of standard Hawking radiation in Eq. (112):

$$\lim_{\chi \rightarrow 0} \Gamma_{\text{AEF}}^{\text{BR}} = \lim_{\chi \rightarrow 0} \exp\left[-\frac{8\pi GE}{\hbar c_*^3(1 + \Upsilon)^2 \varrho} \left(M + \frac{E}{2c_e^2}\right)\right] = \exp\left[-\frac{8\pi GE}{\hbar c^3} \left(M - \frac{E}{2c^2}\right)\right] = \Gamma_{\text{SHR}}^{\text{BR}}. \quad (152)$$

Thus, in the classical limit, the extreme-field degrees of freedom vanish, and the modified expression consistently reduces to the standard Hawking tunnelling probability—confirming the physical validity of the framework.

Based on these results, three primary scenarios emerge that could describe the primordial universe; these will be detailed in the subsequent subsections.

3.1 Scenario 1: The Classical Framework of Pre-Being Fluctuations

During the Pre-Being Fluctuations epoch (PBF) the fluctuations obey the classical framework of special relativity, characterised by the vanishing of the quantised operator $\hat{\chi}$ (i.e. $\chi = 0$). This means the fluctuations adhere to the established laws of classical physics. In this case, taking $\zeta(r_h - \delta) = f(r_h - \delta)$ and using Eq. (114) for $r < r_h$, the normalised temporal component of the jerking metric, Eq. (126), becomes positive:

$$\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle = - (c^2 - \zeta(r_h - \delta)) = -c^2 \left(1 - \frac{f(r_h - \delta)}{c^2} \right) > 0. \quad (153)$$

Hence, no fluctuation can tunnel outward; all remain trapped within the primordial black holes — precisely as the classical picture dictates. A strictly causal framework that respects translational symmetry—in which symmetric fluctuations follow a purely linear distribution—lacks the necessary nonlinearities and couplings to support black-hole formation. Furthermore, the pre-being paradox precludes the first-principle existence of such fluctuations. Consequently, this scenario cannot satisfy the conditions required for the emergence of a universe and remains fundamentally unresolved.

3.2 Scenario 2: Perturbative Framework of Primordial Black Holes

Primordial black hole formation is interpreted within the perturbative framework as arising from extreme-field nonlinear fluctuations in the pre-being epoch. This nonlinearity stems from a perturbatively quantised metric (Proposition 6, Eqs. (76)–(79)), wherein recursive operator-level coupling generates a nonlinear back-reaction among metric components—a result formalised further in Theorem 3 (Eqs. (258)–(270)). At its core, this behaviour originates from the perturbed speed of light in acausal extreme field, promoted to a quantum operator in Proposition 1 and defined in Eq. (84), where v satisfies

$$c_p \gtrsim c, \quad v \gtrsim 0. \quad (154)$$

Consequently, the temporal component of the metric in the extreme-field region near the inner horizon, g_{tt}^p , retains its timelike character. Under this perturbative regime—where $c_* \approx c_p$ marks the perturbative limit of Eq. (134),

the quantum corrections reduce perturbatively to $\Upsilon \equiv \Upsilon_p$ and $\varrho \equiv \varrho_p$, and the covariance term in Eq. (131) is negligible—Eq. (135) therefore reads

$$g_{tt}^p = - \left(1 - \frac{2GM}{rc_p^2} + \Upsilon_p \right) < 0. \quad (155)$$

Eq. (155) therefore signals acausal excitations—the anti-eigen flux—capable of tunnelling through a perturbatively quantum-corrected horizon:

$$r_s^p = \frac{2GM}{c_p^2(1 + \Upsilon_p)}. \quad (156)$$

Equality between the tunnelling probabilities for standard Hawking radiation, Eq. (112), and for the anti-eigen flux, Eq. (151), holds precisely when a perturbative, extreme-field equality is satisfied. This symmetry condition is

$$\frac{\Delta^+}{c_p^3(1 + \Upsilon_p)^2 \varrho_p} = \frac{\Delta^-}{c^3} \iff \Gamma_{\text{AEF}}^{\text{BR}} = \Gamma_{\text{SHR}}^{\text{BR}}. \quad (157)$$

This symmetric tunnelling pattern, given by Eq. (157), establishes a dynamical balance that governs both primordial and present-day astrophysical black holes. It reveals a symmetric correlation between two non-local points—outside and inside the horizon [140]. This symmetry ensures equal tunnelling probabilities for Hawking radiation and anti-eigen flux, thereby regulating information exchange across the horizon through precise symmetry rather than irreversible loss.

This behaviour is thus visualised as a non-classical respiratory process: negative-energy Hawking flux falls inward, while anti-eigen flux tunnels outward with equal probability. The black hole accordingly ‘breathes’—inhaling positive-energy eigen-Hawking radiation and exhaling positive-energy standard Hawking radiation—in a balanced exchange that persists without complete evaporation. Black holes are therefore not perfect information sinks; they conserve entropy [142] and release information as anti-eigen flux through their perturbatively quantum-corrected horizon.

This intuitive picture motivates the following proposition: global purity of the total quantum state, together with reproduction of the Page curve, resolves the black hole information-loss paradox [4, 107].

Proposition 8 (Global Purity and Page-Curve Reproduction). *Consider a black hole and its associated radiation as a closed quantum system, initially*

prepared in a pure state $|\Psi_0\rangle$. Within the perturbative extreme-field regime, the symmetric tunnelling correlations encoded in Eq. (157) ensure that the total system—comprising the black-hole interior B , the radiation sector R , and the anti-eigen flux F —evolves unitarily. These correlations induce pairwise entanglement between exterior radiation modes and interior anti-eigen excitations, allowing a Schmidt decomposition of the total quantum state [117]:

$$|\Psi\rangle = \sum_i c_i |i\rangle_R \otimes |i\rangle_{B+F}, \quad (158)$$

which is manifestly pure. The corresponding density matrix $\rho_{\text{total}} = |\Psi\rangle\langle\Psi|$ satisfies $\rho_{\text{total}}^2 = \rho_{\text{total}}$, establishing the vanishing global von Neumann entropy

$$S_{\text{total}} = -\text{Tr}(\rho_{\text{total}} \ln \rho_{\text{total}}) = 0, \quad (159)$$

in accordance with [156]. This global purity holds regardless of whether the black hole evaporates; any apparent mixedness of the radiation arises solely from tracing over inaccessible interior and anti-eigen degrees of freedom [108, 122]. In particular, although individual outgoing radiation modes may appear locally mixed and thermal [65], this reflects only partial tracing, not a fundamental loss of information [143].

For the radiation subsystem R , the entanglement entropy follows the Page curve: it rises initially as correlations build, peaks near the Page time, and decreases as correlations are transferred outward, providing a clear dynamical signature of unitary evolution [119, 120, 134]. Within the perturbative framework, this behaviour emerges once the exact symmetric tunnelling probabilities in Eq. (157) are perturbed to near symmetry. Concretely, the perturbative quantum corrections induce a controlled deviation from exact tunnelling symmetry, expressed as

$$\Upsilon_p = \sqrt{e_p^{-1} p^3 \frac{\Delta^+}{\Delta^-} - 1} + \xi, \quad (160)$$

where $p \equiv c/c_p$, and ξ quantifies the residual departure from exact symmetry. This deviation reflects the intrinsically probabilistic character of extreme-field perturbations (see Proposition 6). Consequently, from Eqs. (112) and (151), for sufficiently small ξ the perturbations remain nearly classical, and standard Hawking radiation dominates at early times; for larger ξ , the perturbative effects become more pronounced, and the anti-eigen flux grows significant after the Page time. This complementary interplay between the two fluxes naturally reproduces the characteristic Page curve.

However, the theory offers a more coherent picture through its perturbative extreme-field equality, wherein information is recovered continuously via the anti-eigen flux, rather than only after the Page time or at the end of evaporation. In this picture, the Page curve remains meaningful but should be interpreted as a signature of continuous information recovery. This ensures unitarity at all times, preserves global purity ($S_{\text{total}} = 0$), and further suggests that black holes may persist as stable reservoirs of information rather than fully evaporating.

Before addressing the non-perturbative regime, it is useful to examine the perturbative framework through the observational lens of quasi-normal modes (QNMs). These modes characterise the ringdown stage of gravitational wave signals, enabling accurate determination of black hole masses and spins, and offering novel avenues for testing general relativity [25].

Quasi-normal modes are the discrete, damped oscillations that characterise the linear response of a black hole spacetime to perturbations. They are independent of the excitation mechanism and determined solely by the black hole’s macroscopic parameters (mass, charge, and angular momentum) together with the boundary conditions imposed at the horizon and asymptotic infinity—making them sensitive probes of possible deviations from classical geometry [93]. For a Schwarzschild black hole, after separating angular variables and introducing the tortoise coordinate $r_* = r + 2M \ln(r/2M - 1)$, the radial component of a perturbation outside the event horizon satisfies the wave equation

$$\left(-\frac{\partial^2}{\partial r_*^2} + V_\ell(r)\right) \psi_\ell(r) = -\frac{\partial^2}{\partial t^2} \psi_\ell(r), \quad (161)$$

$$\left(-\frac{\partial^2}{\partial r_*^2} + V_\ell(r)\right) \psi_\ell(r) = \omega_{\ell n}^2 \psi_\ell(r). \quad (162)$$

The effective potential for axial perturbations (Regge–Wheeler) is given by

$$V_\ell^{(\text{RW})}(r) = \left(1 - \frac{2M}{r}\right) \left(\frac{\ell(\ell+1)}{r^2} - \frac{2\sigma M}{r^3}\right), \quad (163)$$

where ℓ is the spherical harmonic index, n the overtone number ($n = 0$ denotes the fundamental mode), and $\sigma = 1 - s^2$ labels the perturbation type, with spin weight $s = 0$ (scalar), $s = 1$ (electromagnetic), and $s = 2$ (gravitational), yielding $\sigma = 1, 0, -3$ respectively [42, 93].

Subject to boundary conditions of purely ingoing waves at the event horizon and purely outgoing waves at spatial infinity, the frequency spectrum becomes discrete and complex, $\omega = \omega_R - i\omega_I$, where ω_R determines the oscillation frequency and $\omega_I > 0$ encodes the damping rate due to energy loss through gravitational radiation [25, 93].

In recent years, significant advances have been made in computing QNMs for non-standard geometries. Of particular relevance to the present work is the spectral method developed by Batic, Dutykh and collaborators, which reformulates the perturbation equations as an eigenvalue problem over a compact domain using Chebyshev polynomials [14, 15]. This approach yields high-precision spectra and has been successfully applied to noncommutative geometry-inspired black holes [15, 17], wormholes [14, 18], and Lee-Wick black holes [19]. These studies demonstrate remarkable stability for non-extremal and extremal configurations [15], reveal the emergence of purely imaginary modes in near-extremal regimes [19], and establish that QNMs converge to classical Schwarzschild values for large mass parameters [17]. The framework has also been extended to analyse instabilities in phantom wormholes [16], providing a comprehensive toolkit for probing quantum-gravitational corrections via gravitational wave spectroscopy.

Recent reviews by Berti et al. [26] and Destounis & Duque [51] provide a comprehensive overview of black hole spectroscopy and its observational context in light of current and future gravitational wave detectors. These works examine the interplay between theoretical QNM spectra and gravitational wave measurements, offering essential background for the potential observational signatures of Extreme Field Theory discussed below.

The relevance to extreme field theory is twofold. First, dirty black hole models [17] parallel our operator-valued metric and quantised speed of light through their noncommutative, singularity-regularising structure. Second, astrophysical black holes in our framework occupy the perturbative regime, where symmetric tunnelling and global purity hold (Proposition 8). These systems thereby offer direct observational tests via quasi-normal modes—a connection we now formalise.

Proposition 9 (QNM imprints of extreme-field perturbations). *Perturbative extreme field theory imparts new internal degrees of freedom to the metric—arising from microscopic geometric quantities elevated to quantum operators—permitting internal acausal excitations, the anti-eigen flux, to tunnel outward. Consequently, the black hole’s microscopic quantum structure leaves*

a distinct imprint on its quasi-normal mode spectrum.

While QNMs are classically determined solely by the black hole’s macroscopic hair (M , Q , and J), the perturbed anti-eigen flux introduces subtle corrections to the dissipative dynamics. These corrections arise from three sources: the perturbed speed-of-light component v , the extreme-field perturbative corrections Υ_p , and the interior back-reaction Δ^+ . Collectively, they modify the effective potential governing spacetime perturbations—modifications that emerge intrinsically from perturbations in the metric tensor.

These metric perturbations are obtained by starting from a spherically symmetric Schwarzschild black hole, whose classical line element is

$$ds^2 = g_{\mu\nu}^0 dx^\mu dx^\nu = -e^{\nu(r)} dt^2 + e^{\lambda(r)} dr^2 + r^2(d\theta^2 + \sin^2\theta d\phi^2). \quad (164)$$

Classical linear perturbations then take the standard form

$$g_{\mu\nu} = g_{\mu\nu}^0 + h_{\mu\nu}. \quad (165)$$

In the extreme-field regime, the metric acquires an additional correction of the form

$$g_{\mu\nu}^p \approx g_{\mu\nu}^0 + h_{\mu\nu} + e_{\mu\nu}, \quad (166)$$

where $e_{\mu\nu}$ encodes interior perturbative excitations arising from the extreme field and depends on v , Υ_p , and Δ^+ .

This leads to a concrete observational prediction: the damping factors (imaginary parts) of the QNM frequencies receive small corrections,

$$\delta(\omega_I) = \epsilon \cdot f(v, \Upsilon_p, \Delta^+) + \mathcal{O}(\epsilon^2), \quad (167)$$

where $\epsilon \ll 1$ denotes the perturbative scale, and f encodes the dependence on the extreme-field parameters v , Υ_p , and Δ^+ .

Such modifications fall within reach of high-precision gravitational wave spectroscopy using current and next-generation observatories—LIGO, Virgo, KAGRA, the Einstein Telescope, and LISA. The spectral methods developed by Batic, Dutykh and collaborators [14, 15, 17] enable precise quantification of these deviations, establishing a direct link between the theoretical predictions of Extreme Field Theory and future observational data.

Incorporating the extreme-field correction proportional to v into a toy model produces a monotonic increase in both ω_R and ω_I . At $v = 0.20$, the damping shift reaches $\sim 2.2\%$, surpassing the 1% sensitivity threshold of next-generation observatories.

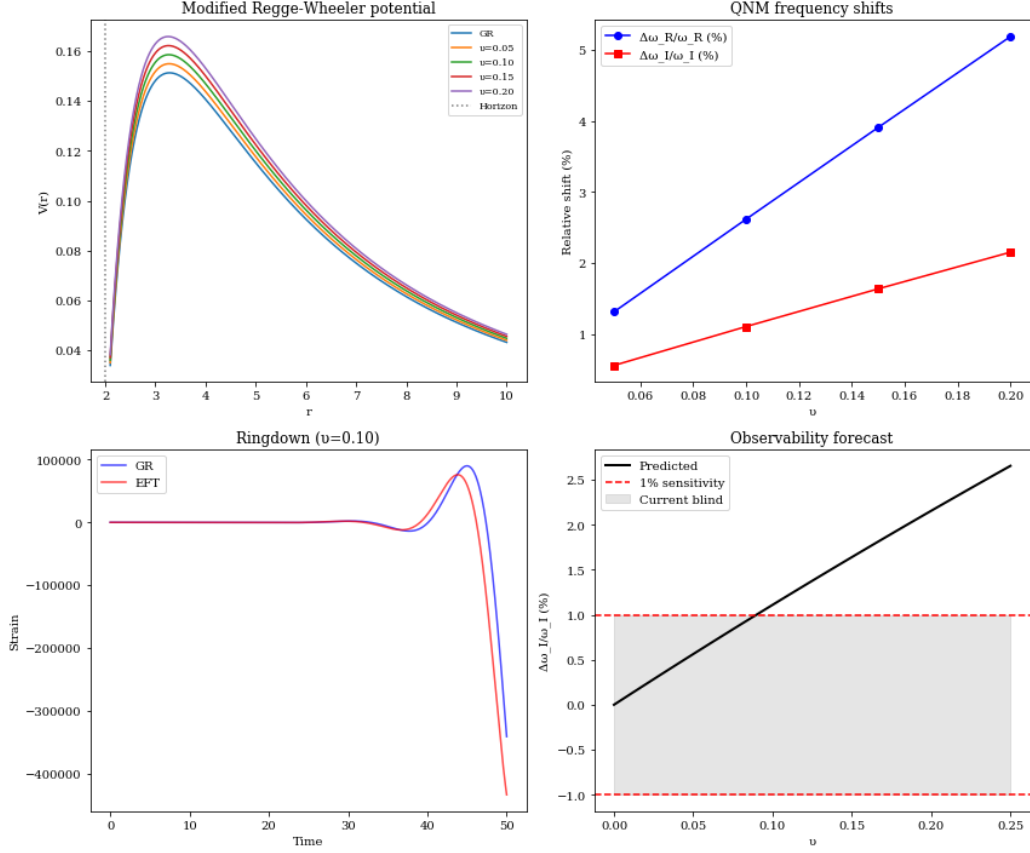


Figure 2: Extreme-field corrections to quasi-normal modes. **(a)** Modified Regge–Wheeler potential for different values of the extreme-field parameter v (amplitude $A = 2$). The correction is localised near the peak and vanishes at the horizon. **(b)** Relative shifts of the real and imaginary parts of the fundamental QNM frequency as functions of v . Both increase monotonically, with $\Delta\omega_R/\omega_R \sim 5.2\%$ and $\Delta\omega_I/\omega_I \sim 2.2\%$ at $v = 0.20$. **(c)** Example ringdown waveforms for GR (blue) and with extreme-field corrections at $v = 0.10$ (red). The faster damping reflects the increased ω_I . **(d)** Observability forecast: fractional change in damping rate $\Delta\omega_I/\omega_I$ vs. v . The grey band marks the 1% sensitivity threshold of next-generation detectors (Einstein Telescope, LISA); for $v \gtrsim 0.08$ the predicted shift exceeds this level.

Table 1: Numerical values of the QNM frequencies and their relative shifts for the extreme-field parameter v , obtained from the toy model with $A = 2$.

v	ω_R	ω_I	$\Delta\omega_R/\omega_R$ (%)	$\Delta\omega_I/\omega_I$ (%)
0.00	0.430845	0.256875	0.0000	0.0000
0.05	0.436517	0.258322	1.3166	0.5634
0.10	0.442131	0.259726	2.6196	1.1099
0.15	0.447688	0.261088	3.9093	1.6402
0.20	0.453188	0.262410	5.1860	2.1549

The toy model in Fig. 2 and Table 1 illustrates how extreme-field corrections affect the quasi-normal modes of a Schwarzschild black hole.

The correction term $\delta V = A v f(r)/r^2$, where $f(r) = 1 - 2M/r$, is chosen to vanish at the horizon and fall off at infinity, mimicking the qualitative behaviour of the operator-valued perturbations $e_{\mu\nu}$ introduced in Proposition 9. The amplitude $A = 2$ is set to produce shifts at the percent level, the typical scale expected from perturbative quantum-gravity effects.

As v increases, both the oscillation frequency ω_R and the damping rate ω_I increase monotonically; at $v = 0.20$ the damping shift reaches 2.15% (Fig. 2b), comfortably above the 1% sensitivity threshold of next-generation gravitational-wave observatories (Einstein Telescope, Cosmic Explorer, LISA). The waveform comparison for $v = 0.10$ (Fig. 2c) shows a visibly faster ring-down, demonstrating that the effect is in principle observable.

Extreme Field Theory therefore predicts inner black-hole degrees of freedom—encoded in extreme-field perturbations and governing anti-eigen flux tunnelling—that are testable via high-precision gravitational-wave spectroscopy.

3.3 Scenario 3: The Non-Perturbative Primordial Black Hole

Following Proposition 8, the perturbative quantisation approach preserves the purity of Hawking radiation globally—excluding mixed and thermal features from the global state while maintaining entropy conservation. Under this framework, a perturbative primordial black hole may reach a threshold at which its acausal fields saturate the lower boundary of non-perturbative quantisation, as illustrated in Fig. 3a. This saturation provides a sufficient

physical condition for an inner thermal transition that triggers a global entropy difference—strictly speaking, a negentropic departure—leading to the formation of a negentropic path towards an extreme-Planckian state. This culminates in a Primordial Planckian Black Hole (PPBH): a Planck-scale seed from which the universe can emerge, as will be discussed in detail in Subsection 3.7 and the main Section 4.

At this threshold, the root mean square of the non-perturbative quantised speed of light in Eq. (87)—denoted c_* in Eq. (134)—becomes significantly enhanced yet bounded, as established in Proposition 2, satisfying the non-perturbative enhancement condition

$$c_* > c, \quad \langle \hat{\chi} \rangle = \chi \gg 0. \quad (168)$$

Furthermore, the non-perturbative quantum corrections Υ_{NP} and ϱ become non-trivial, with the covariance term in Eq. (131) growing significant—ensuring that the temporal component of the metric retains a high probability of preserving its timelike character (i.e., maintaining the Lorentzian signature), thereby upholding the causal structure of spacetime, in agreement with Theorem 1, and analogous to the perturbative result in Eq. (155).

This outcome explicitly violates the perturbative, extreme-field equality presented in Eq. (157). Consequently, an intrinsic spacetime asymmetry emerges, characterised by a dominant tunnelling probability for the anti-eigen flux over that of standard Hawking radiation. This asymmetry is mathematically captured by a non-perturbative, extreme-field inequality, which holds under the strict condition

$$\frac{\Delta^+}{c_*^3(1 + \Upsilon_{\text{NP}})^2\varrho} < \frac{\Delta^-}{c^3} \iff \Gamma_{\text{AEF}}^{\text{BR}} > \Gamma_{\text{SHR}}^{\text{BR}}. \quad (169)$$

Eq. (169) implies that the creation rate \mathbf{R} of eigen-Hawking radiation (EHR) inside the non-perturbative primordial black hole exceeds the creation rate of standard Hawking radiation (SHR) emitted from its outer horizon. This relationship can be written compactly as

$$\mathbf{R}_{\text{EHR}} > \mathbf{R}_{\text{SHR}}. \quad (170)$$

Eqs. (169) and (170) establish the fundamental asymmetry that allows a primordial black hole to collapse to a Planckian-threshold black hole, which at this extreme juncture can recoil and seed large-scale cosmic structure.

