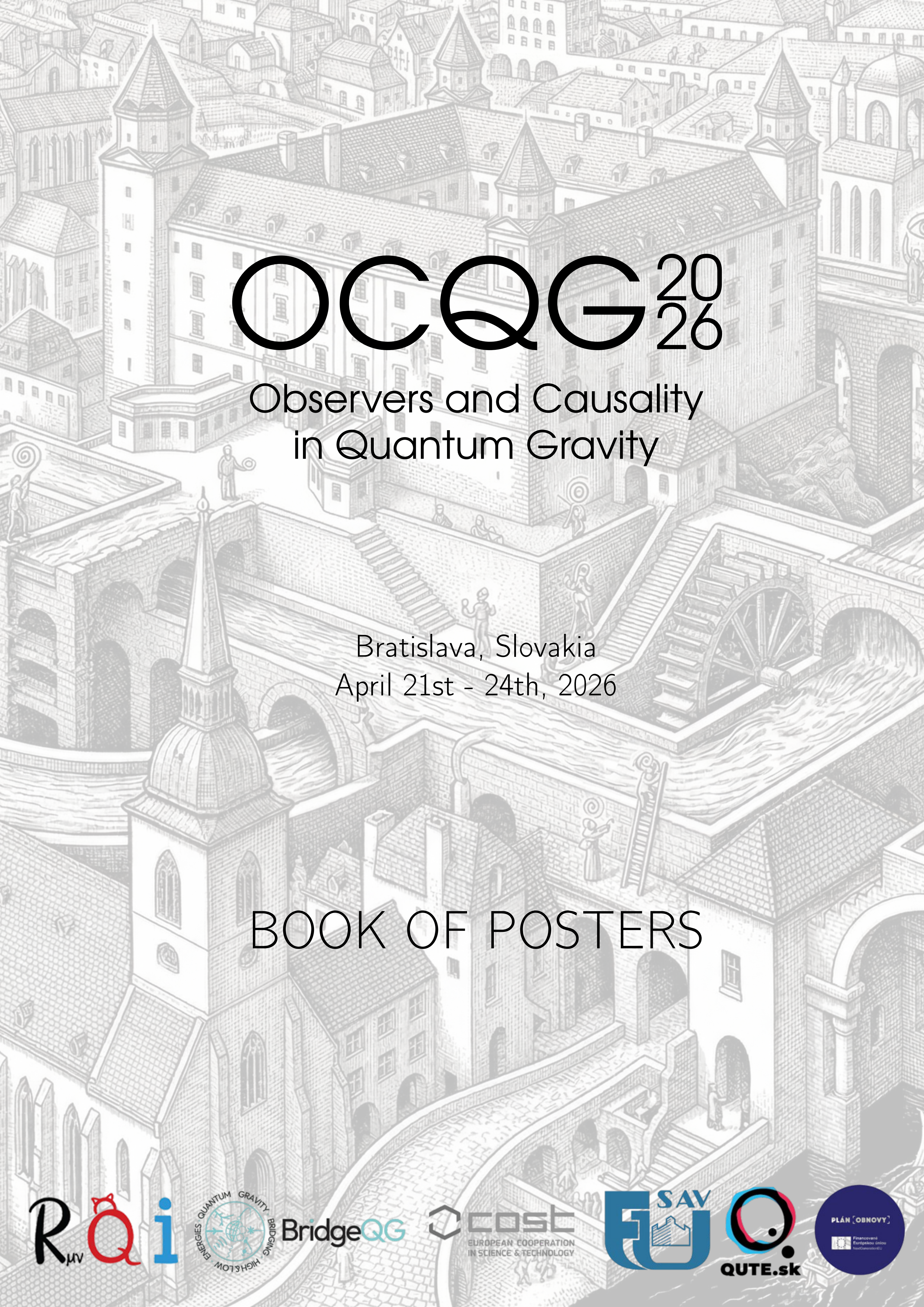

OCQG 2026

Observers and Causality in Quantum Gravity

Bratislava, Slovakia

April 21st - 24th, 2026

BOOK OF POSTERS

Contents

Posters

Alberto Spalvieri

Local Operations and Field Mediated Entanglement without a Local Tensor Product Structure

Local Operations and Field-Mediated Entanglement without Local Tensor Product Structure

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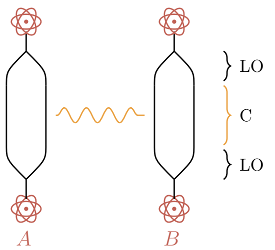
February 2, 2026

Abstract

Quantum information tools rely on subsystem locality, which is obstructed in gauge theories by the absence of a spacetime-local tensor product structure. This limits the direct applicability of results such as the LOCC theorem. We address this in a two-dimensional lattice gauge model of electromagnetism by constructing gauge-invariant local algebras and an operationally meaningful Hilbert space decomposition. Applying this framework to field-mediated entanglement protocols, we show that a discrete analogue of the LOCC theorem holds: entanglement requires genuinely quantum field interactions. This points toward an operational notion of subsystems in gauge theories.