3.4 Hawking Radiation Tunnelling: General Solution and Time Interpretation

In the preceding subsections, we derived the tunnelling probability for Hawking radiation under two assumptions: the high-energy limit near the black hole horizon (Eq. (103)) and the condition that primordial black holes are non-rotating and spherically symmetric, consistent with the Schwarzschild solution (Eq. (104)). To generalise this analysis, we now relax the latter assumption and treat g_{rr} and $-g_{tt}$ as independent metric components, rather than expressing one in terms of the other. Additionally, we incorporate the positive-definite function Ω , which appears under a square root and depends on the angular coordinates, their associated metric components, and the radial coordinate. Previously omitted in Eq. (98) under the assumption of negligible angular contributions, this function allows the radial momentum to be expressed in its general form, thereby generalising Eq. (105) as follows:

$$p_r = \frac{\sqrt{g_{rr}}}{\sqrt{-c^2 g_{tt}}} \Omega E. \quad (171)$$

Therefore, the tunnelling probability for standard Hawking radiation, given in Eq. (112), can be expressed in its general form, Eqs. (88) and (89), as:

$$\Gamma_{\text{SHR}} \sim \exp \left(- \frac{2E}{\hbar} \left[\Im \int \frac{\sqrt{g_{rr}}}{\sqrt{-c^2 g_{tt}}} \Omega dr \right] \right). \quad (172)$$

In a similar vein, for the anti-eigen flux defined by the quantised metric in Eq. (116), we consider the expectation value of the combined radial and temporal metric operators under a single square root. This preserves the full expression $\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle$, without the need to separate out the emergent covariance term, thereby avoiding additional restrictions or hidden assumptions introduced solely to simplify the mathematical treatment. Furthermore, we introduce an angular function $\bar{\Omega}$ that accounts for the dependence on the angular coordinates, their associated metric components, and the radial coordinate. Consequently, the general form of the tunnelling probability for the anti-eigen flux can be written as:

$$\Gamma_{\text{AEF}} \sim \exp \left(- \frac{2E}{\hbar} \left[\Im \int \frac{\sqrt{\langle \hat{j}_{rr} \rangle}}{\sqrt{-\langle \hat{\mathcal{C}}^2 \hat{j}_{tt} \rangle}} \bar{\Omega} dr \right] \right). \quad (173)$$

Although the general tunnelling probabilities for standard Hawking radiation (SHR, Eq. (172)) and the anti-eigen flux (AEF, Eq. (173)) cannot be fully computed due to the relaxation of the Schwarzschild assumption for primordial black holes, the aim of this subsection is not to evaluate the integrals explicitly, but to elucidate their physical significance. Dimensional analysis shows that the imaginary parts of these integrals necessarily have the dimension of time.

For the standard Hawking radiation (SHR), the imaginary part of the integral arises from a simple pole on the real axis at t_h , interpreted as the classical temporal horizon. The contour is taken above the real axis, forming a counterclockwise (CCW) semicircle in the upper half-plane. By the residue theorem, a simple pole enclosed CCW contributes

$$\oint_{\text{CCW}} f(z) dz = +2\pi i \text{Res}[f, z_0]. \quad (174)$$

Since the pole lies on the real axis, a semicircular $i\epsilon$ deformation contributes only half of the residue, yielding π . Hence, for $r = r(t)$, the imaginary part of the integral in Eq. (172) may be evaluated by transforming the integration variable from the radial coordinate to time, giving

$$\Im \int \frac{\sqrt{g_{rr}}}{\sqrt{-c^2 g_{tt}}} \Omega dr = \Im \int \frac{\sqrt{g_{rr}}}{\sqrt{-c^2 g_{tt}}} \Omega \left(\frac{dr}{dt} \right) dt = +\pi f(t_h) \sim \mathbb{T}, \quad (175)$$

where \mathbb{T} carries the dimension of time.

For the anti-eigen flux (AEF), the integral has a pole at τ_h , corresponding to a quantum-escapable temporal horizon, which may differ from the classical one, analogous to the quantum-corrected Schwarzschild horizon in Eq. (141). The contour is taken below the real axis, forming a clockwise (CW) semicircle. For a simple pole enclosed CW, the residue theorem gives

$$\oint_{\text{CW}} f(z) dz = -2\pi i \text{Res}[f, z_0]. \quad (176)$$

By the same reasoning, the semicircular deformation contributes half the residue with opposite sign. Thus, the imaginary part of the integral in Eq. (173) becomes

$$\Im \int \frac{\sqrt{\langle \hat{j}_{rr} \rangle}}{\sqrt{-\langle \hat{C}^2 \hat{j}_{tt} \rangle}} \bar{\Omega} dr = \Im \int \frac{\sqrt{\langle \hat{j}_{rr} \rangle}}{\sqrt{-\langle \hat{C}^2 \hat{j}_{tt} \rangle}} \bar{\Omega} \left(\frac{dr}{dt} \right) dt = -\pi f(\tau_h) \sim -\omega \mathbb{T}, \quad (177)$$

where $\omega \in [0, 1]$ is the extreme-field parameter arising from the quantisation of the speed of light, giving the black hole interior a quantum character capable of supporting acausal anti-eigen excitations.

Comparing Eqs. (175) and (177), an observer at infinity perceives standard Hawking radiation (SHR) as evolving forward in time, while the anti-eigen flux (AEF) effectively evolves backward. For $\omega = 1$, the tunnelling probabilities of SHR and AEF are symmetric, consistent with the perturbative quantisation of the speed of light discussed in Section 3.2 and with the perturbative, extreme-field equality of Eq. (157). This symmetry ensures that tunnelling probabilities remain finite inside the black hole and that the evolution is unitary, $\hat{U}\hat{U}^\dagger = \mathbb{I}$, so that outgoing states remain pure rather than becoming mixed. Consequently, information can escape the black hole through the *temporally escapable quantum horizon*, providing a coherent and symmetric mechanism that preserves both information and unitarity, as demonstrated in Proposition 8.

The uniquely non-perturbative approach developed in Section 3.3 for a primordial black hole establishes a lower bound for non-perturbative quantisation. This bound is achieved if and only if, for such a primordial black hole, the extreme-field parameter satisfies

$$\omega \leq 1. \tag{178}$$

Applying this condition to Eqs. (175) and (177), and inserting them into the tunnelling probabilities of Eqs. (172) and (173), directly implies the emergence of an intrinsic spacetime asymmetry, satisfying Eqs. (169) and (170). This sets the initial conditions for the emergence of the universe, as motivated in 3 and formalised in Theorem 2.

3.5 Interpretation of the Anti-Eigen Flux as Negative-Energy Particles

Remark 5. *The backward-in-time evolution of the anti-eigen flux preserves the probability bounds in Eq. (173) if and only if its energy is interpreted as a negative-energy solution.*

In this subsection, we revisit the thought experiment of Fig. 1 to interpret Remark 5. An observer at infinity imagines a hypothetical sphere whose geometry, if translated in spacetime by the horizon diameter $2r_h$, would ex-

actly coincide with a primordial black hole (PBH). For this observer, standard Hawking radiation (SHR) emerges outward at distance δ from the outer horizon, while its negative-energy counterpart falls inward. Simultaneously, acausal PBH excitations generate anti-eigen flux at the same distance δ from the inner horizon, propagating toward negative infinity.

For the observer, outgoing positive flux balanced by infalling negative energy characterises a perturbative PBH with intact translational symmetry. Violation of this symmetry instead reveals an intrinsic spacetime asymmetry—realised uniquely in the non-perturbative PBH—where enhanced anti-eigen flux tunnelling drives a negentropic trajectory observable toward the Planckian scale.

To formalise this, recall Eq. (172), defining the Bekenstein tunnelling power B as

$$B \equiv \frac{2}{\hbar} \Im \int \frac{\sqrt{g_{rr}}}{\sqrt{-c^2 g_{tt}}} \Omega E dr. \quad (179)$$

As established in Eq. (89), for SHR the integral is evaluated over the radial domain from r_{in} to r_{out} , corresponding to outgoing positive-energy radiation, yielding

$$B_{\text{SHR}} \equiv \frac{2}{\hbar} \Im \int_{r_{\text{in}}}^{r_{\text{out}}} \frac{\sqrt{g_{rr}}}{\sqrt{-c^2 g_{tt}}} \Omega E_{\text{SHR}} dr = B. \quad (180)$$

For the anti-eigen flux, which propagates away from the observer toward negative infinity (opposite to the SHR direction), the integral effectively reverses its boundaries relative to that of SHR. Consequently, from Eq. (173), the Bekenstein tunnelling power for the anti-eigen flux becomes

$$B_{\text{AEF}} \equiv -\frac{2}{\hbar} \Im \int_{r_{\text{in}}}^{r_{\text{out}}} \frac{\sqrt{\langle \hat{j}_{rr} \rangle}}{\sqrt{-\langle \hat{C}^2 \hat{j}_{tt} \rangle}} \bar{\Omega} E_{\text{AEF}} dr \propto -B. \quad (181)$$

The negative sign in Eq. (181) can be interpreted as the action of a parity operator \hat{P} on the SHR Bekenstein tunnelling power B , yielding

$$\hat{P}B = -B. \quad (182)$$

Employing Eq. (181) together with Eq. (177), where the factor ω emerges as the extreme-field parameter, the tunnelling probability for the anti-eigen flux reads

$$\Gamma_{\text{AEF}} \sim e^{-B_{\text{AEF}}} = e^{\omega B}. \quad (183)$$

At first sight, Eq. (183) appears counterintuitive: the tunnelling probability grows exponentially, seemingly violating unitarity. This apparent paradox is resolved by recognising that the Bekenstein tunnelling power associated with the anti-eigen flux must be intrinsically negative. Consistency with the definition of the Bekenstein tunnelling power for standard Hawking radiation in Eq. (180) then requires

$$E_{\text{AEF}} < 0 \quad \iff \quad \Gamma_{\text{AEF}} \sim e^{-B_{\text{AEF}}} = e^{-\omega B}, \quad (184)$$

which restores the correct probabilistic bound and implies that the anti-Eigen flux is necessarily a negative-energy excitation.

This result implies that, for primordial black holes, the anti-eigen flux represents an outward-flowing negative-energy flux, while the eigen-Hawking radiation corresponds to an inward-flowing positive-energy flux. This mechanism prevents PBHs from evaporating via standard Hawking radiation. Within the perturbative framework (Section 3.2), this behaviour ensures the preservation of information, whereas in the non-perturbative framework (Section 3.3), it allows a primordial black hole, saturated at the lower bound of non-perturbative quantisation, to follow a negentropic path without evaporation, thereby establishing the initial conditions for the emergence of the universe. This scenario will be discussed in further detail in Section 3.7.

3.6 The Origin of Matter–Antimatter Asymmetry

The cosmological matter–antimatter asymmetry remains a fundamental open problem: despite the symmetry of microscopic laws, the Universe is overwhelmingly matter-dominated. Within extreme-field theory, this asymmetry originates from the interior of a non-perturbative primordial black hole (NP-PBH), where non-perturbative quantisation yields a dynamical bias favouring matter confinement and antimatter deconfinement—the latter via an enhanced tunnelling probability.

At a fundamental level, two possibilities may account for the observed asymmetry: (i) an intrinsic asymmetry in matter creation, or (ii) symmetric creation of matter and antimatter followed by a biased dynamical separation. The first possibility would require the emergence of matter without corresponding antimatter partners. Such a process cannot arise from vacuum quantum fluctuations and contradicts the foundational principle of extreme-field theory, according to which extreme fluctuations are inevitable in the pre-being epoch.

The second possibility, involving symmetric creation with subsequent dynamical bias, is consistent with the principles of the theory. In this picture, matter–antimatter pairs are created symmetrically, but the extreme-field environment within the NP–PBH induces an asymmetric evolution: matter preferentially accumulates as inward eigen-Hawking radiation, while antimatter is preferentially expelled as anti-eigen flux. This biased separation provides a natural mechanism for generating the observed asymmetry.

To assess whether such a bias could instead arise from causal fluctuations at the outer horizon, we examine standard Hawking radiation. In this case, particle–antiparticle pairs are produced near the horizon, with one partner tunnelling outward while the other falls inward. However, the universal tunnelling probability in Eq. (172) depends only on the energy E and the classical metric components in Eq. (113), both of which are insensitive to particle identity. Consequently, standard Hawking radiation is strictly symmetric between matter and antimatter,

$$\Gamma_{\text{SHR-Matter}} = \Gamma_{\text{SHR-Antimatter}}. \quad (185)$$

This symmetry rules out standard Hawking radiation as the origin of the cosmological asymmetry.

This suggests that the matter–antimatter asymmetry originates within the non-perturbative primordial black hole (NP–PBH), where extreme-field effects break the symmetry via a biased tunnelling mechanism.

Based on the above considerations, we now examine the physical resolution of the matter–antimatter asymmetry by analysing the tunnelling probability of the anti-eigen flux and determining whether it exhibits symmetry or antisymmetry between matter and antimatter. The universal expression for the anti-Eigen flux tunnelling probability follows from Eqs. (173), (177), and (184), and may be written as

$$\Gamma_{\text{AEF}} \equiv \exp\left(-\frac{2}{\hbar} \Im \int \frac{\sqrt{\langle \hat{j}_{rr} \rangle}}{\sqrt{-\langle \hat{C}^2 \hat{j}_{tt} \rangle}} \bar{\Omega} E dr\right) = \exp\left(\frac{-2\omega E^-[-\text{T}]}{\hbar}\right) = e^{-\omega B}. \quad (186)$$

In Eq. (186), E^- denotes the negative-energy solution associated with the anti-eigen flux, while $(-\text{T})$ represents its backward-in-time evolution. Within the non-perturbative framework, the parameter ω emerges intrinsically as an extreme-field fraction satisfying the bound in Eq. (178), with $0 < \omega \leq 1$.

Within the quantum-field-theoretic picture, matter and antimatter can be distinguished by their temporal orientation: matter corresponds to positive-energy states propagating forward in time, while antimatter corresponds to negative-energy states evolving backward in time. In direct analogy with Eq. (186), only the latter—the negative-energy, backward-evolving modes—are permitted to tunnel outward as anti-eigen flux. Conversely, positive-energy, forward-evolving modes cannot tunnel outward; they are confined to propagate inward, seeding the NP–PBH core.

Consequently, we identify the anti-eigen flux with antimatter that escapes via tunnelling, and the eigen-Hawking radiation with matter that remains trapped, collapsing inward instead. We therefore identify this solution as a *PT-symmetric evolution*, in which the two sectors—anti-eigen flux and eigen-Hawking radiation—are related as mirror images under spatial inversion (P) and time reversal (T) [22].

A more general resolution arises when the extreme-field parameter itself acquires an intrinsic asymmetry within the non-perturbative framework. In this case, Eq. (186) allows for distinct values of the extreme-field parameter for matter and antimatter, satisfying the inequality

$$\omega_{\text{antimatter}} < \omega_{\text{matter}}. \quad (187)$$

Substituting Eq. (187) into Eq. (186) immediately yields

$$F_{\text{EHR-Matter}} > F_{\text{EHR-Antimatter}}, \quad (188)$$

$$\Gamma_{\text{AEF-Matter}} < \Gamma_{\text{AEF-Antimatter}}, \quad (189)$$

where F denotes the probability for inward-directed eigen-Hawking radiation and is defined as the complement of the tunnelling probability,

$$F = 1 - \Gamma. \quad (190)$$

It is therefore natural to interpret ω as an *extreme-field eigen-intrinsic asymmetry parameter* that fundamentally differentiates matter from antimatter within the quantum-gravitational interior of the NP–PBH. Under the condition in Eq. (187), a slightly larger value of ω for matter induces an effective confinement: the tunnelling probability is suppressed, while the probability of inward-directed eigen-Hawking radiation is enhanced. Conversely, antimatter—characterised by a marginally smaller ω —experiences an effective deconfinement, manifested as an enhanced tunnelling probability via anti-eigen flux and a reduced inward flux probability.

From first principles, the matter-confinement–antimatter-deconfinement dynamics follows from the extreme-field integral $\mathbf{I} \equiv \Im \int \frac{\sqrt{\langle \hat{j}_{rr} \rangle}}{\sqrt{-\langle \hat{C}^2 \hat{j}_{tt} \rangle}} \bar{\Omega} E dr$ as defined in Eq. (173). The integral therefore satisfies the structure of Eqs. (188) and (189) if and only if it exhibits asymmetric behaviour between matter and antimatter—enhanced for matter and suppressed for antimatter, i.e. $\mathbf{I}_{\text{matter}} > \mathbf{I}_{\text{antimatter}}$. Such asymmetry enhances the tunnelling probability for antimatter while suppressing it for matter, enforcing matter confinement.

This behaviour has two fundamental interpretations. The first involves the covariance contribution to $\langle \hat{C}^2 \hat{j}_{tt} \rangle = \langle \hat{C}^2 \rangle \langle \hat{j}_{tt} \rangle + \text{Cov}(\hat{C}^2, \hat{j}_{tt})$; any marginal difference in covariance between matter and antimatter induces asymmetry. The second relies on the spacetime operator quotient, which via Eq. (177) becomes an extreme-field phase $\varphi(t)$ integrated over time, $\mathbf{I} = \Im \int \varphi(t) dt$. Here, a similarly small asymmetry in the extreme-field phase yields the effect. In either case, a slight enhancement of the integral for matter relative to antimatter leads to matter confinement and antimatter deconfinement.

This mechanism, which inherently favours antimatter tunnelling (deconfinement), raises a natural question: if antimatter indeed tunnelled out with a greater primordial probability, where did it ultimately go?

The answer lies in the fact that antimatter primarily annihilated with the matter component of standard Hawking radiation. The excess antimatter, which could not annihilate immediately, propagated toward the primordial universe’s horizon and eventually resided at the boundary of spacetime. Following the Jerk Tremendous—which triggered an enormous inflationary phase in the early universe—most of this antimatter underwent significant spacelike separation from the newly formed universe. This inflationary expansion propelled antimatter beyond the causal reach of the observable universe, placing it outside its light cone or horizon—permanently inaccessible to any future interaction.

Of the antimatter that remained within the universe—a small fraction compared with the created matter—its contribution became asymmetric: this fraction underwent partial annihilation with a corresponding matter component, while the residual matter seeded the universe, thereby giving rise to the primordial matter–antimatter asymmetry.

A profound and testable prediction of the theory is that the asymmetry was not lost but spatially separated: antimatter resides predominantly in regions beyond the observable horizon, forming domains causally disconnected from our matter-dominated universe.

3.7 The Planckian Extremisation Threshold

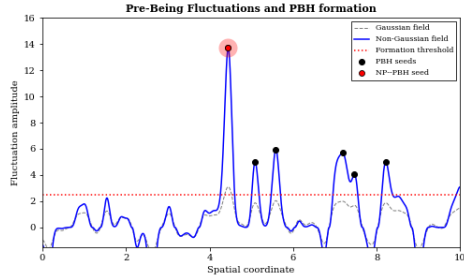
Proposition 10. *In the Pre-Being epoch, if the quantum fluctuations follow a non-Gaussian distribution [23, 32, 33, 80, 113] or other unknown nonlinear distributions, there exist locations where they are naturally amplified into acausal field excitations [38]. These excitations are hypothesised to correspond to perturbative primordial black holes. Locations of maximal amplification act as seeds for primordial black holes, defining the minimal threshold for non-perturbative quantisation. A sufficient condition is that a maximally entangled, non-perturbative primordial black hole saturates this minimal threshold, thereby setting the initial condition for the emergence of the Universe.*

Proposition 11. *The initial cosmological conditions are defined by a non-perturbative primordial black hole evolving along a maximally ordered (negentropic) flow [67, 125], contracted to the Planckian threshold. This establishes a net primordial entropy configuration corresponding to $\Delta S \ll 0$. At this critical juncture, the Planckian extremal negentropic black hole undergoes a tremendous jerk-like recoil process, irreversibly transitioning to seed large-scale cosmic structure.*

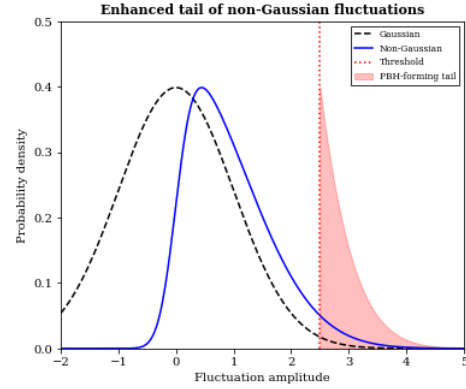
In this subsection we justify the transition from the non-perturbative primordial black hole (NP–PBH) to the primordial Planck black hole (PPBH). As shown in Fig. 3a, the NP–PBH represents the highest non-Gaussian peak—exceeding the statistical threshold (Fig. 3b) and thereby marking the onset of non-perturbative quantisation. A negentropic transition leads to the PPBH, which saturates the upper boundary of this regime at the Planckian extremisation threshold. This progression is motivated in Section 3, theorised in Theorem 2, and formally proposed in Propositions 10 and 11.

Within the non-perturbative framework of Subsection 3.3, we demonstrated that the creation rate of eigen-Hawking radiation exceeds that of standard Hawking radiation from the exterior. In this mechanism, the tunnelling particles forming the anti-eigen flux are negative-energy excitations, as established in Subsection 3.5, while the inward eigen-Hawking flux consists of positive-energy particles that seed the NP–PBH core.

Consequently, NP–PBH thermodynamics is governed by two distinct degrees of freedom: positive-energy eigen-Hawking radiation (EHR) and negative-energy ingoing particles—the negative-Hawking flux (NHF)—which counterbalances standard Hawking radiation. This duality produces a distinctive thermodynamic signature, particularly affecting primordial universe entropy.



(a) Snapshot of primordial fluctuations. The non-Gaussian field (blue) exhibits enhanced peaks compared to a Gaussian realisation (black dashed). Peaks crossing the red dotted threshold (ζ_{th}) become seeds for primordial black holes (PBHs), marked by black circles. The highest peak (red dot with black edge, highlighted by a glow) corresponds to the *non-perturbative PBH seed* (NP-PBH), which saturates the lower boundary of non-perturbative quantisation and provides the initial condition for cosmic emergence.



(b) Probability density functions of fluctuation amplitudes. The Gaussian distribution (black dashed) has an exponentially suppressed tail, while the non-Gaussian distribution (blue solid) exhibits a significantly enhanced tail (red shaded area) above the formation threshold ζ_{th} . This enhancement dramatically increases the likelihood of large fluctuations that collapse into PBHs, with the rarest events reaching the NP-PBH threshold.

Figure 3: Illustration of non-Gaussian fluctuations and primordial black hole seeding, corresponding to Proposition 10. The left panel shows a spatial realisation of the fluctuation field; the right panel demonstrates the statistical enhancement of the tail. Together, these panels illustrate how non-linearities generate PBHs and, in the most extreme case, a single NP-PBH that establishes the initial cosmological conditions described in Proposition 11.

In the standard picture of Hawking radiation, the emission of quanta from a black hole is accompanied by a decrease in the black-hole entropy, associated with the ingoing negative-energy Hawking flux across the horizon. To quantify this entropy change, we invoke the Bekenstein–Hawking relation [20, 73], which relates the black-hole entropy to the area of its event horizon:

$$S_{\text{BH}} = \frac{k_B A}{4\ell_P^2}, \quad \ell_P = \sqrt{\frac{\hbar G}{c^3}}, \quad (191)$$

where k_B is the Boltzmann constant, A denotes the horizon area, and ℓ_P is the Planck length.

For a Schwarzschild black hole, the horizon radius is given by $r_s = 2GM/c^2$, which yields a horizon area of

$$A = 4\pi r_s^2 = \frac{16\pi G^2 M^2}{c^4}. \quad (192)$$

Substituting (192) into (191), the entropy may be expressed explicitly as a function of the black-hole mass:

$$S_{\text{BH}}(M) = \frac{4\pi G M^2 k_B}{\hbar c}. \quad (193)$$

The first law of black-hole thermodynamics [12, 163],

$$dM = T_H dS, \quad (194)$$

provides the fundamental relation between mass and entropy. Multiplying the left-hand side by c^2 to restore energy units and applying the Bekenstein–Hawking entropy formula, Eq. (193), then yields the Hawking temperature,

$$T_H = \left(\frac{dS}{d(Mc^2)} \right)^{-1} = \frac{\hbar c^3}{8\pi G M k_B}. \quad (195)$$

Emission of a quantum of energy E reduces the black-hole mass from M to $M - E/c^2$. In the fixed-background (purely thermal) approximation, the corresponding entropy change follows from expanding Eq. (193) to first order in E :

$$\begin{aligned} \Delta S &\equiv S_{\text{BH}}\left(M - \frac{E}{c^2}\right) - S_{\text{BH}}(M) \\ &\simeq -\frac{dS}{dM} \frac{E}{c^2} = -\frac{8\pi G M E k_B}{\hbar c^3}. \end{aligned} \quad (196)$$

Within the semiclassical tunnelling framework, the emission probability is determined by the change in black-hole entropy via the detailed-balance relation [123]

$$\Gamma \sim \exp\left(\frac{\Delta S}{k_B}\right), \quad (197)$$

which follows from the microcanonical identification of the number of black-hole states with $\exp(S_{\text{BH}}/k_B)$. Substituting (196) into (197) yields the standard thermal result

$$\Gamma \sim \exp\left(-\frac{8\pi G M E}{\hbar c^3}\right), \quad (198)$$

equivalent to the Boltzmann factor $\exp(-E/k_B T_H)$.

To incorporate black-hole back-reaction consistently, as established in Eq. (112), energy conservation must be enforced during the emission process. This can be achieved by evaluating the change in black-hole entropy at the reduced mass, using Eq. (193):

$$\begin{aligned} \Delta S &= S_{\text{BH}}\left(M - \frac{E}{c^2}\right) - S_{\text{BH}}(M) = \frac{4\pi G k_B}{\hbar c} \left[\left(M - \frac{E}{c^2}\right)^2 - M^2 \right] \quad (199) \\ &\approx \frac{4\pi G k_B}{\hbar c} \left[M^2 - 2M \frac{E}{c^2} + \frac{E^2}{c^4} - M^2 \right] = -\frac{8\pi G E k_B}{\hbar c^3} \left(M - \frac{E}{2c^2}\right). \quad (200) \end{aligned}$$

Following the same reasoning as in Eq. (112), one can define

$$\Delta^- \equiv M - \frac{E}{2c^2}. \quad (201)$$

Evaluating the entropy change in this way, rather than expanding about a fixed background, naturally incorporates the leading-order effect of back-reaction. The corresponding tunnelling probability becomes

$$\Gamma_{\text{SHR}}^{\text{BR}} \sim \exp\left[\frac{-8\pi G E \Delta^-}{\hbar c^3}\right], \quad (202)$$

which is in agreement with the Parikh–Wilczek tunnelling formulation of Hawking radiation [121], as derived in Eq. (112).

Accordingly, the back-reaction gives rise to a negative entropy change, originating from the negative-Hawking flux (NHF), and can be expressed as:

$$\Delta S_{\text{NHF}}^- = -\frac{8\pi G E k_B \Delta^-}{\hbar c^3}, \quad (203)$$

demonstrating that back-reaction induces a deviation from strict thermality while remaining consistent with semiclassical black-hole thermodynamics.

Analogously, we examine the entropy change associated with positive-energy eigen-Hawking radiation (EHR). The corresponding anti-eigen flux obeys the back-reaction-corrected tunnelling probability given by Eq. (151):

$$\Gamma_{\text{AEF}}^{\text{BR}} \sim \exp\left[-\frac{8\pi GE}{\hbar c_*^3(1 + \Upsilon_{\text{NP}})^2 \varrho} \left(M + \frac{E}{2c_e^2}\right)\right] = \exp\left[\frac{-8\pi GE \Delta^+}{\hbar c_*^3(1 + \Upsilon_{\text{NP}})^2 \varrho}\right]. \quad (204)$$

For the positive-energy eigen-Hawking radiation emitted inward, the associated black-hole entropy must increase, in accordance with the second law of thermodynamics [43]. This thermodynamic requirement imposes a detailed-balance condition on the tunnelling probability. However, when the entropy change is positive, the standard detailed-balance relation in Eq. (197) would lead to an unbounded exponential enhancement of the probability—thereby violating probabilistic consistency.

To preserve probability normalisation while maintaining thermodynamic consistency, the detailed-balance relation must therefore be generalised and expressed in terms of the negative modulus of the entropy change:

$$\Gamma \sim \exp\left(-\frac{|\Delta S|}{k_B}\right). \quad (205)$$

Equating the tunnelling probability in Eq. (204) with the generalised detailed-balance form (205) yields the positive entropy change associated with eigen-Hawking radiation,

$$\Delta S_{\text{EHR}}^+ = \frac{8\pi GE k_B \Delta^+}{\hbar c_*^3(1 + \Upsilon_{\text{NP}})^2 \varrho}. \quad (206)$$

From the intrinsic spacetime asymmetry illustrated in Fig. 1—discussed further in the non-perturbative framework of Subsection 3.3 and encoded mathematically by the non-perturbative, extreme-field inequality of Eq. (169)—we obtain a corresponding inequality between the entropic variations of Eqs. (203) and (206). This inequality reads

$$|\Delta S_{\text{NHF}}^-| > |\Delta S_{\text{EHR}}^+|, \quad (207)$$

where $\Delta S_{\text{NHF}}^- < 0$ is the negentropic change induced by the negative-Hawking flux, and $\Delta S_{\text{EHR}}^+ > 0$ is the positive entropy contribution from the eigen-Hawking radiation.

For the non-perturbative primordial black hole (NP–PBH), the total entropy change is therefore negative:

$$\Delta S_{\text{NP-PBH}} = \Delta S_{\text{NHF}}^- + \Delta S_{\text{EHR}}^+ < 0. \quad (208)$$

Consequently, the final state of the NP–PBH possesses lower entropy than its initial state. We term this phenomenon *non-perturbative entropy contraction*.

The result appears counter-intuitive at first glance. According to Eq. (170), the creation rate of eigen-Hawking radiation exceeds that of standard Hawking radiation (and hence also exceeds the rate of the compensating negative Hawking flux). Naively, one would expect that adding more positive-energy states should increase the entropy. However, in the non-perturbative extreme-field regime, the creation-rate inequality of Eq. (170)—which defines the intrinsic spacetime asymmetry—is realised *if and only if* the fundamental asymmetry condition (169) holds. This condition demands that the effective exponent in the tunnelling probability for the anti-eigen flux be less negative than that for standard Hawking radiation, after accounting for the modified speed of light c_* , the quantum corrections Υ_{NP} , and the back-reaction term Δ^+ . Since the exponent is proportional to the entropy change (up to the Boltzmann constant), as shown in Eq. (205), the less negative exponent directly translates into a smaller magnitude of negentropy, thereby explaining the net entropy contraction relative to the standard Hawking process.

A deeper understanding follows from the analysis in Subsection 3.5. There it is shown that the negentropic change due to the negative-Hawking flux equals the negative of the Bekenstein tunnelling exponent (up to a factor of k_B), while the positive entropic change from the eigen-Hawking radiation corresponds to the positive exponent multiplied by a contraction factor ω . Crucially, within the non-perturbative framework, ω —explicitly given by Eq. (178)—acts as a suppression parameter that reduces the positive entropy contribution. Thus, although a net addition of positive-energy states occurs, extreme-field physics enforces an overall entropy contraction, maintaining the NP–PBH in a more ordered, negentropic state.

This entropy-contraction mechanism represents a distinctive thermodynamic signature of the non-perturbative regime and provides a rigorous basis for the transition to a primordial Planck black hole. The analysis thereby justifies Proposition 11: in light of Eq. (191), the total net negentropic change induces an area contraction, forcing the NP–PBH to follow a negentropic flow until it reaches a Planckian threshold where its area approaches $\sim \ell_P^2$.

Consequently, the cumulative area contraction is parametrically large:

$$|\Delta A| \gg \ell_P^2, \quad (209)$$

signifying a macroscopic reduction of the horizon.

At this juncture, the net negentropic state of the threshold object—now identified as the primordial Planck black hole (PPBH)—implies, via Eq. (191), a large negative entropy:

$$S_{\text{PPBH}} \ll 0. \quad (210)$$

Substituting Eq. (210) into the generalised detailed-balance relation (205) yields an extreme suppression of the tunnelling probability for both the anti-eigen flux and standard Hawking radiation:

$$\Gamma_{\text{PPBH}} \sim 0. \quad (211)$$

We identify this condition as Planckian extremisation: the system simultaneously saturates the fundamental upper bound of causal horizon fluctuations and maximises interior acausal field excitations. In this saturated regime, quantum fluctuations cease oscillatory behaviour and converge to a unified Planckian state. The resulting primordial Planck black hole (PPBH) thus constitutes a terminal gravitational configuration, where further Hawking emission is dynamically prohibited. Consequently, this extreme negentropic state—corresponding to zero accessible microstates and extraordinary sensitivity to perturbations—represents the upper limit of the non-perturbative regime. It is rendered highly unstable by an intrinsic non-perturbative eigen back-reaction, which is justified in Subsection 4.3 (Theorem 3) and physically realised as an effective infinite resistance that prevents the extreme Planckian configuration from collapsing indefinitely to a singularity.

Having established the terminal, extremised nature of the primordial Planck black hole (PPBH), this framework motivates an examination of the complete structure of the primordial universe. We propose that the primordial universe consists of three interacting components: (i) classical causal pre-being fluctuations (PBFs), which generate a positive entropic contribution (Subsection 3.1); (ii) perturbative primordial black holes with approximately fixed entropy (Subsection 3.2); and (iii) the non-perturbative fluxes—the negentropic anti-eigen flux and the entropic standard Hawking radiation (Subsection 3.3)—both of which originate from the non-perturbative primordial black hole as it approaches the PPBH configuration.

Recalling the hierarchy established in Eq. (170), the relative flux rates can be expressed as

$$F_{\text{AEF}} > F_{\text{SHR}}. \quad (212)$$

Eq. (212) indicates that, in the primordial universe, the rate of the negative-energy anti-eigen flux exceeds that of the positive-energy standard Hawking radiation as measured by an observer at infinity. Consequently, the tunnelling of negative-energy particles associated with the anti-eigen flux contributes a net negative entropy,

$$\Delta S_{\text{AEF}} < 0, \quad (213)$$

whereas the same observer records standard Hawking radiation as contributing a positive entropy,

$$\Delta S_{\text{SHR}} > 0. \quad (214)$$

The non-perturbative extreme-field asymmetry between the anti-eigen and standard Hawking flux rates—captured in Eqs. (169) and (170), and equivalently in the flux notation of Eq. (212)—implies that the negentropic contribution generated by the anti-eigen flux progressively dominates over the combined positive entropic contributions from standard Hawking radiation and from classical pre-being fluctuations. This leads to a net entropy imbalance,

$$|\Delta S_{\text{AEF}}| > |\Delta S_{\text{SHR}}| + |\Delta S_{\text{PBFs}}|. \quad (215)$$

Physically, Eq. (215) implies that the anti-eigen flux progressively suppresses the random microstates generated by standard Hawking radiation and causal pre-being fluctuations, drawing the primordial energy distribution into an increasingly uniform configuration within the PPBH. This process effectively reverses the conventional thermodynamic arrow—a phenomenon we term the *time–entropy arrow reversal*. Consequently, from the perspective of an observer at infinity, the primordial universe does not appear to expand outwardly, but rather evolves inwardly towards an increasingly ordered, low-entropy state.

The central result emerging from Eqs. (208) and (215) is that, during the pre-being epoch, the primordial universe undergoes a continuous net decrease in total entropy. This behaviour reflects the combined action of time–entropy arrow reversal at the cosmological level and entropy contraction within the non-perturbative primordial black hole. Consequently, the energy content of the universe is progressively drawn toward a unique Planckian core, forming

an exceptionally ordered macrostate identified as the PPBH. The net entropic evolution of the primordial universe may therefore be expressed as

$$\Delta S_{\text{universe}} = \Delta S_{\text{PPBH}} + \Delta S_{\text{AEF}} + \Delta S_{\text{SHR}} + \Delta S_{\text{PBFs}} < 0. \quad (216)$$

In this extreme limit, the PPBH functions as the dominant universal attractor, drawing in all available energy. Analogous to a singular point in configuration space, it absorbs highly random, non-uniform primordial microstates and transforms them into a coherent, uniform macrostate. Consequently, the total entropy of the primordial universe is reduced, culminating in a highly negentropic state physically embodied by the PPBH. This extreme ordering can be quantified via Boltzmann’s entropy relation [29, 129]:

$$S = k_B \ln W, \quad (217)$$

where W denotes the number of accessible microstates.

In conventional physics, a state with vanishing microstates is regarded as unattainable, as residual energy and disorder persist even at absolute zero. However, within the non-perturbative extreme-field regime, the intrinsic spacetime asymmetry—quantified in Eqs. (169) and (170)—enforces an effective reversal of the entropy arrow, driving the system toward progressively lower randomness and higher order. In this regime, the limit of vanishing microstates becomes physically accessible, yielding a primordial universe that is intrinsically negentropic. Specifically, as $W \rightarrow 0$,

$$\lim_{W \rightarrow 0} \Delta S_{\text{universe}} = \lim_{W \rightarrow 0} k_B \ln W = -\infty. \quad (218)$$

At this juncture of extreme order—where discretisation saturates to a continuum, gravitational coupling becomes indistinguishable from quantum coupling, and the resulting deterministic state is rendered inherently unstable—such a continuum deterministic configuration is prohibited at the Planckian threshold. Consequently, a non-perturbative eigen back-reaction of the field on itself becomes inevitable, as established in Theorem 3, driving an extremisation phase transition that forces the system from the Planckian threshold into an entropic large-scale regime where gravitational coupling becomes asymptotically free. The detailed dynamics of this extremisation phase transition—the physics of the quantum-gravitational jerk bridging the micro-Planckian and macro-cosmological regimes—are derived in Section 4.

4 Main

4.1 The Cosmological Initial Asymmetries

Symmetries in physics represent fundamental invariances in nature’s laws, while their breaking—where systems governed by symmetric laws transition to asymmetric states—has profound implications across physics. This is evidenced by chiral symmetry breaking in quantum chromodynamics (QCD) [115], CP violation in weak interactions [144], and electroweak symmetry breaking via the Higgs mechanism [78]. In cosmology, early-universe phase transitions accompanied by symmetry breaking shaped matter–antimatter asymmetry and cosmic structure formation [91]. Building on this foundation, the theory predicts that intrinsically asymmetric physics governed the Pre-Being Epoch, seeding the processes that culminated in cosmic genesis.

A primary breaking event in the primordial universe was the violation of causality [3, 58, 98, 111], which ordinarily ensures that effects follow their causes. In the Pre-Being Epoch, however, spacetime—the necessary substrate for fluctuations—could not emerge causally. The formation of primordial black holes (PPBs), with non-causal core fields exhibiting superluminal fluctuations, therefore preceded their causal triggers: the pre-being fluctuations (PBFs) near the nascent horizons. This counterintuitive ordering can be understood through gravitational time dilation, $dt = d\tau/(-g_{tt})$, as measured by a higher observer [132, 155]. For such an observer, the emergence of the PPBs occurs before the PBFs, i.e., $dt_{\text{PPB}} < dt_{\text{PBF}}$, because the temporal component characterising the perturbative PPB core, $(-g_{tt}^p)$ (Eq. (155)), marginally exceeds the standard Schwarzschild metric component, g_{tt} , which characterises the causal fluctuations (Eq. (91)). Consequently, this apparent violation of causality was not merely allowed but fundamentally required for the formation of a spacetime substrate capable of supporting subsequent physical processes. This effect is intimately tied to the intrinsic spacetime asymmetry described in Section 3: inside the non-perturbative primordial black hole (NP–PBH) horizon, spacetime is quantised non-perturbatively, while the exterior remains a classical continuum. This asymmetry enables the eigen-Hawking radiation (EHR) rate inside the NP–PBH to exceed that of standard Hawking radiation (SHR) outside (Eq. (170))—a result encoded in Eq. (169). Its solution confirms the extreme-field parameters remain bounded but not unity (Eqs. (177) and (178)), establishing initial conditions for an apparent thermodynamic violation.

The matter–antimatter asymmetry originates from further asymmetries rooted in the non-perturbative quantisation framework, as detailed in Subsection 3.6. Two distinct mechanisms emerge from this framework: (i) PT-symmetric evolution [22], which permits negative-energy backward-evolving modes to tunnel outward as anti-eigen flux (antimatter), while confining positive-energy forward-evolving eigen-Hawking modes (matter) inward; and (ii) a more general extreme-field solution, in which the asymmetry parameter ω enhances matter confinement while marginally suppressing antimatter deconfinement. Both mechanisms naturally produce a matter-dominated universe consistent with the observed baryon asymmetry [144].

Quantum vacuum fluctuations normally exhibit translational symmetry—they are statistically uniform across space. During the Pre-Being Epoch, however, fluctuations displayed significant spatial heterogeneity following a nonlinear, non-Gaussian distribution [46, 94, 103, 160]. Fluctuations were most intense near perturbative primordial black-hole horizons, weaker in distant regions, and extreme within the non-perturbative primordial black hole—the seed that saturated the lower boundary of non-perturbative quantisation and initiated the negentropic flow toward the Primordial Planck Black Hole (PPBH).

4.2 The Quantum-Gravitational Conjugate

Section 1 establishes the speed of light as an operator in extreme-field physics, where its classical constancy is modified by a significant quantum-gravitational conjugate, $\hat{\chi}$, in the non-perturbative formulation of Section 3.3. By Proposition 2 (Eq. (42)) and Eq. (87), this operator takes the form

$$\hat{\mathcal{C}} = \sqrt{c^2 + \hat{Z}^\mu \nabla_\mu}, \quad (219)$$

where ∇_μ denotes the spacetime covariant derivative, whose action on a test function is given explicitly by

$$\nabla_\mu \phi^\nu = \partial_\mu \phi^\nu + \Gamma_{\mu\lambda}^\nu \phi^\lambda, \quad (220)$$

with $\Gamma_{\mu\lambda}^\nu$ the Christoffel symbols, and \hat{Z}^μ the quantum operator encoding non-perturbative gravitational fluctuations.

Historically, the invariance of c follows from the electromagnetic properties of the vacuum, specifically the permittivity ϵ_0 and permeability μ_0 of free space [109]:

$$c = \frac{1}{\sqrt{\epsilon_0 \mu_0}}. \quad (221)$$

These constants characterise the vacuum impedance to electric and magnetic fields; their constancy follows from the translational symmetry, isotropy, and homogeneity of flat spacetime. Combining them into a single homogeneous causality constant χ_0 and setting $c = \sqrt{\chi_0}$ yields, via Eq. (221), the definition

$$\chi_0 \equiv \frac{1}{\epsilon_0 \mu_0}. \quad (222)$$

In extreme electromagnetic field regimes, ϵ_0 and μ_0 are rendered dynamical [2, 5, 59, 105]. This breaks translational symmetry, forcing χ_0 to acquire an operator character. Strong fields polarise the vacuum, varying ϵ_0, μ_0 ; similarly, strong curvature modifies light propagation [84]. These considerations motivate promoting the speed of light to an operator. Such an operator description offers key advantages over a classical variable- c model: it incorporates quantum fluctuations and couples consistently to quantum fields and curvature, aligning with quantum-gravity expectations. We thus promote $\chi_0 = c^2$ to an extreme-field operator $\hat{\chi}^* = \hat{\mathcal{C}}^2$, defined as

$$\hat{\mathcal{C}}^2 \equiv c^2 + \hat{\chi}. \quad (223)$$

However, classical physics must be recovered at the boundary between the extreme and classical regimes, implying the condition

$$\lim_{\text{boundary}} \hat{\chi} = 0, \quad \lim_{\text{boundary}} \hat{\mathcal{C}} = c. \quad (224)$$

For a black hole, this boundary is the event horizon. The Schwarzschild radius, $r_s = 2GM/c^2$, remains fixed because the speed of light returns to its classical value there. Consequently, quantum-gravitational states $|\phi\rangle$ must be annihilated by the extreme-field operator at the horizon:

$$\lim_{r \rightarrow r_s} \hat{\chi}|\phi\rangle = 0. \quad (225)$$

Within black-hole interiors, locations where spacetime deviates from classical smoothness and becomes purely fluctuating yield a distinct geometry—the Extreme Hermitian Geometry—that establishes the extreme conditions for Hermiticity. There, the vanishing boundary condition together with $\hat{\chi} = \hat{Z}^\mu \nabla_\mu$ (with $\hat{Z}^\mu = i\hbar\alpha \partial^\mu$ from Eq. (245)) guarantees conditional Hermiticity iff $\hbar\alpha \partial^\nu \nabla_\nu$ is anti-Hermitian. This occurs in two ways:

1. $\partial^\nu \nabla_\nu$ itself satisfies anti-Hermiticity under a conditional anticommutative geometry;

2. the combination $\alpha\partial^\nu\nabla_\nu$ yields an odd number of integrations by parts, rendering it anti-Hermitian.

The second possibility requires interpreting the running coupling as the operator $\alpha = \alpha_0\partial_u$ (Proposition 2, Eq. (45)). Building on this, we then examine the compound action of $\alpha\partial^\mu\nabla_\mu$ on states in the Hilbert space equipped with the inner product

$$\langle\phi|\psi\rangle = \int d^4x\sqrt{-g}\phi^*\psi. \quad (226)$$

Applying integration by parts, imposing vanishing boundary conditions on the fields and their derivatives, and using the noncommutativity identity $CBA = ABC - A[B, C] - [A, C]B - C[A, B]$ yields

$$\begin{aligned} \alpha_0\langle\phi|\partial_u\partial^\mu\nabla_\mu\psi\rangle &= -\alpha_0\langle\nabla_\mu\partial^\mu\partial_u\phi|\psi\rangle \\ &= -\alpha_0\langle\partial_u\partial^\mu\nabla_\mu\phi|\psi\rangle \end{aligned} \quad (227)$$

$$+ \alpha_0\left\langle\left(\partial_u[\partial^\mu, \nabla_\mu] + [\partial_u, \nabla_\mu]\partial^\mu + \nabla_\mu[\partial_u, \partial^\mu]\right)\phi\right|\psi\rangle. \quad (228)$$

Eq. (227) therefore admits conditional anti-Hermiticity at particular extreme-field locations, which we refer to as points of Hermiticity, where the compound commutator terms in Eq. (228) vanish. Acting on a test function, this configuration defines the commutative extreme Hermitian geometry identity:

$$\left(\left(\partial_u[\partial^\mu, \nabla_\mu] + [\partial_u, \nabla_\mu]\partial^\mu + \nabla_\mu[\partial_u, \partial^\mu]\right)\Phi\right)^\sigma = 0. \quad (229)$$

This condition, however, is complex and may conflict with Eq. (242) if the emerging coupling is a genuine coordinate operator. An alternative is to pursue the first possibility: that the coupling is strictly coordinate-independent and that $\partial^\nu\nabla_\nu$ satisfies anti-Hermiticity under a conditional anticommutative geometry, expressed as

$$\langle\phi|\partial^\mu\nabla_\mu\psi\rangle = \langle\nabla_\mu\partial^\mu\phi|\psi\rangle = -\langle\partial^\mu\nabla_\mu\phi|\psi\rangle \quad (230)$$

$$+ \langle\{\nabla_\mu, \partial^\mu\}\phi|\psi\rangle. \quad (231)$$

Using Eq. (47) together with the covariant derivative acting on a test function (Eq. (220)), one obtains the explicit form of $\{\nabla_\mu, \partial^\mu\}$ appearing in Eq. (231):

$$\begin{aligned} \left(\{\nabla_\nu, \partial^\nu\}\Phi\right)^\sigma &= \left((\partial_\nu g^{\nu\rho})\partial_\rho\Phi^\sigma + \partial^\nu\partial_\nu\Phi^\sigma + \Gamma_{\nu\rho}^\sigma\partial^\nu\Phi^\rho\right) \\ &\quad + \left(\partial^\nu\partial_\nu\Phi^\sigma + (\partial^\nu\Gamma_{\nu\lambda}^\sigma)\Phi^\lambda + \Gamma_{\nu\lambda}^\sigma\partial^\nu\Phi^\lambda\right) \end{aligned} \quad (232)$$

$$= (\partial_\nu g^{\nu\rho})\partial_\rho\Phi^\sigma + 2\partial^\nu\partial_\nu\Phi^\sigma + 2\Gamma_{\nu\lambda}^\sigma\partial^\nu\Phi^\lambda + (\partial^\nu\Gamma_{\nu\lambda}^\sigma)\Phi^\lambda. \quad (233)$$

Eq. (230) therefore yields conditional anti-Hermiticity upon vanishing of the anticommutator term in Eq. (231). Using the explicit form in Eq. (233), this defines the anticommutative extreme Hermitian geometry identity:

$$(\partial_\nu g^{\nu\rho})\partial_\rho\Phi^\sigma + 2\partial^\nu\partial_\nu\Phi^\sigma + 2\Gamma_{\nu\lambda}^\sigma\partial^\nu\Phi^\lambda + (\partial^\nu\Gamma_{\nu\lambda}^\sigma)\Phi^\lambda = 0. \quad (234)$$

When either condition—Eq. (229) or Eq. (234)—holds, $\hbar\alpha\partial^\mu\nabla_\mu$ becomes anti-Hermitian:

$$(\hbar\alpha\partial^\mu\nabla_\mu)^\dagger = -\hbar\alpha\partial^\mu\nabla_\mu. \quad (235)$$

Consequently, the operator $\hat{\chi} = i\hbar\alpha\partial^\mu\nabla_\mu$ is Hermitian, the factor i ensuring

$$\hat{\chi}^\dagger = (i\hbar\alpha\partial^\mu\nabla_\mu)^\dagger = -i\hbar(\alpha\partial^\mu\nabla_\mu)^\dagger = -i\hbar(-\alpha\partial^\mu\nabla_\mu) = i\hbar\alpha\partial^\mu\nabla_\mu = \hat{\chi}, \quad (236)$$

yielding real eigenvalues and expectation values as required for an observable.

Should both conditions fail, we instead employ the antisymmetric construction (Remark 2, Eq. (44)), which yields a globally Hermitian operator without conditional assumptions. This construction, together with the Extreme Hermitian Geometry, will be addressed in a forthcoming paper.

Turning to the perspective of a quantum-gravitational state, the vanishing boundary condition (Eq. (225)) can be reinterpreted as the freezing of its quantum-gravitational evolution, which implies

$$\lim_{r \rightarrow r_s} \nabla_\mu |\phi\rangle = 0. \quad (237)$$

This observation validates the structure of $\hat{\chi}$ given in Eq. (42) (Proposition 2):

$$\hat{\chi} \equiv \hat{Z}^\mu \nabla_\mu. \quad (238)$$

From a deeper physical perspective, Eq. (225), together with the definition of $\hat{\chi}$ in Eq. (238), demands reinterpretation as a parallel-transport condition along an operator-defined quantum direction. At the horizon, the state $|\phi\rangle$ becomes covariantly constant along \hat{Z}^μ , signalling the freezing of quantum-gravitational evolution. The horizon thereby acts as a natural filter: the state freezes, classical degrees of freedom propagate outward, and probabilistic quantum features are suppressed. Conversely, within the black-hole interior, the non-vanishing condition forces non-parallel transport along a closed loop of the quantised interior—a behaviour mirroring the Berry geometric phase in adiabatic quantum cycles [24] and thereby encoding the geometric structure of the interior.

The resulting operator $\hat{\mathcal{C}}$ is inherently non-local—a feature that is not an artefact but a necessity for resolving singularities and incorporating space-time entanglement and holography [86,152]. In quantum gravity, such non-locality can transiently modify classical causality, as anticipated in the Pre-Being Paradox introduced in Section 2: there, the final state—primordial black holes in an extreme non-causal field—precedes its cause, namely the causal pre-being fluctuations. This reversed chronology, rooted in the non-local dynamics of the quantised speed of light, underscores the consistency of the operator framework in Eq. (219) for extreme-field physics.

4.3 Quantum-Gravitational Noncommutative Algebra and the Emergence of a Unified Coupling

A consistent quantum-gravitational field theory requires spacetime quantisation [48, 101, 141]. In extreme field theory, this emerges from promoting the speed of light to a quantum operator—an additional term encoding quantum-gravitational effects while preserving classical coordinate labels.

Principle 1 promotes the metric to a quantum operator: the line element becomes operator-valued while coordinates remain classical. Spatial coordinates may later acquire operator character, but time stays a classical external parameter. This structure induces noncommutativity between spacetime coordinates and the quantum-gravitational conjugate.

To formalise this structure, we examine the fundamental commutation relations between the spacetime coordinate operators \hat{x}^μ [116] and the quantum-gravitational conjugate $\hat{Z}^\nu \nabla_\nu$, which arises naturally from quantising the metric in the extreme-field limit, as demonstrated in Proposition 6:

$$[\hat{x}^\mu, \hat{Z}^\nu \nabla_\nu]. \quad (239)$$

We begin with metric-compatible geometry, where the classical Riemannian framework [97,138] may admit quantum modifications. Covariant derivative preserves the metric, $\nabla_\lambda g_{\mu\nu} = 0$, so indices are raised and lowered by $V_\mu = g_{\mu\nu} V^\nu$ and $V^\mu = g^{\mu\nu} V_\nu$. Metric symmetry $g_{\mu\nu} = g_{\nu\mu}$ guarantees invariance under parallel transport: $\nabla_\lambda (g_{\mu\nu} V^\mu W^\nu) = 0$. Christoffel symbols are symmetric, $\Gamma_{\mu\nu}^\lambda = \Gamma_{\nu\mu}^\lambda$. Within this framework, Eq. (239) expands the commutation relation between spacetime and the quantum-gravitational conjugate using the distributive property, yielding:

$$[\hat{x}^\mu, \hat{Z}^\nu \nabla_\nu] = [\hat{x}^\mu, \hat{Z}^\nu] \nabla_\nu + \hat{Z}^\nu [\hat{x}^\mu, \nabla_\nu]. \quad (240)$$

We commence by examining the first term on the right-hand side of Eq. (240):

$$[\hat{x}^\mu, \hat{Z}^\nu] \nabla_\nu. \quad (241)$$

To ensure the units are consistent, we introduce a constant α with units [Length²/(Mass \times Time)]. This gives us the commutation relation:

$$[\hat{x}^\mu, \hat{Z}^\nu] = i\hbar\alpha g^{\mu\nu}. \quad (242)$$

This directly mirrors the canonical commutation relation between the four-position operator \hat{x}^μ and the four-momentum operator \hat{p}_ν in standard quantum theory, with \hat{Z}^ν interpreted as the canonical four-momentum conjugate to \hat{x}^μ , scaled by the emergent extreme-field constant α .

By substituting Eq. (242) into Eq. (241), we obtain:

$$[\hat{x}^\mu, \hat{Z}^\nu] \nabla_\nu = i\hbar\alpha g^{\mu\nu} \nabla_\nu, \quad (243)$$

and contracting the covariant derivative with the metric tensor yields:

$$[\hat{x}^\mu, \hat{Z}^\nu] \nabla_\nu = i\hbar\alpha \nabla^\mu. \quad (244)$$

In this analysis, we define \hat{Z}^μ as:

$$\hat{Z}^\mu \equiv i\hbar\alpha \partial^\mu. \quad (245)$$

We now turn to the second commutator on the right-hand side of Eq. (240):

$$\hat{Z}^\nu [\hat{x}^\mu, \nabla_\nu] = \hat{Z}^\nu (\hat{x}^\mu \nabla_\nu - \nabla_\nu \hat{x}^\mu). \quad (246)$$

Substituting the differential form of \hat{Z}^ν from Eq. (245), together with the explicit covariant derivative, expands the operator commutator at the level of abstract algebra. This expansion draws on the canonical commutation relations of quantum mechanics and the tensorial properties of general relativity, treating the operators independently of any specific test function:

$$\hat{Z}^\nu [\hat{x}^\mu, \nabla_\nu] = i\hbar\alpha \partial^\nu (\hat{x}^\mu \nabla_\nu - \nabla_\nu \hat{x}^\mu) \quad (247)$$

$$= i\hbar\alpha (\partial^\nu \hat{x}^\mu \nabla_\nu - \partial^\nu \nabla_\nu \hat{x}^\mu) \quad (248)$$

$$= i\hbar\alpha ((\partial^\nu \hat{x}^\mu) \nabla_\nu + \hat{x}^\mu (\partial^\nu \nabla_\nu) - (\partial^\nu \nabla_\nu) \hat{x}^\mu - \nabla_\nu (\partial^\nu \hat{x}^\mu)) \quad (249)$$

$$= i\hbar\alpha (g^{\nu\mu} \nabla_\nu + \hat{x}^\mu (\partial^\nu \nabla_\nu) - (\partial^\nu \nabla_\nu) \hat{x}^\mu - \nabla_\nu g^{\nu\mu}) \quad (250)$$

$$= i\hbar\alpha (\nabla^\mu + \hat{x}^\mu (\partial^\nu \nabla_\nu) - (\partial^\nu \nabla_\nu) \hat{x}^\mu) \quad (251)$$

$$= i\hbar\alpha (\nabla^\mu + \hat{x}^\mu \partial^\nu \partial_\nu + \hat{x}^\lambda \partial^\nu \Gamma_{\nu\lambda}^\mu - \partial^\nu \partial_\nu \hat{x}^\mu - (\partial^\nu \Gamma_{\nu\lambda}^\mu) \hat{x}^\lambda) \quad (252)$$

$$= i\hbar\alpha (\nabla^\mu + \hat{x}^\mu \hat{\square} + \hat{x}^\lambda (\partial^\nu \Gamma_{\nu\lambda}^\mu) - \hat{\square} \hat{x}^\mu - (\partial^\nu \Gamma_{\nu\lambda}^\mu) \hat{x}^\lambda) \quad (253)$$

$$= i\hbar\alpha (\nabla^\mu + [\hat{x}^\mu, \hat{\square}] + [\hat{x}^\lambda, \partial^\nu \Gamma_{\nu\lambda}^\mu]). \quad (254)$$

Under metric compatibility, expanding Eqs. (249)–(251) via the product rule isolates the relevant contribution. Eqs. (252)–(253) introduce Christoffel symbols and their derivatives, allowing the d’Alembertian $\hat{\square} = \partial^\nu \partial_\nu$ to be identified, thereby yielding a non-trivial quantum geometry. Eq. (254) then combines covariant derivatives and commutators, identifying the reciprocal-space operator $\hat{\mathcal{K}}^\mu$ and the spacetime-geometry operator $\hat{\mathcal{G}}^\mu$ as follows:

$$\hat{\mathcal{K}}^\mu \equiv [\hat{x}^\mu, \hat{\square}], \quad (255)$$

$$\hat{\mathcal{G}}^\mu \equiv [\hat{x}^\lambda, \partial^\nu \Gamma_{\nu\lambda}^\mu]. \quad (256)$$

We define the metric-compatible tensor operator

$$\Psi^\mu = 2\nabla^\mu + \hat{\mathcal{K}}^\mu + \hat{\mathcal{G}}^\mu, \quad (257)$$

which unifies the covariant derivative ∇^μ , the reciprocal-space operator $\hat{\mathcal{K}}^\mu$, and the spacetime-geometry operator $\hat{\mathcal{G}}^\mu$. In this framework, the commutator between the quantised spacetime coordinate and its quantum-gravitational conjugate takes the simplified form

$$[\hat{x}^\mu, \hat{Z}^\nu \nabla_\nu] = i\hbar\alpha\Psi^\mu. \quad (258)$$

Eq. (258) is a cornerstone of the theory, encoding noncommutative extreme-field behaviour and providing a natural quantum-gravitational regularisation mechanism that resolves black hole singularities—a sharp contrast to classical metrics like the Schwarzschild solution in Eq. (65), whose r^{-1} terms diverge as $r \rightarrow 0$, producing curvature singularities.

The quantisation of the metric tensor, established in Proposition 6, naturally gives rise to the noncommutativity in Eq. (258) when the spacetime operator is restricted to the radial coordinate. From this, the metric-compatible operator Ψ^μ emerges. Assuming Ψ^μ is bounded on the Hilbert space of quantum-gravitational states, there exists a finite constant M —its operator norm—such that

$$\|\Psi^\mu V^\sigma\| \leq M \|V^\sigma\|, \quad \forall V^\sigma. \quad (259)$$

This bound is mathematically underpinned by the Cauchy–Schwarz inequality, which guarantees that the action of Ψ^μ on any finite-norm state cannot diverge. Under this condition, the action of Ψ^μ produces a well-defined state with finite norm:

$$\Psi^\mu V^\sigma \neq 0, \quad \|\Psi^\mu V^\sigma\| < \infty. \quad (260)$$

This boundedness is crucial, as it ensures that all components of the metric remain finite even at $r = 0$, thereby providing a regular, singularity-free description of the black hole interior.

It is therefore essential to verify that the metric-compatible tensor operator Ψ^μ defined in Eq. (257) is both non-vanishing and finite. This requires analysing the combined contributions of the covariant-derivative term ∇^μ , the reciprocal-space operator $\hat{\mathcal{K}}^\mu$ [Eq. (255)], and the spacetime-geometry operator $\hat{\mathcal{G}}^\mu$ [Eq. (256)].

We first write the covariant derivative acting on a vector V^σ explicitly, including the factor of 2 appearing in Eq. (257):

$$2\nabla^\mu V^\sigma = 2g^{\mu\nu}\nabla_\nu V^\sigma = 2\partial^\mu V^\sigma + 2g^{\mu\nu}\Gamma_{\nu\lambda}^\sigma V^\lambda. \quad (261)$$

Next, the action of the reciprocal-space operator on V^σ is given by

$$\hat{\mathcal{K}}^\mu V^\sigma = [\hat{x}^\mu, \hat{\square}]V^\sigma = \hat{x}^\mu \partial^\nu \partial_\nu V^\sigma - \partial^\nu \partial_\nu (\hat{x}^\mu V^\sigma) \quad (262)$$

$$= \hat{x}^\mu \partial^\nu \partial_\nu V^\sigma - (\partial^\nu \partial_\nu \hat{x}^\mu) V^\sigma - 2(\partial^\nu \hat{x}^\mu) \partial_\nu V^\sigma - \hat{x}^\mu \partial^\nu \partial_\nu V^\sigma = -2\partial^\mu V^\sigma. \quad (263)$$

Using the identities $\partial^\nu \hat{x}^\mu = g^{\mu\nu}$ and $\partial^\nu \partial_\nu \hat{x}^\mu = 0$, the term $-2\partial^\mu V^\sigma$ appearing in Eq. (263) is exactly cancelled by its positive counterpart in Eq. (261).

Finally, the spacetime-geometry operator acting on V^σ yields

$$\hat{\mathcal{G}}^\mu V^\sigma = [\hat{x}^\lambda, \partial^\nu \Gamma_{\nu\lambda}^\mu]V^\sigma = \hat{x}^\lambda (\partial^\nu \Gamma_{\nu\lambda}^\mu V^\sigma) - \partial^\nu (\Gamma_{\nu\lambda}^\mu \hat{x}^\lambda V^\sigma) \quad (264)$$

$$= \hat{x}^\lambda (\partial^\nu \Gamma_{\nu\lambda}^\mu) V^\sigma + \hat{x}^\lambda \Gamma_{\nu\lambda}^\mu (\partial^\nu V^\sigma) - (\partial^\nu \Gamma_{\nu\lambda}^\mu) \hat{x}^\lambda V^\sigma - \Gamma_{\nu\lambda}^\mu (\partial^\nu \hat{x}^\lambda) V^\sigma - \Gamma_{\nu\lambda}^\mu \hat{x}^\lambda (\partial^\nu V^\sigma) \quad (265)$$

$$= [\hat{x}^\lambda, (\partial^\nu \Gamma_{\nu\lambda}^\mu)]_{\text{mult}} V^\sigma + [\hat{x}^\lambda, \Gamma_{\nu\lambda}^\mu]_{\text{mult}} (\partial^\nu V^\sigma) - \Gamma_{\nu\lambda}^\mu g^{\lambda\nu} V^\sigma = (\hat{\mathfrak{R}}^\mu - \Gamma_{\nu\lambda}^\mu g^{\lambda\nu}) V^\sigma. \quad (266)$$

The subscript ‘‘mult’’ denotes purely multiplicative commutators, which we define as the reduced Riemann–Ricci operator:

$$\hat{\mathfrak{R}}^\mu \equiv [\hat{x}^\lambda, g^{\rho\nu} \partial_\rho \Gamma_{\nu\lambda}^\mu]_{\text{mult}} + [\hat{x}^\lambda, \Gamma_{\nu\lambda}^\mu]_{\text{mult}} g^{\beta\nu} \partial_\beta. \quad (267)$$

The action of Eq. (258) on a quantum gravitational state V^σ can therefore be explicitly verified by substituting Eqs. (261), (263), and (266). This yields the explicit form of the metric-compatible tensor operator acting on V^σ :

$$\Psi^\mu V^\sigma = 2g^{\mu\nu}\Gamma_{\nu\lambda}^\sigma V^\lambda + [\hat{\mathfrak{R}}^\mu - \Gamma_{\nu\lambda}^\mu g^{\lambda\nu}] V^\sigma. \quad (268)$$

From the explicit forms of $\hat{\mathcal{K}}^\mu$ and $\hat{\mathcal{G}}^\mu$ (Eqs. (263) and (266)), we find

$$\left[\hat{\mathfrak{R}}^\mu - \Gamma_{\nu\lambda}^\mu g^{\lambda\nu} \right] V^\sigma \neq -2g^{\mu\nu} \Gamma_{\nu\lambda}^\sigma V^\lambda, \quad (269)$$

which implies that Eq. (268) is non-vanishing. Hence, the metric-compatible tensor operator Ψ^μ exhibits a non-trivial quantum-geometric structure. A full demonstration that this structure satisfies Eq. (260)—thereby resolving the physical singularity by ensuring that $\Psi^\mu V^\sigma$ remains finite at $r = 0$ —will be presented in a forthcoming paper. Nevertheless, a highly counterintuitive feature emerges from the metric-compatible tensor operator Ψ^μ in Eq. (268), together with the reduced Riemann–Ricci operator in Eq. (267). Here, the inverse metric components $g^{\mu\nu}$ arise recursively from the metric, in conjunction with the Christoffel symbols $\Gamma_{\nu\lambda}^\sigma$ and their derivatives $\partial_\rho \Gamma_{\nu\lambda}^\mu$. This recursive operator-level coupling generates a nonlinear back-reaction among the metric, its inverse, the connection, and the differential structure of spacetime. The feedback loop acts as an effective “infinite resistance”, preventing the geometry from approaching classically singular configurations. This self-regulating mechanism stabilises the spacetime and establishes the structural basis for the resolution of physical singularities, formalised in Theorem 3.

Theorem 3 (Self-Stabilisation of the Quantum Spacetime Differential Structure). *Let $\hat{g}_{\mu\nu}$ denote the quantum metric operator from Proposition 6, where the radial operator and the speed-of-light operator satisfy the noncommutative algebra of Eq. (258), with the spacetime operator restricted to the quantised radial degree of freedom. In the extreme-field regime of black-hole interiors, $\hat{g}_{\mu\nu}$ satisfies the implicit operator relation*

$$\hat{g}_{\mu\nu} = g_{\mu\nu}(\hat{r}, \hat{g}^{\alpha\beta}, \Gamma_{\nu\lambda}^\sigma, \partial_\rho \Gamma_{\nu\lambda}^\mu), \quad (270)$$

thereby defining a closed, nonlinear operator system that couples the metric, its inverse, the affine connection, and their derivatives. This relation elevates quantisation beyond the metric alone to a quantum theory of the full differential structure of spacetime. As $\|\hat{g}_{\mu\nu}\|$ approaches classically singular configurations, the coupled operators $\hat{g}^{\alpha\beta}$, $\Gamma_{\nu\lambda}^\sigma$, and $\partial_\rho \Gamma_{\nu\lambda}^\mu$ respond oppositely, generating a stabilising seesaw mechanism that dynamically suppresses divergences. This recursive hierarchy converges to a unique finite fixed point at which the metric is self-consistently determined—realising background independence—and all curvature invariants, including \widehat{R} and $\widehat{R_{\mu\nu\rho\sigma} R^{\mu\nu\rho\sigma}}$, are bounded operators, thereby providing intrinsic, non-perturbative regularisation of black-hole singularities without cut-offs or ad hoc modifications.

Remark 6 (The Physical Singularity as Temporal Collapse and the Extreme-Field Seesaw as Temporal Opposition). *In the classical Schwarzschild solution, the radial coordinate becomes timelike within the black-hole interior (Eq. (114): $r < r_h$), enforcing an inevitable one-directional trajectory toward the singularity—an inescapable moment in time. Theorem 3 shows that extreme-field regularisations permit a seesaw mechanism opposing this collapse, yielding finite fixed points consistent with Subsection 4.2, where extreme points of Hermiticity arise as conditional solutions to the Hermitian quantum-gravitational conjugate equations (226)–(236), supporting the conservative structure established in Remarks 1 and Theorem 1. This mechanism may also underpin the negative-energy anti-eigen flux identified in Subsections 3.4 and 3.5, describing entities that propagate backward in time, evading the singularity and tunnelling through the quantum-corrected horizon.*

To illustrate the seesaw regularisation mechanism from Theorem 3 and Remark 6, we now demonstrate that the expectation value of the Ricci scalar—a fundamental curvature invariant—remains finite in all physical states.

Proposition 12 (Boundedness of the Ricci scalar expectation value). *In the extreme-field regime described by the quantised Einstein equations (58), and under the assumptions of Proposition 2 and Remark 1, the expectation value of the Ricci scalar operator \hat{R} is finite in any physical state for which the stress–energy expectation values are finite. Explicitly,*

$$|\langle \hat{R} \rangle| < \infty. \quad (271)$$

Proof. We begin with the quantised Einstein field equations, as detailed in Remark 3,

$$\hat{R}_{\mu\nu} - \frac{1}{2} \hat{g}_{\mu\nu} \hat{R} = \kappa \hat{T}_{\mu\nu}, \quad \kappa = \frac{8\pi G}{c^4}. \quad (272)$$

Given that, in the extreme-field regime, classical quantities are reinterpreted as operators acting on a common Hilbert space (as established in Section 1), contracting Eq. (272) with the inverse quantised metric $\hat{g}^{\mu\nu}$ and invoking the operator invertibility condition $\hat{g}^{\mu\nu} \hat{g}_{\mu\nu} = \mathbb{I}$ (Eq. (24)), whose trace in 4D spacetime equals 4, yields

$$\hat{R} - 2\hat{R} = \kappa \hat{g}^{\mu\nu} \hat{T}_{\mu\nu} \implies \hat{R} = -\kappa \hat{T}, \quad (273)$$

where the trace operator is defined as $\hat{T} \equiv \hat{g}^{\mu\nu} \hat{T}_{\mu\nu}$. Eq. (273) is therefore an exact operator identity derived directly from the quantum Einstein equations, with no approximation involved.

Now consider an arbitrary physical state $|\phi\rangle$. By Proposition 2, Remark 1, and Theorem 3, together with the aforementioned quantum gravitational noncommutative algebra derived in this subsection, the operators $\hat{g}^{\mu\nu}$ possess bounded expectation values. That is, there exist constants $C_{\mu\nu}$ such that $|\langle\hat{g}^{\mu\nu}\rangle| \leq C_{\mu\nu}$ for all μ, ν .

Moreover, in any physically admissible state, the expectation values of the stress–energy tensor components $\langle\hat{T}_{\mu\nu}\rangle$ are required to be finite—a standard consistency condition in quantum field theory, where renormalisation guarantees finiteness. It follows that there exist constants $D_{\mu\nu}$ such that $|\langle\hat{T}_{\mu\nu}\rangle| \leq D_{\mu\nu}$. This conclusion is supported by Remark 3 (see Eq. (56)), which shows that the stress–energy tensor, when expressed as a functional of the quantised metric, is inherently finite for a finite quantised metric, as demonstrated in Theorem 3.

To estimate $\langle\hat{T}\rangle$, we address the operator ordering in the product $\hat{g}^{\mu\nu}\hat{T}_{\mu\nu}$ by adopting the symmetric (Weyl) ordering Sym [162]. This prescription maps real classical products to Hermitian operators, thereby ensuring real expectation values. Applying the triangle inequality for operator norms,

$$|\langle\hat{T}\rangle| = |\langle\text{Sym}(\hat{g}^{\mu\nu}\hat{T}_{\mu\nu})\rangle| \leq \|\text{Sym}(\hat{g}^{\mu\nu}\hat{T}_{\mu\nu})\|_{\text{op}} \leq \sum_{\mu,\nu} \|\hat{g}^{\mu\nu}\|_{\text{op}} \|\hat{T}_{\mu\nu}\|_{\text{op}} < \infty. \quad (274)$$

The finiteness of the operator norms follows from boundedness in the extreme-field regime, as established in Proposition 2, Remark 1, Remark 3, and Theorem 3.

Applying the trace relation (273) together with the bound (274) yields

$$|\langle\hat{R}\rangle| = \kappa |\langle\hat{T}\rangle| < \infty. \quad (275)$$

Consequently, the expectation value of the Ricci scalar remains finite in all physical states, including regimes where the classical curvature would diverge.

Higher-order invariants (e.g., $\langle\hat{R}_{\mu\nu}\hat{R}^{\mu\nu}\rangle$, Kretschmann scalar) can be analysed similarly via metric–Riemann combinations, but a complete treatment requires specifying the quantum state and is reserved for future work. \square

Remark 7. *Future work will also investigate whether Einstein’s constant $\kappa = 8\pi G/c^4$ acquires an operator character in the extreme-field regime, with $1/c^4$ promoted to $\hat{\mathcal{C}}^{-4}$ in accordance with Proposition 6. From Eq. (219), the inverse squared speed-of-light operator expands perturbatively as*

$$\hat{\mathcal{C}}^{-2} = \frac{1}{c^2} \left[1 + \frac{\hat{Z}^\mu \nabla_\mu}{c^2} \right]^{-1} = \frac{1}{c^2} \sum_{j=0}^{\infty} (-1)^j \left(\frac{\hat{Z}^\mu \nabla_\mu}{c^2} \right)^j. \quad (276)$$

Introducing the dimensionless operator $\epsilon \equiv \hat{Z}^\mu \nabla_\mu / c^2$, the expansion of $\hat{\mathcal{C}}^{-4}$ follows directly as

$$\hat{\mathcal{C}}^{-4} = \frac{1}{c^4} \sum_{j=0}^{\infty} (-1)^j (j+1) \epsilon^j = \frac{1}{c^4} + \frac{1}{c^4} \sum_{j=1}^{\infty} (-1)^j (j+1) \epsilon^j. \quad (277)$$

Truncating this series at order $j = p$ and defining κ^p as the perturbative correction term from Eq. (277) (multiplied by $8\pi G$), the right-hand side of Eq. (58) becomes the classical stress-energy tensor augmented by perturbative extreme-field corrections:

$$\hat{\kappa} \hat{T}_{\mu\nu} \approx \kappa T_{\mu\nu}^{\text{classical}} + \kappa^p T_{\mu\nu}^{\text{quantum}} + \mathcal{O}(\epsilon^{p+1}) + \text{cross terms}. \quad (278)$$

Considering the Einstein tensor operator $\hat{G}_{\mu\nu}$ —the left-hand side of Eq. (58)—it depends purely on the metric tensor and encodes the spacetime geometry. Expanding the quantised metric into a classical part plus quantum corrections, as explored in Eqs. (128), (129), and (130), yields a decomposition into a classical component with extreme-field operator-valued corrections:

$$\hat{G}_{\mu\nu} \approx G_{\mu\nu}^{\text{classical}} + G_{\mu\nu}^{\text{quantum}}. \quad (279)$$

Together with Eq. (278), this recovers the classical Einstein field equations (Eq. (57)) augmented by extreme-field quantum corrections, in full consistency with the quantised formulation established in Remark 3.

Our attention now turns to the constant factor $i\hbar\alpha$, and in particular to the unique extreme-field coupling α —a quantity that encapsulates fundamental quantum-gravitational effects. Dimensional consistency requires α to carry units of [Length²/(Mass · Time)]. Within the non-perturbative extreme-field framework of the primordial Planck black hole (PPBH), α can be expressed in terms of the fundamental Planck units: the Planck length l_{Pl} , Planck time t_{Pl} , and Planck mass M_{Pl} , defined as [83, 118, 128]:

$$l_{\text{Pl}} = \sqrt{\frac{\hbar G}{c^3}}, \quad t_{\text{Pl}} = \sqrt{\frac{\hbar G}{c^5}}, \quad M_{\text{Pl}} = \sqrt{\frac{\hbar c}{G}}. \quad (280)$$

Accordingly, α adopts the Planck unit combination [Planck Length²/(Planck Mass · Planck Time)], yielding

$$\alpha = \frac{\ell_{\text{Pl}}^2}{M_{\text{Pl}} t_{\text{Pl}}} = \frac{\frac{\hbar G}{c^3}}{\sqrt{\frac{\hbar c}{G}} \cdot \sqrt{\frac{\hbar G}{c^5}}} = \frac{G}{c}. \quad (281)$$

Eq. (281) relates the dimensional constant α to G and c , encoding the fundamental spacetime properties of extreme-field physics. The commutation relation between $\hat{Z}^\nu \nabla_\nu$ and \hat{x}^μ then takes the form:

$$[\hat{x}^\mu, \hat{Z}^\nu \nabla_\nu] = i \frac{\hbar G}{c} \Psi^\mu. \quad (282)$$

A comprehensive quantum-gravity framework, however, requires extending beyond Riemannian geometry to incorporate intrinsic non-metricity, i.e., $\nabla_\lambda g_{\mu\nu} \neq 0$ [47, 77, 81]. In this setting, the commutator of the quantum-gravitational conjugate—written as $\hat{Z}_\nu \nabla^\nu$ to avoid explicit dependence on the inverse metric and \hat{x}^μ —takes the form:

$$[\hat{x}^\mu, \hat{Z}_\nu \nabla^\nu] = i \frac{\hbar G}{c} \Xi^\mu + \mathbf{N}^\mu, \quad (283)$$

where Ξ^μ —the non-metricity quantum-gravitational operator generalising Ψ^μ from Eq. (282)—includes non-metric effects and additional corrections, compactly denoted \mathbf{N}^μ , which may arise from

$$\partial_\nu \hat{x}^\mu = \delta_\nu^\mu + \mathfrak{E}_\nu^\mu, \quad \frac{\hbar G}{c} + \text{constant}, \quad \mathcal{M}[g_{\mu\nu}, \nabla_\lambda g_{\mu\nu}, g_{\mu\nu} V^\mu], \quad (284)$$

where \mathfrak{E}_ν^μ denotes an extreme-field correction, along with a constant term and the metric-dependent functional $\mathcal{M}[g_{\mu\nu}, \nabla_\lambda g_{\mu\nu}, g_{\mu\nu} V^\mu]$, encoding quantum-gravitational effects from non-metricity and non-Riemannian geometry.

Using Eqs. (243)–(249) within the primary commutation relation of Eq. (240), the action of Ξ^μ on a test function f can be explicitly examined, yielding:

$$\begin{aligned} \Xi^\mu f &= \nabla^\mu f + (\partial_\nu \hat{x}^\mu) \nabla^\nu f + \hat{x}^\mu (\partial_\nu \nabla^\nu) f + \hat{x}^\mu \nabla_\nu (\partial_\nu f) - (\partial_\nu \nabla^\nu) \hat{x}^\mu f - \nabla^\nu (\partial_\nu \hat{x}^\mu) f - \nabla^\nu \hat{x}^\mu (\partial_\nu f) \\ & \quad (285) \end{aligned}$$

$$= \left(\nabla^\mu + [\hat{x}^\mu, \partial_\nu \nabla^\nu]_{\text{mult}} + [\hat{x}^\mu, \nabla^\nu]_{\text{mult}} \partial_\nu \right) f. \quad (286)$$

A key feature of Eq. (286) is that it provides a consistent generalised approach to the singularity problem, treating black-hole cores as intrinsically non-Riemannian within a non-metricity geometry. The non-metricity operator Ξ^μ (equivalently \mathbf{N}^μ) acts on a quantum-gravitational state $f \equiv f(g_{\mu\nu})$ in a highly non-trivial manner, reflecting the enriched differential structure of spacetime and encoding its quantum-geometric properties—thereby supporting the resolution of the classical singularity through the self-stabilisation of the quantum spacetime, as formalised in Theorem 3.

Another feature of the non-metricity model is the annihilation of quantum-gravitational states at the boundary between extreme and classical fields (recall Eq. (237)). This arises from the intrinsic properties of these states. For a black hole, the boundary is the horizon, yielding the condition

$$\lim_{r \rightarrow r_s} \nabla_\mu |\phi\rangle = 0. \quad (287)$$

Within this model, which strictly enforces the non-metricity condition

$$\nabla_\lambda g_{\mu\nu}^* \neq 0, \quad (288)$$

the metric $g_{\mu\nu}^*$ characterises the extreme-field regime. Since quantum-gravitational states depend explicitly on the metric and its derivatives, the covariant derivative does not annihilate them in this regime:

$$\nabla_\mu |\phi(g_{\alpha\beta}^*)\rangle \equiv \lim_{r \rightarrow r_s^-} \nabla_\mu |\phi(g_{\alpha\beta})\rangle \neq 0. \quad (289)$$

At and beyond the horizon ($r \geq r_s$), the metric becomes metric-compatible, restoring classical Riemannian geometry. In this regime, the covariant derivative annihilates quantum-gravitational states:

$$\lim_{r \rightarrow r_s^+} \nabla_\mu |\phi(g_{\alpha\beta})\rangle = 0, \quad (290)$$

capturing the extreme–classical transition within non-metricity.

Finally, we examine more deeply the coupling constant $\alpha = G/c$ (Eq. (281), introduced in Eq. (242)). The commutation relation Eq. (282) (comprehensive form Eq. (283)) predicts that in the extreme-field regime, α encodes the fundamental interplay between quantum mechanics and gravity. Yet $\alpha = G/c$ captures only basic gravitational coupling—a partial description. A complete framework, particularly in extreme-field regimes, requires promoting α to a running coupling $\alpha(\mu)$, which is justified by the Hermiticity condition for the quantum-gravitational conjugate: this demands that α emerge as a constant modulated by a differential operator, as proposed in Proposition 2 (Eq. (45)). To avoid conflict with Eq. (242), this condition is mitigated by rendering it coordinate-independent in Subsection 4.2 (Eq. (230)), where the operator structure remains satisfied for $\alpha_0 \equiv G/c$ with ∂_u as a scale operator: its action on a quantum-gravitational state generates the effective running scale, captured by the dimensionless coupling $Q(\mu)$. In direct analogy to the

renormalisation-group running of gauge couplings [52, 63, 68], we therefore propose the running coupling

$$\alpha(\mu) = \frac{G Q(\mu)}{c}, \quad (291)$$

where $Q(\mu)$ captures the effective unified quantum-gravitational running,

$$Q(\mu) = h(\alpha_g(\mu), \alpha_w(\mu), \alpha_e(\mu), \alpha_s(\mu)). \quad (292)$$

The function h encapsulates the mutual running of the dimensionless gravitational (α_g), weak (α_w), electromagnetic (α_e) and strong (α_s) couplings.

The running coupling $\alpha(\mu)$ defined in Eq. (291) provides a unified description of the fundamental interactions at the Planckian limit, where a single symmetric framework is believed to merge the four forces. This unification is signalled by the convergence

$$\alpha_g \approx \alpha_w \approx \alpha_s \approx \alpha_e \sim \mathcal{O}(1), \quad (293)$$

rendering the forces effectively indistinguishable. Subsequent cosmic inflation broke this symmetry, leading to gravitational decoupling as α_G became exponentially attenuated relative to the quantum gauge couplings. Consequently, gravity was rendered negligible at quantum scales, its influence confined to macroscopic domains while the other forces govern particle interactions.

At the Planckian extremisation threshold, the universe enters a state of extreme negentropy wherein quantum fluctuations freeze, and the primordial Planck black hole crystallises as a deterministic state—initiating the extremisation phase transition of Subsection 3.7. The onset of the classical regime then restores the constancy of the speed of light:

$$\hat{\mathcal{C}} \rightarrow c. \quad (294)$$

Following the extremisation phase transition, in the low-energy regime $\mu \rightarrow 0$, the quantum gauge couplings approach their reference values, while the gravitational coupling vanishes:

$$\lim_{\mu \rightarrow 0} \alpha_g(\mu) = 0, \quad \lim_{\mu \rightarrow 0} \alpha_w(\mu) = \alpha_w, \quad \lim_{\mu \rightarrow 0} \alpha_e(\mu) = \alpha_e, \quad \lim_{\mu \rightarrow 0} \alpha_s(\mu) = \alpha_s. \quad (295)$$

Consequently, the unified coupling itself vanishes,

$$\lim_{\mu \rightarrow 0} \alpha(\mu) = \frac{G}{c} \lim_{\mu \rightarrow 0} Q(\mu) = 0, \quad (296)$$

signalling complete gravitational decoupling and the emergence of the four distinct forces from their primordial symmetry.

4.4 The Extreme-Field Equation

The extreme-field equation generalises the energy–momentum dispersion relation in curved spacetime to describe black-hole extreme-field dynamics.

In curved spacetime, the invariant mass formula is expressed as [11, 110]:

$$m^2 c^2 = -g^{\mu\nu} p_\mu p_\nu, \quad (297)$$

which can be expanded to:

$$m^2 c^2 = -g^{00} p_0^2 - g^{ii} p_i p_i. \quad (298)$$

The temporal component of Eq. (298) corresponds to the energy:

$$p_0 = -\frac{E}{c}. \quad (299)$$

The spatial component corresponds to the spatial momentum vector p_i , which its magnitude can be written as follows:

$$p_i p_i = \mathbf{p}^2. \quad (300)$$

Substituting Eqs. (299) and (300) into Eq. (298), and incorporating energy considerations, we scale the expression by c^2 , yielding:

$$m^2 c^4 = -g^{00} E^2 - g^{ii} c^2 \mathbf{p}^2. \quad (301)$$

At the Planckian extremisation threshold, the Primordial Planck Black Hole (PPBH) attains an unstable yet deterministic state—a coherent configuration emerging from the coalescence of discrete extreme quantum states into a single continuum-like structure. This deterministic configuration naturally provides the primordial seed for cosmic emergence.

To elucidate this behaviour, we consider, for a discrete extreme excitation state, the operator analogue of the curved-spacetime invariant-mass relation, Eq. (297). In this formulation, the metric tensor and the speed of light are promoted to the non-perturbative operators $\hat{j}^{\mu\nu}$ and $\hat{\mathcal{C}}$, respectively, as established in Propositions 4 and 6:

$$m^2 \hat{\mathcal{C}}^2 = -\hat{j}^{\mu\nu} \hat{p}_\mu \hat{p}_\nu. \quad (302)$$

By analogy with the preceding derivation, the operator form of Eq. (301) reads

$$m^2 \hat{\mathcal{C}}^4 = -\hat{j}^{00} \hat{E}^2 - \hat{\mathcal{C}}^2 \hat{j}^{ii} \hat{\mathbf{p}}^2. \quad (303)$$

The energy contribution is accounted for by rearranging Eq. (303) into its expectation-value form:

$$\langle -\hat{j}^{00}\hat{\mathbf{E}}^2 \rangle = \langle \hat{\mathcal{C}}^2 \hat{j}^{ii} \hat{\mathbf{p}}^2 \rangle + m^2 \langle \hat{\mathcal{C}}^4 \rangle. \quad (304)$$

The quantity $\langle -\hat{j}^{00}\hat{\mathbf{E}}^2 \rangle$ is thus interpreted as the squared energy of an extreme-field excitation state, denoted U^2 . This energy is assumed to comprise three contributions: the vacuum ground-state fluctuation ϵ , the virtual-particle energy V , and a universal parameter \mathbf{u} encompassing any further sources. Accordingly, we define

$$\sqrt{\langle -\hat{j}^{00}\hat{\mathbf{E}}^2 \rangle} \equiv U = V + \epsilon + \mathbf{u}. \quad (305)$$

Consequently, upon identifying $\hat{\mathfrak{D}} \equiv \hat{j}^{ii}$ as the spatial oscillator operator, Eq. (304) can be reformulated as

$$U^2 = \langle \hat{\mathcal{C}}^2 \hat{\mathfrak{D}} \hat{\mathbf{p}}^2 \rangle + m^2 \langle \hat{\mathcal{C}}^4 \rangle. \quad (306)$$

Using the explicit form of Eq. (305), the extreme-field equation, Eq. (306), can be rewritten as

$$(V + \epsilon + \mathbf{u})^2 = \langle \hat{\mathcal{C}}^2 \hat{\mathfrak{D}} \hat{\mathbf{p}}^2 \rangle + m^2 \langle \hat{\mathcal{C}}^4 \rangle. \quad (307)$$

We now substitute into Eq. (307) the non-perturbative operator speed of light, $\hat{\mathcal{C}}$, defined in Eq. (219) and incorporating the quantum-gravitational conjugate that characterises the full non-perturbative quantum-gravity dynamics. This substitution yields:

$$(V + \epsilon + \mathbf{u})^2 = \left\langle \left(\sqrt{c^2 + \hat{Z}^\mu \nabla_\mu} \right)^2 \hat{\mathfrak{D}} \hat{\mathbf{p}}^2 \right\rangle + m^2 \left\langle \left(\sqrt{c^2 + \hat{Z}^\mu \nabla_\mu} \right)^4 \right\rangle. \quad (308)$$

Simplifying the squares and fourth powers:

$$(V + \epsilon + \mathbf{u})^2 = \langle (c^2 + \hat{Z}^\mu \nabla_\mu) \hat{\mathfrak{D}} \hat{\mathbf{p}}^2 \rangle + m^2 \langle (c^2 + \hat{Z}^\mu \nabla_\mu)^2 \rangle. \quad (309)$$

Next, expanding the right-hand side:

$$(V + \epsilon + \mathbf{u})^2 = c^2 \langle \hat{\mathfrak{D}} \hat{\mathbf{p}}^2 \rangle + \langle \hat{Z}^\mu \nabla_\mu \hat{\mathfrak{D}} \hat{\mathbf{p}}^2 \rangle + m^2 c^4 + 2m^2 c^2 \langle \hat{Z}^\mu \nabla_\mu \rangle + m^2 \langle (\hat{Z}^\mu \nabla_\mu)^2 \rangle. \quad (310)$$

Expanding the LHS and rearranging the RHS, one gets:

$$V^2 + \epsilon^2 + \mathbf{u}^2 + 2V\epsilon + 2V\mathbf{u} + 2\epsilon\mathbf{u} = c^2 \langle \hat{\mathcal{D}}\hat{\mathbf{p}}^2 \rangle + m^2 c^4 \\ + \langle \hat{Z}^\mu \nabla_\mu \hat{\mathcal{D}}\hat{\mathbf{p}}^2 \rangle + 2m^2 c^2 \langle \hat{Z}^\mu \nabla_\mu \rangle + m^2 \langle (\hat{Z}^\mu \nabla_\mu)^2 \rangle. \quad (311)$$

For real particles, the on-shell energy-momentum dispersion relation satisfies:

$$H^2 = c^2 \langle \hat{\mathcal{D}}\hat{\mathbf{p}}^2 \rangle + m^2 c^4. \quad (312)$$

Thus, considering virtual particles with energy V , the energy of the vacuum ground state ϵ , the universal parameter \mathbf{u} and real particles with on-shell energy H , Eq. (311) can be reformulated as:

$$V^2 + \epsilon^2 + \mathbf{u}^2 + 2V\epsilon + 2V\mathbf{u} + 2\epsilon\mathbf{u} = H^2 + \langle \hat{Z}^\mu \nabla_\mu \hat{\mathcal{D}}\hat{\mathbf{p}}^2 \rangle + 2m^2 c^2 \langle \hat{Z}^\mu \nabla_\mu \rangle + m^2 \langle (\hat{Z}^\mu \nabla_\mu)^2 \rangle. \quad (313)$$

Next, we write Eq. (313) in the explicit quantum operator form:

$$\hat{V}^2 + \hat{\epsilon}^2 + \hat{\mathbf{u}}^2 + 2\hat{V}\hat{\epsilon} + 2\hat{V}\hat{\mathbf{u}} + 2\hat{\epsilon}\hat{\mathbf{u}} = \hat{H}^2 + \hat{Z}^\mu \nabla_\mu \hat{\mathcal{D}}(i\hbar\nabla)^2 + 2m^2 c^2 \hat{Z}^\mu \nabla_\mu + m^2 (\hat{Z}^\mu \nabla_\mu)^2, \quad (314)$$

where ∇^2 is the Laplacian operator. Using the quantum momentum operator $i\hbar\nabla$, we proceed to move \hat{H}^2 to the left-hand side of Eq. (314) as follows:

$$\hat{V}^2 - \hat{H}^2 + \hat{\epsilon}^2 + \hat{\mathbf{u}}^2 + 2\hat{V}\hat{\epsilon} + 2\hat{V}\hat{\mathbf{u}} + 2\hat{\epsilon}\hat{\mathbf{u}} = -\hbar^2 \hat{Z}^\mu \nabla_\mu \hat{\mathcal{D}}\nabla^2 + 2m^2 c^2 \hat{Z}^\mu \nabla_\mu + m^2 (\hat{Z}^\mu \nabla_\mu)^2. \quad (315)$$

In quantum gravity, the extreme field $\Phi_{\mathcal{E}}$ is a scalar field incorporating couplings to curvature, gauge fields, quantised geometry, and extreme non-commutative quantum-gravitational effects. Its most general expression is

$$\Phi_{\mathcal{E}} \equiv \Phi(\phi, \{b_i\}, X), \quad (316)$$

where ϕ is the classical scalar, $\{b_i\}$ encodes internal degrees of freedom (basis, gauge, quantum states), and X subsumes further extreme unspecified contributions. Within this framework, the covariant derivative acquires deeper meaning: in classical general relativity, ∇_μ describes the evolution of vectors and their basis, but for the extreme field $\Phi_{\mathcal{E}}$ it is elevated to a tensorial operator $\hat{\nabla}_\mu$ governing field propagation and internal structure—basis, couplings to quantised geometry, and field–spacetime interplay. Its action on the extreme field is central and must not vanish at the extremisation

phase transition, where the non-perturbative evolution of the PPBH's internal structure drives boundary term preservation—forming the foundation for the quantum-gravitational jerking, as established in Subsection 4.5.

Consider now the Primordial Planck Black Hole (PPBH) and the evolution of its extreme states, depicted as a lattice of fluctuating spacetime points. Inward-pointing arrows represent the positive-energy eigen-Hawking radiation flowing toward regions of increasing curvature, while the negative-energy anti-eigen flux propagates in the opposite direction, escaping through the quantum-corrected horizon. At the Planckian threshold, an extremisation phase transition becomes inevitable: the PPBH recoils suddenly—the quantum-gravitational jerk—propelling it toward a large-scale structure. A fundamental condition for this transition is that the extreme field must not vanish at the boundary, owing to the non-perturbative eigen back-reaction of the field upon itself. This ensures that inflation follows inexorably.

Building on this picture, we now apply Eq. (315) to $\Phi_{\mathcal{E}}$ (Eq. (316)), which encodes the non-perturbative extreme-field physics that drives it:

$$\left[\hat{V}^2 - \hat{H}^2 + \hat{\epsilon}^2 + \hat{u}^2 + 2\hat{V}\hat{\epsilon} + 2\hat{V}\hat{u} + 2\hat{\epsilon}\hat{u} \right] \Phi_{\mathcal{E}} = \left[-\hbar^2 \hat{Z}^\mu \nabla_\mu \hat{\mathcal{D}} \nabla^2 + 2m^2 c^2 \hat{Z}^\mu \nabla_\mu + m^2 (\hat{Z}^\mu \nabla_\mu)^2 \right] \Phi_{\mathcal{E}}. \quad (317)$$

We characterise the difference $(\hat{V}^2 - \hat{H}^2)$ in Eq. (317) as follows:

$$\hat{\mathfrak{W}}^2 \equiv \hat{V}^2 - \hat{H}^2, \quad (318)$$

where $\hat{\mathfrak{W}}^2$ denotes the operator corresponding to the squared extreme-field energy gap. Therefore, Eq. (317) can be reformulated as:

$$\left[\hat{\mathfrak{W}}^2 + \hat{\epsilon}^2 + \hat{u}^2 + 2\hat{V}\hat{\epsilon} + 2\hat{V}\hat{u} + 2\hat{\epsilon}\hat{u} \right] \Phi_{\mathcal{E}} = \left[-\hbar^2 \hat{Z}^\mu \nabla_\mu \nabla^2 + 2m^2 c^2 \hat{Z}^\mu \nabla_\mu + m^2 (\hat{Z}^\mu \nabla_\mu)^2 \right] \Phi_{\mathcal{E}}. \quad (319)$$

One might define the intrinsic interaction terms on the left hand side as:

$$\hat{\mathcal{I}} = \hat{V}\hat{\epsilon} + \hat{V}\hat{u} + \hat{\epsilon}\hat{u}. \quad (320)$$

Recalling the definition of the quantum operator \hat{Z}^μ from Eq. (245), we can reinterpret the quantum gravitational conjugate as:

$$\hat{Z}^\mu \nabla_\mu = i\hbar\alpha \partial^\mu \nabla_\mu = \partial^\mu \hat{\mathcal{P}}_\mu = \hat{\mathcal{F}}. \quad (321)$$

We define the quantum-gravitational momentum operator $\hat{\mathcal{P}}_\mu \equiv i\hbar\alpha \nabla_\mu$. By analogy with classical mechanics, where force is the time derivative of momentum, we introduce the conjugate force term $\hat{\mathcal{F}} = \partial^\mu \hat{\mathcal{P}}_\mu$. Eqs. (320)

and (321) then allow Eq. (319) to be recast as:

$$\left[\hat{\mathfrak{W}}^2 + \hat{\epsilon}^2 + \hat{\mathbf{u}}^2 + 2\hat{\mathcal{L}} \right] \Phi_{\mathcal{E}} = \left[-\hbar^2 \hat{\mathcal{F}} \nabla^2 + 2m^2 c^2 \hat{\mathcal{F}} + m^2 \hat{\mathcal{F}}^2 \right] \Phi_{\mathcal{E}}. \quad (322)$$

Consequently, the left-hand side of Eq. (322) constitutes a composite system comprising the vacuum, virtual and real particles, a universal parameter, and all interaction terms. This system encodes complete information about the nature of a physical entity, which motivates the introduction of a quiddity operator, $\hat{\mathbb{Q}}$. The term quiddity denotes the inherent essence or “whatness” of an entity; here the operator encapsulates the essential nature of the created state. We define the quiddity operator as:

$$\hat{\mathbb{Q}} \equiv \hat{\mathfrak{W}}^2 + \hat{\epsilon}^2 + \hat{\mathbf{u}}^2 + 2\hat{\mathcal{L}}. \quad (323)$$

The right-hand side of Eq. (322) describes the intrinsic quantum-gravitational geometry of the system, represented by the quantum-gravitational geometric operator $\hat{\mathbb{G}}$, defined as:

$$\hat{\mathbb{G}} \equiv -\hbar^2 \hat{\mathcal{F}} \hat{\mathcal{D}} \nabla^2 + 2m^2 c^2 \hat{\mathcal{F}} + m^2 \hat{\mathcal{F}}^2. \quad (324)$$

Using the quiddity operator $\hat{\mathbb{Q}}$ from Eq. (323) and the quantum-gravitational geometric operator $\hat{\mathbb{G}}$ from Eq. (324), we recast Eq. (322) into its final form—the extreme-field equation. This equation governs the emergence of an entity whose properties are dictated by the quantum-gravitational geometry encoded in $\hat{\mathbb{G}}$. We thus obtain:

$$\hat{\mathbb{Q}} \Phi_{\mathcal{E}} = \hat{\mathbb{G}} \Phi_{\mathcal{E}}. \quad (325)$$

Eq. (325) represents the compact formulation of the extreme-field equation: the entity generated by the action of the quiddity operator on the extreme field is intrinsically linked to a unique geometry encoded within the quantum-gravitational geometric operator. This equation therefore establishes that every physical entity possesses a distinct quiddity, fundamentally tied to its underlying quantum-gravitational geometry. Consider, for example, a physical entity **A**: when produced via the action of the quiddity operator on the extreme field, it acquires a unique geometric signature that emerges from the action of the geometric operator on the field—a defining hallmark of its quiddity that distinguishes it from all other entities. Should two entities, **B** and **A**, share precisely the same geometric signature, they are, in essence, the same unique entity. Within the non-perturbative framework, the equation thus implies a singular cosmic origin: the Universe emerges from a single Planckian state—namely, the primordial Planck-black-hole geometry.

4.5 The Quantum-Gravitational Jerking

Building on the Planckian extremisation threshold (Subsection 3.7), we examine the extreme-field equation Eq. (325) (explicitly given in Eq. (322)). At this threshold, primordial negentropy is extremised:

$$\Delta S \ll 0, \quad (326)$$

giving rise to the Primordial Planck Black Hole (PPBH)—a deterministic yet unstable Planck-scale state. This configuration drives an extremisation phase transition that underlies cosmic emergence. The transition is characterised by three physical limits:

$$\lim_{\Delta S \rightarrow -\infty} (\hat{\epsilon}^2 + \hat{u}^2) \Phi_{\mathcal{E}} \approx \hat{u}^2 \Phi_{\mathcal{E}}, \quad (327)$$

$$\lim_{\Delta S \rightarrow -\infty} \mathfrak{W}^2 \Phi_{\mathcal{E}} = 0, \quad (328)$$

$$\lim_{\Delta S \rightarrow -\infty} \hat{\mathcal{I}} \rightarrow \hat{\mathcal{P}}|_{\text{MAX}}. \quad (329)$$

Eq. (326) governs the extremisation of primordial fluctuations—particularly those within the PPBH—yielding a deterministic configuration at the Planckian limit. This state is intrinsically unstable due to the infinite resistance to singularity formation established by Theorem 3. The extremisation condition further entails Eq. (327), signifying the dominance of the universal parameter \hat{u} and the effective suppression of the vacuum contribution. Eq. (328) implies an additional consequence: the extreme-field energy gap vanishes at this threshold. Virtual particles thereby become real, precipitating an inevitable transition to a causal, non-Planckian regime. From the definition of the squared energy-gap operator in Eq. (318), this condition yields

$$\hat{V}^2 = \hat{H}^2. \quad (330)$$

Extreme-field excitations (virtual particles) thus converge toward a classical state analogous to real particles. This transition marks an unstable ordered Planckian configuration—one that behaves deterministically rather than probabilistically. Condition (329) signals that primordial universal interactions have reached an extremised tipping point, whereupon their nature fundamentally changes. Interactions cease, leaving the PPBH core as an ordered Planckian energy-density state, with an intrinsic non-perturbative pressure $\hat{\mathcal{P}}$ maximised beyond gravitational-collapse thresholds.

A pressure gradient in this unstable state results from the seesaw-like mechanism generating infinite resistance to singularity formation. This recoil—the inverse of black-hole formation and, equivalently, of Page-curve dynamics (Proposition 8)—reverses the entropic arrow. Whereas conventional collapse requires gravity to overcome outward pressure, the non-perturbative regime at the negentropic Planckian limit reverses this process (Theorem 2). As criticality is approached, the non-perturbative outward pressure exceeds the inward gravitational collapse force, triggering a violent recoil that initiates the universe’s emergence. Classically, this jerk limit corresponds to the condition

$$P_{\text{out}} > P_{\text{in}}, \quad (331)$$

where P_{out} denotes the outward Jerk Tremendous pressure and P_{in} the effective inward gravitational pressure (which may approach the Planck pressure). This inequality defines the critical threshold for the jerk event, thereby reversing gravitational collapse and seeding the creation of the universe.

Consequently, the inverse black-hole formation mechanism—embodied in the invertibility of Page-curve dynamics—physically interprets cosmic genesis within extreme field theory, realised via the quantum-gravitational jerking mechanism. Modelled by a decoupling jerk operator \mathcal{J} acting as the inverse of \hat{Z}^μ in $\hat{Z}^\mu \nabla_\mu$, this mechanism disentangles the quantum and gravitational sectors: quantum dynamics remain confined to the Hilbert space, while gravity emerges classically through a purely classical contribution. This decoupling catalysed the *Jerk Tremendous*—the recoil that drove rapid inflation—suggesting the form

$$\mathcal{J} = \nabla_\omega \hat{A}^\omega. \quad (332)$$

A key requirement is the selective separation of classical and quantum effects by the decoupling jerk operator. On the first-order conjugate $\hat{Z}^\nu \nabla_\nu$, it suppresses quantum contributions, leaving a classical gravitational term with residual quantum corrections. On the second-order conjugate $(\hat{Z}^\nu \nabla_\nu)^2$, it generates purely quantum phenomena confined to the Hilbert space. This behaviour is formalised by the decoupling relation

$$\hat{A}^\mu \hat{Z}^\nu = g^{\mu\nu} G, \quad (333)$$

establishing that \hat{A}^μ inverts \hat{Z}^ν , yielding a purely classical contribution.

To further elucidate these properties, we now revisit the operator form of Z^μ from Eq. (245):

$$\hat{Z}^\mu = i\hbar\alpha\partial^\mu. \quad (334)$$

The decoupling relation in Eq. (333) necessitates that the operator \hat{A}^μ act as the inverse of \hat{Z}^μ , which we express mathematically as:

$$G^{-1}\hat{A}^\mu = (\hat{Z}^\mu)^{-1}. \quad (335)$$

The inverse $(\hat{Z}^\mu)^{-1}$ is dictated by the structure of \hat{Z}^μ itself. As a first-order partial differential operator (Eq. (334)), its inverse \hat{A}^μ must reverse the action of \hat{Z}^μ on any function f . This action is expressed mathematically as

$$\hat{Z}^\mu f = i\hbar\alpha g^{\mu\nu} \frac{\partial f}{\partial x^\nu}. \quad (336)$$

The operator \hat{A}^μ must therefore satisfy $\hat{A}^\mu \hat{Z}_\mu f = Gf$. Using the covariant form $\hat{Z}_\mu = i\hbar\alpha\partial_\mu$ from Eq. (334), this condition becomes

$$\hat{A}^\mu (i\hbar\alpha\partial_\mu) f = Gf. \quad (337)$$

This implies that \hat{A}^μ must invert \hat{Z}_μ , i.e., act as an integral operator that undoes differentiation. We therefore propose an operator of the form:

$$\hat{A}^\mu = -i\frac{G}{\hbar\alpha} \int dx^\mu. \quad (338)$$

To verify Eq. (333), we apply \hat{A}^μ to $\hat{Z}^\nu f$. Using the definition of \hat{A}^μ from Eq. (338), we obtain

$$\hat{A}^\mu \hat{Z}^\nu f = -i\frac{G}{\hbar\alpha} \int dx^\mu \left(i\hbar\alpha g^{\mu\nu} \frac{\partial f}{\partial x^\mu} \right) = G \int dx^\mu g^{\mu\nu} \partial_\mu f. \quad (339)$$

Treating the integral as indefinite and integrating by parts gives

$$\hat{A}^\mu \hat{Z}^\nu f = G \left(g^{\mu\nu} f - \int dx^\mu (\partial_\mu g^{\mu\nu}) f + b^{\mu\nu} \right), \quad (340)$$

with $g^{\mu\nu} f$ as the boundary term, the integral capturing metric variations, and $b^{\mu\nu}$ arbitrary constants independent of the integration variable.

It follows that the decoupling relation (333) is globally satisfied if and only if the first term—the boundary term—in Eq. (340) survives while the others vanish. In quantum field theory, however, boundary terms are conventionally taken to vanish under the assumption of sufficiently rapid field decay at infinity [102]. Adopting this assumption, each term in Eq. (340) would

individually vanish: $g^{\mu\nu} f \rightarrow 0$, the metric-variation integral decays during extremisation, and the constants $b^{\mu\nu}$ vanish. Eq. (340) would then reduce identically to zero, thereby contradicting the requirement of Eq. (333).

The necessary condition for resolving this contradiction is that the boundary term must not be assumed to vanish—thereby contravening the standard convention of quantum field theory. This condition obtains when the system is open or dynamically coupled to its boundaries [88]. In our framework, we mirror this by requiring that the emerged unified state of the PPBH evolve at the boundaries, i.e. $\nabla_\mu \Phi_\mathcal{E} \neq 0$, thereby contravening the freezing of quantum-gravitational states at the boundary established in Subsection 4.2 (Eq. (237)). This condition is justified in Subsection 4.3 (Theorem 3) and is physically realised as an effective infinite resistance that prevents the extreme Planckian configuration from collapsing indefinitely to a singularity; consequently, the field admits significant evolution via an intrinsic non-perturbative eigen back-reaction. This is precisely the extremisation phase transition, wherein the non-vanishing of boundary terms—as discussed in the context of correlation functions—constitutes the very origin of the transition [136].

Should the boundary contribution in Eq. (340) survive as the intrinsic origin of the extremisation phase transition, the quantum-gravitational jerking mechanism thereby satisfies the decoupling relation of Eq. (333), marking the onset of full decoupling of quantum gravitational effects.

The quantum-gravitational jerking mechanism is therefore established as the natural mechanism explaining the non-perturbative extreme evolution of the universe’s origin—from the unstable Planckian state of the Primordial Planck Black Hole to a non-Planckian classical structure. Physically, we interpret this jerk as the action of the decoupling jerk operator on the quantum-gravitational conjugate, thereby suppressing quantum effects and giving rise to asymptotically free gravitational contributions. Within the compact formulation of the extreme-field equation, Eq. (325), this conjugate is encoded in the quantum-gravitational geometric operator $\hat{\mathbb{G}}$ (defined in Eq. (319)) as:

$$\hat{\mathbb{G}} \Phi_\mathcal{E} \equiv \left(-\hbar^2 \hat{Z}^\mu \nabla_\mu \hat{\mathcal{D}} \nabla^2 + 2m^2 c^2 \hat{Z}^\mu \nabla_\mu + m^2 (\hat{Z}^\mu \nabla_\mu)^2 \right) \Phi_\mathcal{E}. \quad (341)$$

Applying the decoupling jerk operator, defined in Eq. (332), to Eq. (325) elucidates the physical implications of the quantum-gravitational jerk:

$$\mathcal{J} \hat{\mathbb{Q}} \Phi_\mathcal{E} = \mathcal{J} \hat{\mathbb{G}} \Phi_\mathcal{E}. \quad (342)$$

With $\hat{\mathbb{G}}$ explicitly given by Eq. (341), we obtain:

$$\mathcal{J} \hat{\mathbb{Q}} \Phi_{\mathcal{E}} = \mathcal{J} \left(-\hbar^2 \hat{Z}^{\mu} \nabla_{\mu} \hat{\mathfrak{D}} \nabla^2 + 2m^2 c^2 \hat{Z}^{\mu} \nabla_{\mu} + m^2 (\hat{Z}^{\mu} \nabla_{\mu})^2 \right) \Phi_{\mathcal{E}}. \quad (343)$$

Central to the proposed formalism, the decoupling jerk operator is defined by Eqs. (332) and (338) as

$$\mathcal{J} = -i \frac{G}{\hbar \alpha} \nabla_{\omega} \int dx^{\omega}, \quad (344)$$

while \hat{Z}^{μ} (Eq. (245)) is given by $\hat{Z}^{\mu} = i\hbar\alpha\partial^{\mu}$.

The quantum-gravitational jerking thus takes the explicit form:

$$\begin{aligned} \mathcal{J} \left(-\hbar^2 \hat{Z}^{\mu} \nabla_{\mu} \hat{\mathfrak{D}} \nabla^2 + 2m^2 c^2 \hat{Z}^{\mu} \nabla_{\mu} + m^2 (\hat{Z}^{\mu} \nabla_{\mu})^2 \right) \Phi_{\mathcal{E}} &= \hbar^2 \frac{iG}{\hbar \alpha} (i\hbar\alpha) \nabla_{\omega} \int dx^{\omega} \partial^{\mu} (\nabla_{\mu} \hat{\mathfrak{D}} \nabla^2 \Phi_{\mathcal{E}}) \\ &\quad - i \frac{G}{\hbar \alpha} (i\hbar\alpha) (2m^2 c^2) \nabla_{\omega} \int dx^{\omega} \partial^{\mu} (\nabla_{\mu} \Phi_{\mathcal{E}}) \\ &\quad - i \frac{G}{\hbar \alpha} m^2 \nabla_{\omega} \int dx^{\omega} (Z^{\mu} \nabla_{\mu})^2 \Phi_{\mathcal{E}}. \end{aligned} \quad (345)$$

We first consider the first-order quantum-gravitational jerking, represented by the first two terms of Eq. (345), which depend linearly on the quantum-gravitational conjugate $\hat{Z}^{\mu} \nabla_{\mu}$. Introducing

$$f_{\mu} \equiv \nabla_{\mu} \hat{\mathfrak{D}} \nabla^2 \Phi_{\mathcal{E}}, \quad h_{\mu} \equiv \nabla_{\mu} \Phi_{\mathcal{E}}, \quad \mathbb{K} \equiv (2m^2 c^2 - \hbar^2) G, \quad (346)$$

Eq. (345) can be rewritten as

$$\mathcal{J} \left(-\hbar^2 \hat{Z}^{\mu} \nabla_{\mu} \nabla^2 \Phi_{\mathcal{E}} + 2m^2 c^2 \hat{Z}^{\mu} \nabla_{\mu} \Phi_{\mathcal{E}} \right) = -\hbar^2 G \nabla_{\omega} \int dx^{\omega} \partial^{\mu} f_{\mu} + 2m^2 c^2 G \nabla_{\omega} \int dx^{\omega} \partial^{\mu} h_{\mu}. \quad (347)$$

Applying Eq. (340) yields

$$-\hbar^2 G \nabla_{\omega} \left(g^{\omega\mu} f_{\mu} - \int dx^{\omega} (\partial_{\omega} g^{\omega\mu}) f_{\mu} \right) + 2m^2 c^2 G \nabla_{\omega} \left(g^{\omega\mu} h_{\mu} - \int dx^{\omega} (\partial_{\omega} g^{\omega\mu}) h_{\mu} \right) + \mathbb{K} \nabla_{\omega} b^{\omega}. \quad (348)$$

We reaffirm the fundamental requirement of the quantum-gravitational jerk: the boundary terms must not vanish—they persist as the origin of

the extremisation phase transition, thereby partially fulfilling the decoupling relation of Eq. (333). Physically, this requires that the unified PPBH state evolve significantly, exhibiting infinite resistance to singularity formation, which is equivalent to the field acting non-perturbatively upon itself via an eigen back-reaction during the transition, as established in Theorem 3.

For Eq. (333) to be fully satisfied, the constants b^ω must vanish under covariant differentiation, while the subleading terms $\partial_\omega g^{\omega\mu}$ decay. These \hbar^2 -scaled decaying contributions, denoted \hat{c} , are incorporated into the dynamical Cartan field (Eq. (384)) as signatures of early-universe dynamics. The survival of the boundary terms emerges as the indispensable condition for the decoupling relation and, consequently, for the onset of inflation.

With f_μ and h_μ defined in Eq. (346), the first-order quantum-gravitational jerking follows from its strictly boundary contributions in Eq. (348):

$$\mathcal{J} \left(-\hbar^2 \hat{Z}^\mu \nabla_\mu \hat{\mathcal{D}} \nabla^2 \Phi_\mathcal{E} + 2m^2 c^2 \hat{Z}^\mu \nabla_\mu \Phi_\mathcal{E} \right) = -\hbar^2 G \nabla_\omega \left(g^{\omega\mu} \nabla_\mu \hat{\mathcal{D}} \nabla^2 \right) \Phi_\mathcal{E} + 2m^2 c^2 G \nabla_\omega (g^{\omega\mu} \nabla_\mu) \Phi_\mathcal{E} \quad (349)$$

$$= -\hbar^2 G \nabla_\omega (\nabla^\omega \hat{\mathcal{D}} \nabla^2) \Phi_\mathcal{E} + 2m^2 c^2 G \nabla_\omega \nabla^\omega \Phi_\mathcal{E}. \quad (350)$$

The product rule applied to the first term gives

$$-\hbar^2 G \nabla_\omega (\nabla^\omega \hat{\mathcal{D}} \nabla^2 \Phi_\mathcal{E}) = -\hbar^2 G [(\nabla_\omega \nabla^\omega) \hat{\mathcal{D}} \nabla^2 \Phi_\mathcal{E} + \nabla^\omega (\nabla_\omega \hat{\mathcal{D}}) \nabla^2 \Phi_\mathcal{E} + \nabla^\omega \hat{\mathcal{D}} (\{\nabla_\omega, \nabla^2\} \Phi_\mathcal{E})]. \quad (351)$$

Employing the curved-spacetime d'Alembertian operator [61],

$$\hat{\square} = \nabla^\omega \nabla_\omega, \quad (352)$$

and substituting Eqs. (351) and (352) into Eq. (350) yields

$$-\hbar^2 G [(\hat{\square}) \hat{\mathcal{D}} \nabla^2 + \nabla^\omega (\nabla_\omega \hat{\mathcal{D}}) \nabla^2 + \nabla^\omega \hat{\mathcal{D}} (\{\nabla_\omega, \nabla^2\})] \Phi_\mathcal{E} + 2m^2 c^2 G \hat{\square} \Phi_\mathcal{E}. \quad (353)$$

The curved-spacetime Klein–Gordon equation reads [61]:

$$\hat{\square} \Phi_\mathcal{E} = \frac{m^2 c^2}{\hbar^2} \Phi_\mathcal{E}. \quad (354)$$

Substituting this into Eq. (353) yields

$$-\hbar^2 G [\hat{\square} \hat{\mathcal{D}} \nabla^2 + \nabla^\omega (\nabla_\omega \hat{\mathcal{D}}) \nabla^2 + \nabla^\omega \hat{\mathcal{D}} (\{\nabla_\omega, \nabla^2\})] \Phi_\mathcal{E} + \frac{2Gm^4 c^4}{\hbar^2} \Phi_\mathcal{E}. \quad (355)$$

We now consider the third term in Eq. (345), which governs the second-order quantum-gravitational jerking. Employing Eq. (343) together with the quantum-gravitational force operator $\hat{\mathcal{F}} = \hat{Z}^\mu \nabla_\mu$ from Eq. (321), we obtain:

$$\mathcal{J} \left(m^2 (\hat{Z}^\mu \nabla_\mu)^2 \Phi_\varepsilon \right) = m^2 \mathcal{J} \hat{\mathcal{F}} \hat{\mathcal{F}} \Phi_\varepsilon. \quad (356)$$

Using the identity $\mathcal{J} \hat{\mathcal{F}} = \hat{\mathcal{F}} \mathcal{J} + [\mathcal{J}, \hat{\mathcal{F}}]$ to account for operator noncommutativity, we rewrite Eq. (356) as

$$m^2 \mathcal{J} \hat{\mathcal{F}} \hat{\mathcal{F}} \Phi_\varepsilon = m^2 \hat{\mathcal{F}} \mathcal{J} \hat{\mathcal{F}} \Phi_\varepsilon + m^2 [\mathcal{J}, \hat{\mathcal{F}}] \hat{\mathcal{F}} \Phi_\varepsilon \quad (357)$$

$$= m^2 \hat{Z}^\mu \nabla_\mu \left(-i \frac{G}{\hbar \alpha} \right) (i \hbar \alpha) \nabla_\omega \int dx^\omega \partial^\mu (\nabla_\mu \Phi_\varepsilon) + m^2 [\mathcal{J}, \hat{\mathcal{F}}] \hat{\mathcal{F}} \Phi_\varepsilon \quad (358)$$

$$= i \hbar \alpha G m^2 \hat{q} \Phi_\varepsilon. \quad (359)$$

In Eq. (359), the common factor $i \hbar \alpha G m^2$ has been extracted from both terms, and the operator \hat{q} is defined abstractly so as to incorporate the remaining operator-valued structure of the equation.

Combining Eqs. (355) and (359) therefore gives the expression for the quantum gravitational jerk operation:

$$\begin{aligned} \mathcal{J} \left(-\hbar^2 \hat{Z}^\mu \nabla_\mu \hat{\Delta} \nabla^2 + 2m^2 c^2 \hat{Z}^\mu \nabla_\mu + m^2 (\hat{Z}^\mu \nabla_\mu)^2 \right) \Phi_\varepsilon &= \left(\frac{2Gm^4 c^4}{\hbar^2} \right) \Phi_\varepsilon \\ - \hbar^2 G [\hat{\square} \hat{\Delta} \nabla^2 + \nabla^\omega (\nabla_\omega \hat{\Delta}) \nabla^2 + \nabla^\omega \hat{\Delta} (\{\nabla_\omega, \nabla^2\})] \Phi_\varepsilon &+ i \hbar \alpha G m^2 \hat{q} \Phi_\varepsilon. \end{aligned} \quad (360)$$

Analysis of Eq. (360) reveals that the first two terms are real, while the final term remains imaginary owing to the factor i . The quantum-gravitational jerk thus yields both real and imaginary contributions. Among the real components, the first scales as $1/\hbar^2$, enabling growth away from the quantum limit; it therefore governs gravitational dynamics and possible inflationary mechanisms. The second real component (scaling with \hbar^2) together with the imaginary term, which scales as \hbar^{2+q} (with q arising from \hat{q} and the Planckian mass in Eq. (359) contributing an additional \hbar from its square), encapsulates purely quantum phenomena. During the extremisation phase transition, this separation becomes manifest: gravitational effects emerge into real spacetime, while quantum effects remain confined to the Hilbert space. Employing the definition of the quantum-gravitational geometric operator from

Eq. (341), the quantum-gravitational jerk operation can be expressed as:

$$\mathcal{J} \hat{\mathbb{G}} \Phi_{\mathcal{E}} = (\hat{\beta} + i \hat{\kappa}) \Phi_{\mathcal{E}}, \quad (361)$$

where $\hat{\beta}$ (β is the initial letter of the Greek term βαρύτητα, meaning 'gravity') and $\hat{\kappa}$ (κ is the initial letter of the Greek term κβαντικός, meaning 'quantum'), correspond to the real and imaginary components, respectively. The real component, $\hat{\beta}$, which pertains to gravitational effects, is formulated as:

$$\hat{\beta} \Phi_{\mathcal{E}} \equiv \Re \left[\mathcal{J} \hat{\mathbb{G}} \Phi_{\mathcal{E}} \right] = \left[\frac{2Gm^4 c^4}{\hbar^2} - \hbar^2 G \left[\hat{\square} \hat{\mathcal{D}} \nabla^2 + \nabla^\omega (\nabla_\omega \hat{\mathcal{D}}) \nabla^2 + \nabla^\omega \hat{\mathcal{D}} (\{\nabla_\omega, \nabla^2\}) \right] \right] \Phi_{\mathcal{E}}. \quad (362)$$

Likewise, the imaginary component, $\hat{\kappa}$, which is associated with quantum effects, is given by:

$$\hat{\kappa} \Phi_{\mathcal{E}} \equiv \Im \left[\mathcal{J} \hat{\mathbb{G}} \Phi_{\mathcal{E}} \right] = \hbar \alpha G m^2 \hat{q} \Phi_{\mathcal{E}}. \quad (363)$$

Combining Eqs. (362), (363), and (342) delivers the quantum-gravitational jerk of the primordial extreme-field quiddity:

$$\mathcal{J} \hat{\mathbb{Q}} \Phi_{\mathcal{E}} = \Re \left[\mathcal{J} \hat{\mathbb{G}} \Phi_{\mathcal{E}} \right] + i \Im \left[\mathcal{J} \hat{\mathbb{G}} \Phi_{\mathcal{E}} \right] = (\hat{\beta} + i \hat{\kappa}) \Phi_{\mathcal{E}}. \quad (364)$$

Eq. (364)—the *Equation of Creation*—describes the quantum-gravitational jerk as a key trigger for decoupling gravity from the quantum realm. The second real high-derivative component of $\hat{\beta}$ (scaling with \hbar^2 in Eq. (362)) resides in the gravitational domain, as quantum-gravitational effects persist into the Planckian regime preceding the Jerk Tremendous. Complete decoupling is governed by the attenuation energy scale μ and the departure from Planckian limits. As established in Subsection 4.3 (Eqs. (294)–(296)), the extremisation phase transition drives couplings away from the Planckian limit, attenuating the unified interaction. Decoupling proceeds: the first real term of $\hat{\beta}$ —scaling as $1/\hbar^2$ —becomes stable and non-Planckian. With the PPBH mass being Planckian, $m_{\text{Pl}}^4 = \hbar^2 c^2 / G^2$ (Subsection 4.6, Eq. (367)); devoid of the unified coupling, it persists as a constant. The Planckian real terms (scaling as \hbar^2) and the imaginary component $\hat{\kappa}$ (scaling as \hbar^{2+q}), which retains the unified coupling, vanish. This asymptotic limit thereby achieves the definitive decoupling required for cosmic genesis, now described solely by an intrinsic constant of creation:

$$\lim_{\hbar \rightarrow 0} \left[\lim_{\mu \rightarrow 0} \mathcal{J} \hat{\mathbb{Q}} \Phi_{\mathcal{E}} \right] = \lim_{\hbar \rightarrow 0} \left[\lim_{\mu \rightarrow 0} (\hat{\beta} + i \hat{\kappa}) \Phi_{\mathcal{E}} \right] \approx \frac{2c^6}{G} \Phi_{\mathcal{E}}. \quad (365)$$

4.6 The Tremendous Inflation

Cosmic inflation theory describes an early epoch of exponential expansion that smoothed the geometry of the early universe [69, 100]. Within extreme field theory, inflation emerges from the collective dynamics of $\hat{\beta}$ and $\hat{\kappa}$ in the creation equation (Eq. (364)), derived directly from the quantum-gravitational jerk (Subsection 4.5).

To elucidate the inflationary mechanism, we first examine the constant term in Eq. (362), which embodies purely classical phenomena at the extremisation threshold. Governed by $\hat{\beta}$, this term dominates the classical inflationary epoch and is essential for understanding both the universe's tremendous expansion and subsequent structure formation. This asymptotic dominance follows from Eq. (365). The constant term is defined as:

$$\mathbb{J} \equiv \frac{2Gm^4c^4}{\hbar^2}. \quad (366)$$

This constant \mathbb{J} is henceforth referred to as the constant of creation. At the Planckian extremisation threshold, the mass term m in Eq. (366) is identified with the Planck mass of the primordial Planck black hole, $M_{\text{Pl}} = \sqrt{\frac{\hbar c}{G}} = 2.176 \times 10^{-8}$ kg. To elucidate the physical significance of this fundamental constant, we replace m with the Planck mass M_{Pl} and substitute the standard values: the speed of light $c = 2.99792458 \times 10^8$ m/s, the reduced Planck constant $\hbar = 1.0545718 \times 10^{-34}$ J · s, and the gravitational constant $G = 6.67430 \times 10^{-11}$ m³kg⁻¹s⁻². This substitution yields the following expression:

$$\mathbb{J} = \frac{2GM_{\text{Pl}}^4c^4}{\hbar^2} = \frac{2G(\frac{\hbar c}{G})^2c^4}{\hbar^2} = \frac{2c^6}{G} = \frac{(2) \cdot (2.99792458 \times 10^8 \text{ m/s})^6}{(6.6743 \times 10^{-11} \text{ m}^3\text{kg}^{-1}\text{s}^{-2})} = 2.175 \times 10^{61} \text{ J} \cdot \text{m/s}^2. \quad (367)$$

The result from Eq. (367) has dimensions of [Energy × Acceleration], analogous to Newton's second law $F = ma$. By this correspondence, we identify the derived constant as 'Newton's Universal Extreme Force' (\mathbb{F}_{NUEF}). Normalised by c^2 , it yields a constant with dimensions of mass times acceleration, whose magnitude is 2.4×10^{44} N—exceeding all known physical forces. This colossal force could have driven a primordial acceleration \mathbb{A} of the Planck black hole field immediately following the Jerk Tremendous, making \mathbb{F}_{NUEF} a primary driver of cosmic creation. The force is quantified as:

$$\mathbb{F}_{\text{NUEF}} \equiv \frac{\mathbb{J}}{c^2} = M_{\text{Pl}} \frac{2GM_{\text{Pl}}^3c^4}{\hbar^2} = M_{\text{Pl}} \mathbb{A} = \frac{2c^4}{G} = 2.421 \times 10^{44} \text{ N}. \quad (368)$$

It is instructive to compute the outward pressure corresponding to Newton's Universal Extreme Force, Eq. (368), by dividing it by the Planck area l_{Pl}^2 (where $l_{\text{Pl}}^2 = \hbar G/c^3 = 2.6121 \times 10^{-70} \text{ m}^2$) [83]. This defines Pascal's Universal Extreme Pressure:

$$\mathbb{P}_{\text{PUEP}} = \frac{\mathbb{F}_{\text{NUEF}}}{l_{\text{Pl}}^2} = \frac{2.4 \times 10^{44}}{2.6121 \times 10^{-70}} = 9.266 \times 10^{113} \text{ Pa.} \quad (369)$$

Comparing this value with the Planck pressure,

$$P_{\text{Pl}} = \frac{c^7}{\hbar G^2} = 4.633 \times 10^{113} \text{ Pa,} \quad (370)$$

we find

$$\mathbb{P}_{\text{PUEP}} > P_{\text{Pl}}. \quad (371)$$

This inequality signals inverse gravitational collapse and opposition to singularity formation, in agreement with Theorem 3: Pascal's Universal Extreme (outward) Pressure surpasses the inward Planck pressure of the primordial Planck black hole. At the extreme negentropy threshold, this pressure gradient triggers the Jerk Tremendous—a violent recoil that ejects the system from the Planckian regime into a large-scale, non-Planckian structure.

Building on this foundation, we may reinterpret the constant of creation \mathbb{J} in Eq. (366) at a deeper level. While initially interpreted with dimensions [Energy \times Acceleration], Eq. (367) reveals it can be expressed as $\text{J} \cdot \text{s} \times \text{m/s}^3$ —the product of action (\mathbf{S} , in $\text{J} \cdot \text{s}$) and jerk (the third time derivative of position, in m/s^3 , denoted \mathbf{T} from Τριτοταχύτητα). Thus \mathbb{J} corresponds to a static *jerk tremendous* term in the quantum-gravitational action $\mathbf{S}[\Phi_{\mathcal{E}}]$, encoding the primordial universe's dynamics, and motivates the definition of a new unit:

$$\text{hawking} \equiv [\text{J}\cdot\text{s}] \cdot [\text{m/s}^3]. \quad (372)$$

The constant of creation thereby appears as the static term in the real component of Eq. (364), situated alongside dynamic terms arising from the extreme phase transition in Eqs. (362) and (363). This illuminates a deeper aspect of quantum-gravitational decoupling: the decoupling jerk operator acting on the extreme-field quiddity in Eq. (364) generates both classical and quantum phenomena. This action now corresponds to jerking the quantum-gravitational action $\hat{\mathbf{T}}\mathbf{S}[\Phi_{\mathcal{E}}]$, as modelled by Eq. (373). This interpretation implies that, following an infinitesimal post-jerk interval beyond the

Planckian regime, the action grows as the jerk subsides, eventually satisfying $\mathbf{S} \gg \hbar$ —the hallmark of classicality. With every term in the creation equation bearing the unit [Hawking] (hereafter H), the equation inherently characterises the jerk dynamics of the primordial extreme-field action. This framework consequently captures the primordial universe’s evolution, encoding the quantum-to-classical transition via jerk attenuation and action amplification, expressed as:

$$\mathcal{J} \hat{\mathcal{Q}} \Phi_{\mathcal{E}} \equiv \hat{\mathbf{T}} \mathbf{S}[\Phi_{\mathcal{E}}] = \hat{\mathbf{T}} \int d^4x \mathcal{L}_{\text{QG}}. \quad (373)$$

The quantum-gravitational Lagrangian density \mathcal{L}_{QG} encodes the dynamics of the extreme field, with dependencies on $\Phi_{\mathcal{E}}$, its derivatives, curvature terms, and coupling constants.

From this analysis, the constant of creation emerges as the static term arising from the jerking of the quantum-gravitational action during the extremisation phase transition:

$$\mathbf{T} \int d^4x \mathcal{L}_{\text{QG}} \Big|_{\text{static}} \equiv \mathbb{J}. \quad (374)$$

This constant—henceforth interpreted as the quantum-gravitational jerk constant—encapsulates the fundamental physics underlying all constants of creation. Consequently, all basic constants of the universe can be collectively derived from \mathbb{J} as follows:

$$\text{The Quantum-Gravitational Jerk Constant} \equiv \mathbb{J} = 2.175 \times 10^{61} \text{ H}, \quad (375)$$

$$\text{Power of Creation} \equiv \mathbb{W} = \frac{\mathbb{J}}{c} = 7.257 \times 10^{52} \text{ W}, \quad (376)$$

$$\text{Force of Creation} \equiv \mathbb{F}_{\text{NUEF}} = \frac{\mathbb{J}}{c^2} = 2.421 \times 10^{44} \text{ N}, \quad (377)$$

$$\text{Pressure of Creation} \equiv \mathbb{P}_{\text{PUEP}} = \frac{\mathbb{J}}{c^2 l_{\text{Pl}}^2} = 9.266 \times 10^{113} \text{ Pa}, \quad (378)$$

$$\text{Primordial Extreme Acceleration} \equiv \mathbb{A} = \frac{\mathbb{J}}{M_{\text{Pl}} c^2} = 1.112 \times 10^{52} \text{ m/s}^2, \quad (379)$$

$$\text{Primordial Extreme Density} \equiv \mathbb{D} = \frac{\mathbb{J}}{\mathbb{A} c^2 l_{\text{Pl}}^3} = 5.155 \times 10^{96} \text{ kg/m}^3, \quad (380)$$

$$\text{Primordial Extreme Temperature} \equiv \mathbb{T} = \frac{\mathbb{J}}{\mathbb{A} k_B} = 1.417 \times 10^{32} \text{ K}. \quad (381)$$

Eqs. (375)–(381) together encapsulate the cosmic creation framework of extreme field theory, according to which the Universe emerged from a primordial Planck black hole (PPBH) that underwent a quantum-gravitational jerk of approximately 2.175×10^{61} H. At the extremisation threshold, the PPBH exhibited a primordial density of 5.155×10^{96} kg/m³, emitted an immense power of 7.257×10^{52} W—precisely double the Planck power $W_{\text{Pl}} = c^5/G = 3.628 \times 10^{52}$ W—and attained a temperature of 1.417×10^{32} K.

This extreme configuration was driven by Newton’s Universal Extreme Force, quantified as 2.421×10^{44} N—precisely twice the Planck force $F_{\text{Pl}} = c^4/G = 1.210 \times 10^{44}$ N. This force generated Pascal’s Universal Extreme Pressure of 9.266×10^{113} Pa, double the Planck pressure $P_{\text{Pl}} = c^7/(\hbar G^2) = 4.633 \times 10^{113}$ Pa, and induced a primordial extreme acceleration of 1.112×10^{52} m/s², twice the Planck acceleration $a_{\text{Pl}} = \sqrt{c^7/(\hbar G)} = 5.561 \times 10^{51}$ m/s². These colossal Planckian quantities were fundamental to the creation of our universe, culminating in the Jerk Tremendous that brought the universe into existence [83, 118, 128].

We now analyse Newton’s Universal Extreme Force, \mathbb{F}_{NUEF} , defined in Eq. (368). As a result of quantum gravitational jerking, this force emerges as a fundamental constant. The constancy of \mathbb{F}_{NUEF} requires that as the primordial extreme acceleration decays, mass-energy grows inversely to conserve the force:

$$\mathbb{F}_{\text{NUEF}} = e^{H(t)} M_{\text{Pl}} \mathbb{A} e^{-H(t)}. \quad (382)$$

Eq. (382) indicates that the primordial extreme acceleration decays exponentially, $\mathbb{A}(t) \sim \mathbb{A} e^{-H(t)}$, while the corresponding mass-energy grows exponentially, $M(t) \sim e^{H(t)} M_{\text{Pl}}$, ensuring that the force remains constant. We reserve a detailed analysis of the evolution function $H(t)$ for future investigation.

An alternative representation of Newton’s Universal Extreme Force is given by

$$\mathbb{F}_{\text{NUEF}} = \mathfrak{F}^+ + \mathfrak{F}^-, \quad (383)$$

where \mathfrak{F}^- may be associated with dark energy and the resulting accelerated expansion of the universe. In this interpretation, a negative force density manifests as negative pressure in the stress–energy tensor, thereby producing a repulsive gravitational contribution. Dark energy can thus be viewed as a field whose intrinsic negative pressure acts as a cosmological-scale repulsive force, driving the universe’s acceleration. Both formulations, Eqs. (382) and (383), will be examined in greater detail in future work.

Having established the constants of creation as emerging from the static contribution of the equation of creation (Eq. (374)), we now consider its dynamical counterpart. These arise from Planckian-scaled components generated by the quantum-gravitational jerk operator (Eq. (360)): the second high-derivative term scales with \hbar^2 (more generally \hbar^{2+d} for some d that may arise from derivative orders) and incorporates the postulated decaying terms $\hat{\omega}$ from Eq. (348), together with a generalised noncommutativity correction $\hat{\sigma}$ for \hat{p}^2 and \hat{C}^2 from Eq. (303) (both scaling with \hbar^2). Meanwhile, the purely imaginary component $\hat{\kappa}$ from Eq. (363) scales with \hbar (more precisely \hbar^{2+q} where q originates from the operator \hat{q} in Eq. (359) and the Planck mass squared contributes an additional \hbar). Acting collectively on $\Phi_{\mathcal{E}}$, these contributions are jointly defined as the Dynamical Cartan Field (DCF):

$$\hat{D} \Phi_{\mathcal{E}} \equiv \left[-\hbar^2 G [\hat{\square} \hat{\mathcal{D}} \nabla^2 + \nabla^\omega (\nabla_\omega \hat{\mathcal{D}}) \nabla^2 + \nabla^\omega \hat{\mathcal{D}} (\{\nabla_\omega, \nabla^2\}) + \hat{\omega} + \hat{\sigma}] + i \hbar \alpha G m^2 \hat{q} \right] \Phi_{\mathcal{E}}. \quad (384)$$

The dynamical Cartan field \hat{D} decomposes into real and imaginary components. The real \hbar^2 -scaled part consists of higher-derivative dynamical terms: $\hat{\square}$, $\hat{\mathcal{D}}$ (spatial metric), ∇^2 , ∇_ω , and an abstract anticommutator term $((\nabla_\omega \nabla^2) \Phi_{\mathcal{E}} + \nabla^2 (\nabla_\omega \Phi_{\mathcal{E}}))$ arising from the product rule in Eq. (351). These are accompanied by the decaying-geometry operator $\hat{\omega}$, which captures metric variations $\partial_\omega g^{\omega\mu}$ from the first-order quantum-gravitational jerk (Eq. (348)) that decay as the metric approaches constancy.

The correction $\hat{\sigma}$ originates from the noncommutativity of \hat{p}^2 and \hat{C}^2 (Eq. (303)). If these operators commuted, $\hat{\sigma}$ would vanish; their noncommutativity thus forces its introduction. To preserve conventional ordering, we define

$$\hbar^2 \hat{\sigma} \equiv \mathcal{J} \left[\hat{j}^{ii} \hat{\mathbf{p}}^2, \hat{C}^2 \right] = \mathcal{J} \left[\hat{\mathcal{D}} \hat{\mathbf{p}}^2, \hat{C}^2 \right]. \quad (385)$$

The imaginary component is exactly the operator $\hat{\kappa}$ from Eq. (363), corresponding to the imaginary part of the equation of creation (Eq. (361)). Acting collectively on $\Phi_{\mathcal{E}}$, these elements together constitute \hat{D} .

By construction, the dynamical Cartan field \hat{D} in Eq. (384) encodes purely quantum effects and must therefore vanish in the classical limit, acting as an annihilation operator on the extreme field:

$$\lim_{\hbar \rightarrow 0} \hat{D} \Phi_{\mathcal{E}} = 0. \quad (386)$$

In the Planckian inflationary epoch following the Jerk Tremendous, however, \hat{D} becomes significant. Its higher-derivative terms—fundamental at

this scale [89, 114]—together with its purely imaginary damping component [34, 35], govern the field’s dynamics. This mechanism smooths spacetime geometry, contributes to renormalisation, and counterbalances the rapid acceleration driven by Newton’s Universal Extreme Force (NUEF, Eq. (377)). Consequently, it stabilises the exponential decay of the primordial acceleration \mathbb{A} (Eq. (379)) and the inverse growth of energy and mass (Eq. (382)), thereby regulating cosmic expansion and enabling the transition to weak-field physics and large-scale structure formation. A comprehensive analysis of the full dynamics of $\hat{\mathbb{D}}$ is reserved for forthcoming work.

Another aspect of the DCF is its conjectured non-torsion-free nature. Should the field violate the lower-index symmetry of the Christoffel symbols during the inflationary epoch, such that

$$\Gamma_{\mu\nu}^{\lambda} \neq \Gamma_{\nu\mu}^{\lambda}, \quad (387)$$

then this asymmetry indicates both the presence of curvature and an interaction with spin. The torsion tensor, which encodes this interaction, is defined by the antisymmetric part of the connection:

$$T_{\mu\nu}^{\lambda} = \Gamma_{\mu\nu}^{\lambda} - \Gamma_{\nu\mu}^{\lambda}. \quad (388)$$

Within this framework, the DCF integrates both spin and curvature, consistent with the principles of Einstein–Cartan theory [165–167], which extends General Relativity by incorporating torsion arising from the intrinsic spin of matter. During inflation, the DCF is conjectured to have exhibited a pronounced asymmetry—coupling preferentially to matter while partially decoupling from antimatter—thereby contributing to a secondary mechanism for the observed matter–antimatter imbalance. The primary mechanism, established in Subsection 3.6, arises from non-perturbative extreme-field asymmetric tunnelling probabilities: matter tunnelling from the non-perturbative primordial black hole interior is attenuated as eigen-Hawking radiation, while antimatter tunnelling is enhanced as anti-eigen flux. Together, these two mechanisms offer a complementary resolution to the matter–antimatter asymmetry within the extreme-field framework.

The dynamical Cartan field (Eq. (384)) and the constant of creation (Eq. (366)) thus unify the quantum-gravitational jerk with the subsequent inflationary epoch. Their complementary mechanisms—the jerk driving the emergence of the universe from the Planck scale, while damping stabilises expansion—smooth spacetime, enable classicality, ensure matter persistence, and guide cosmic evolution.

4.7 Ansatz for an Asymptotically Free Gravitational Universe

Extreme field theory concludes that the universe originated from a *Jerk Tremendous* of an extremely negentropic primordial Planck black hole. This framework, encapsulated in Eqs. (375)–(381), is consistent with observational evidence for a hot, dense Planckian origin followed by rapid expansion and cooling [40, 96, 126, 139].

This framework defines a *Planckian-Fixed Universe*, in which the decoupling jerk operator exclusively decouples the universal gravitational constant G (Eq. (333)) and annihilates quantum operators. While this elegantly describes cosmic emergence, the initial decoupling jerk operator (Eqs. (332) and (338)) admits a complementary refinement—a partial decoupling of both G and $\alpha_g(\mu)$ —thereby achieving a unique separation of gravitational and quantum effects. For this analysis, the operator must satisfy three conditions: it must be third-order in derivatives; it must attenuate the quantum coupling while inversely amplifying the gravitational coupling so as to preserve the constants of creation; and it must uniquely invert the quantum component of the gravitational conjugate \hat{Z}^μ up to the covariant derivative. This derivative completes the conjugate, and together they generate the d'Alembert operator—the mass-creating operator. We therefore posit the following ansatz for the decoupling jerk operator:

$$\mathcal{J} = \left(\frac{-ic}{\hbar} \right) \frac{\partial \ln}{\partial \alpha_g(\mu)} \frac{\partial}{\partial \alpha_q(\mu)} \nabla_\omega \int dx^\omega. \quad (389)$$

While this structure could generate additional terms beyond those in Eq. (360) through coupling derivatives acting on the field itself, in this subsection we focus exclusively on the first classical static term—the constant of creation defined in Eq. (366).

The physical justification for the decoupling jerk operator lies in its action on the quantum-gravitational conjugate $i\hbar\alpha(\mu)\partial^\mu\nabla_\mu$ (Eq. (321)). This action decomposes the unified coupling $\alpha(\mu) \equiv \frac{G}{c}\alpha_g(\mu)\alpha_q(\mu)$ into distinct gravitational (α_g) and quantum (α_q) components. The operator's primary role is therefore to disentangle these phenomena, isolating gravitational interactions as:

$$\frac{\partial \alpha(\mu)}{\partial \alpha_q(\mu)} = \frac{G}{c}\alpha_g(\mu). \quad (390)$$

Furthermore, gravitational phenomena must exhibit asymptotic freedom. By analogy with QCD—where quarks become free in the ultraviolet—gravity must become free in the infrared, while preserving the constants of creation. This requires:

$$\frac{G}{c} \frac{\partial \ln \alpha_g(\mu)}{\partial \alpha_g(\mu)} = \frac{G}{c} \frac{1}{\alpha_g(\mu)} \sim \frac{G}{c} \frac{1}{\beta(\alpha_g)}. \quad (391)$$

In quantum field theory, the beta function (Gell-Mann–Low function) for the gravitational coupling is defined as

$$\beta(\alpha_g) \equiv \frac{\partial \alpha_g(\mu)}{\partial \ln \mu}. \quad (392)$$

Originally formulated by Gell–Mann and Low [63] and widely used in quantum field theory [49, 127], this definition has been extended to gravitational interactions within the asymptotic safety framework [137, 158]. Since $\alpha_g(\mu)$ depends on the energy scale μ , one can define—up to a constant Planckian factor \mathbf{P} —the relation:

$$\frac{\partial \ln \alpha_g(\mu)}{\partial \alpha_g(\mu)} = \mathbf{P} \beta(\alpha_g)^{-1}. \quad (393)$$

To complete the physical justification of the decoupling jerk operator (Eq. (389)), we recall that at the extremisation phase transition, quantum fluctuations extremised and converged to a continuum-like state. Mathematically, this is represented by the action of the integral operator $\int dx^\omega$ on the quantum component of the conjugate \hat{Z}^μ , where integration over spacetime embodies the transition of quantum fluctuations out of the quantum regime, averaging their behaviour into a classical description that permits the negentropic Primordial Planck Black Hole (PPBH) to converge into a single, intrinsically unstable state with extreme resistance to singularity formation. This yields a natural response we term the *non-perturbative eigen back-reaction*, originating from the self-action of the non-inverted covariant derivative and reflecting its geometric invariance during the transition. At the Planckian threshold, the covariant derivative encoded the emerging spacetime curvature, now a classical entity. As quantum interactions ceased, the extreme field assumed a classical PPBH state—analogueous to a quantum Planckian particle embedded in a curved spacetime. The non-perturbative eigen back-reaction thus marked the precursor to the PPBH recoil, driven by the mass-generation mechanism of the Klein–Gordon equation, where the operator $\nabla^\mu \nabla_\mu$ acting

on the scalar field $\Phi_{\mathcal{E}}$ (Eq. (352)) yields mass (Eq. (354)) as a manifestation of this process.

Therefore, the action of the ansatz decoupling jerk operator in Eq. (389) on the quantum operator \hat{Z}^ν , as introduced in Eq. (333), is modified by the inverse of the beta function, according to Eqs. (389), (391), (392), and (393):

$$\hat{A}^\mu \hat{Z}^\nu = \frac{\mathbb{P} G}{\beta(\alpha_g)} g^{\mu\nu}. \quad (394)$$

Applying Eq. (394) to Eq. (345) shows the universe emerged with evolving, not fixed, constants. For a primordial Planck black hole, the creation constants in Eqs. (375)–(379) scale inversely with the gravitational beta function:

$$\mathbb{J}(\mu) \sim \frac{\mathbb{J}}{\beta(\alpha_g)}, \quad \mathbb{W}(\mu) \sim \frac{\mathbb{W}}{\beta(\alpha_g)}, \quad \mathbb{F}(\mu) \sim \frac{\mathbb{F}_{\text{NEUF}}}{\beta(\alpha_g)}, \quad \mathbb{P}(\mu) \sim \frac{\mathbb{P}_{\text{PUEP}}}{\beta(\alpha_g)}, \quad \mathbb{A}(\mu) \sim \frac{\mathbb{A}}{\beta(\alpha_g)}. \quad (395)$$

The infrared attenuation $\beta(\alpha_g) \rightarrow 0$ defines an *Asymptotically Free Gravitational Universe (AFGU)*, in which gravity completely decouples from quantum fluctuations at large scales—explaining both its apparent weakness in the quantum regime and its dominance in classical cosmology.

Moreover, the second-order quantum-gravitational term—defined by the action of the ansatz decoupling jerk operator in Eq. (389) on the third term in Eq. (345), which is of order $\alpha(\mu)^2$ —persists as $\mu \rightarrow 0$ because it yields $\alpha_q(\mu)/\alpha_g(\mu) \sim \beta(\alpha_g)^{-1}$. This ratio runs dynamically in the infrared and does not vanish, as shown in Eq. (365).

Furthermore, the AFGU framework may provide a natural resolution to the cosmological constant problem [106, 159]. The problem stems from the enormous discrepancy between the observed vacuum energy density and the value predicted by quantum field theory. Observations yield

$$\rho_\Lambda^{\text{obs}} \simeq (10^{-3} \text{ eV})^4, \quad (396)$$

while QFT vacuum fluctuations up to a UV cutoff scale Λ_{UV} yield

$$\rho_{\text{vac}}^{\text{QFT}} \simeq \frac{1}{16\pi^2} \int^{\Lambda_{\text{UV}}} k^3 dk \sim \Lambda_{\text{UV}}^4, \quad (397)$$

which for a Planck-scale cutoff produces a discrepancy of order 10^{122} .

In an AFGU, this discrepancy may be naturally alleviated. According to Eq. (395), the fundamental constants of creation evolve dynamically, scaling

inversely with the gravitational beta function:

$$\mathbb{C}(\mu) \propto \frac{1}{\beta(\alpha_g)}. \quad (398)$$

This behaviour becomes particularly significant when combined with the dynamical acceleration that redefines Newton’s Universal Extreme Force, \mathbb{F}_{NUEF} , in Eq. (382). This relation shows that the running acceleration dramatically suppresses the observed local energy scale via the tiny gravitational beta function $\beta(\alpha_g)$, offering a natural mechanism to neutralise Planck-scale vacuum energy at cosmological scales:

$$\mathbb{A}(\mu) \sim \frac{\mathbb{A}}{\beta(\alpha_g)} \sim \frac{\mathbb{F}_{\text{NUEF}}}{\beta(\alpha_g)M} \sim \frac{1}{\beta(\alpha_g)E}. \quad (399)$$

Therefore, in an AFGU, there is no true Planckian cutoff. As indicated by Eqs. (373) and (395), continuous gravitational jerks—remnants of the Jerk Tremendous event—prevent the ultraviolet scale from reaching a strict Planckian limit. Instead, this scale becomes gravitationally attenuated to an AFGU effective cutoff:

$$\Lambda_{\text{eff}}^{\text{AFGU}} \sim \Lambda_{\text{UV}} \int \int P^{-1} \beta(\alpha_g) d\phi^b d\psi^f, \quad (400)$$

where the gravitational beta function has been integrated over all bosonic and fermionic fields in the universe. This yields an averaged effective cutoff that is expected to be far below the Planck scale. Consequently, the corresponding effective QFT vacuum energy—analogueous to Eq. (397)—scales as the quartic power of this attenuated cutoff:

$$\rho_{\text{vac}}^{\text{eff}} \sim (\Lambda_{\text{eff}}^{\text{AFGU}})^4. \quad (401)$$

Under the ansatz of an AFGU, the contribution $\int \int P^{-1} \beta(\alpha_g) d\phi^b d\psi^f$ is parametrically suppressed. Consequently, its effect on the vacuum sector becomes negligible, allowing the effective vacuum energy density to dynamically relax toward the observed value given in Eq. (396).

This suppression stabilises the vacuum energy at Planckian scales, circumventing the large energy densities associated with microscopic volumes while preserving translational symmetry. Crucially, it explains why cosmic emergence in the non-perturbative extreme field theory arises uniquely from the Primordial Planck Black Hole (PPBH), a process prohibited in the asymptotically free gravitational regime. Asymptotic gravitational freedom thus provides a unified framework that simultaneously addresses inflation, late-time acceleration, and the cosmological constant problem.

5 Conclusion

The central achievement of this work is Extreme Field Theory—an effective operator-based framework describing quantum-geometric corrections in black hole interiors—wherein the metric and speed of light acquire operator-valued structure, quantising geometry at high curvature and revealing an intrinsic extreme-field noncommutativity. This induces a highly non-trivial quantum spacetime geometry, fundamentally departing from the classical picture and laying the groundwork for addressing the singularity problem. Recursive feedback between the metric, its inverse, and the connection self-stabilises the quantum differential structure, suggesting a seesaw-like mechanism that dynamically prevents singularities.

Turning to the primordial universe, the theory predicts that primordial black holes can arise acausally, preceding the causal pre-being fluctuations traditionally assumed to generate them—provided such fluctuations follow a nonlinear (non-Gaussian) distribution. This non-causal ordering thus provides the spacetime arena necessary for conventional fluctuations to occur.

In the perturbative regime—applicable both to acausal primordial black holes and to contemporary astrophysical black holes—the quantum metric exhibits a probabilistic quantum-geometric behaviour that preserves the underlying spacetime structure. This feature allows for a consistent resolution of the information paradox, as quantum-corrected escapable horizons, distinct from their classical counterparts, emerge and permit a non-vanishing tunnelling probability for acausal anti-eigen flux excitations from the black hole interior. In this regime, the tunnelling probabilities of standard Hawking radiation and anti-eigen flux are symmetric, thereby ensuring information preservation and maintaining unitarity.

A primordial black hole saturating the lower boundary of non-perturbative quantisation provides a sufficient initial condition for cosmic genesis. In this non-perturbative phase, an intrinsic spacetime asymmetry emerges, such that the creation rate of interior eigen-Hawking radiation exceeds that of standard Hawking radiation. This imbalance drives the system along a negentropic trajectory toward an increasingly ordered Planck-scale configuration, identified as the *Primordial Planck Black Hole*. Along this trajectory, an internal asymmetry emerges: the extreme-field eigen-intrinsic asymmetry parameter favours matter confinement, while antimatter deconfines via anti-eigen flux tunnelling, predicting an accumulation of antimatter beyond the observable horizon.

At the Planckian extremisation threshold, tunnelling probabilities are suppressed as quantum fluctuations extremise and discretisation effects converge to a continuum, driving the system towards a deterministic, negentropic, and intrinsically unstable state. This configuration consequently necessitates an extremisation-induced phase transition leading to macroscopic structure formation. At this critical juncture, a quantum-gravitational jerk—termed the *Jerk Tremendous*—occurs within the PPBH core, manifesting as a violent recoil from the extremal Planckian configuration. This jerking mechanism represents the quantum-gravitational disconjugation of the unified coupling, signalling gravity’s decoupling from the quantum regime, the onset of inflation, and the birth of the Universe as an asymptotically free gravitational system, naturally resolving the cosmological constant problem.

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The author declares no conflict of interest.

Author Contributions

The author is solely responsible for all aspects of this work, including conceptualisation, theoretical development, analysis, and manuscript preparation.

Data Availability

No datasets were generated or analysed during the current study; therefore, data sharing is not applicable.

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