BMV: Field-Mediated Entanglement

Protocol: LO (path splitting and recombination) + C (field-mediated evolution). By LOCC theorem [1], LO + classical C cannot create entanglement. If entanglement is observed, the mediator is non-classical [2, 3].



Gauge-Theory Obstruction

In QED, Gauss law becomes an operator constraint on physical states:

$$\boldsymbol{\nabla}\cdot\hat{\mathbf{E}}(\mathbf{x})\,|\Psi\rangle_{\text{phy}}=0\quad\forall\,\mathbf{x}.$$

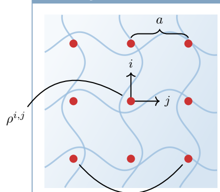
Thus $\mathcal{H}_{\text{phy}}$ is not a simple tensor product over space points:

$$\mathcal{H}=\cdots\otimes\;\underbrace{\mathcal{H}^E_{\mathbf{x}-\mathbf{h}}\;\otimes\;\mathcal{H}^E_{\mathbf{x}}}_{\nabla}\;\otimes\;\underbrace{\phantom{\mathcal{H}}}_{\nabla}\mathcal{H}^E_{\mathbf{x}+\mathbf{h}}\;\otimes\Rightarrow.\quad\mathcal{H}_{\text{phy}}\neq\bigotimes_{\mathbf{x}}\mathcal{H}^E_{\mathbf{x}}$$

so “local ops” and “entanglement” need a refined, operational definition.

1. **The LOCC framework and the corresponding theorem cannot be applied**
2. **The mechanism of generating matter superposition and measuring entanglement is unclear**

2D Toy Lattice Model



$$\bar{\partial}_x f^{\,i,j}=\frac{1}{2a}(f^{\,i,j+1}-f^{\,i,j-1})$$

Lattice and discrete derivatives.

Electromagnetism	Toy model
$\mathbf{A}(\mathbf{x})$	$\mathbf{q}^{i,j}=(q_x^{i,j},q_y^{i,j})$
$-\mathbf{E}(\mathbf{x})$	$\mathbf{p}^{i,j}=(p_x^{i,j},p_y^{i,j})$
$\mathcal{G}:=\nabla\cdot\mathbf{E}(\mathbf{x})$	$\mathcal{C}^{i,j}=\bar{\partial}_x p_x^{i,j}+\bar{\partial}_y p_y^{i,j}$
$\mathbf{B}(\mathbf{x})=\nabla\times\mathbf{A}(\mathbf{x})$	$b^{i,j}=\bar{\partial}_x q_y^{i,j}-\bar{\partial}_y q_x^{i,j}$
$H=\frac{1}{2}\int d\mathbf{x}(\lvert\mathbf{E}(\mathbf{x})\rvert^2+\lvert\mathbf{B}(\mathbf{x})\rvert^2)$	$H=\frac{1}{2}\sum_{i,j}[(p_x^{i,j})^2+(p_y^{i,j})^2+(b^{i,j})^2]$

$$H_\rho=H,\qquad \mathcal{C}_\rho^{i,j}:=\mathcal{C}^{i,j}+\rho^{i,j}=0$$

Quantization

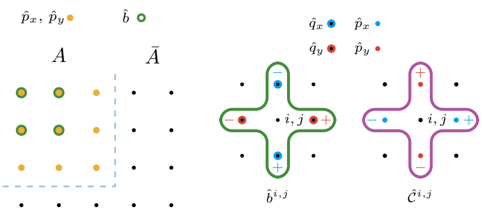
Canonical commutators: $[\hat{q}_s^{i,j},\hat{p}_r^{n,m}]=i\hbar\,\delta_{sr}\,\delta^{i,n}\delta^{j,m}$.

Gauge Constraint: $\left(\bar{\partial}\cdot\hat{\mathbf{p}}^{i,j}+\hat{\rho}^{i,j}\right)|\Psi\rangle_{\text{phy}}=0,\qquad \hat{\rho}^{i,j}=|1\rangle\langle 1|_{i,j}\,.$

Vacuum solution for a generic semiclassical source state $\sum_s c(s)\,|s\rangle$ with $|s\rangle:=\bigotimes_{i,j}|0\vee 1\rangle_{i,j}$, is a coherent superposition of shifted Gaussians [4]:

$$|\Upsilon\rangle=\eta\sum_s\alpha_s\,|s\rangle^m\,|\Psi_{0,\rho(s)}\rangle^f\,.$$

Local Algebra on a Region A



Generators

With matter: add the $\hat{\rho}^{i,j}$'s

Local Algebra of A $\rightarrow\quad \mathcal{A}_A:\{\hat{O}\,|\,[\hat{O},\hat{\mathcal{C}}^{i,j}]=0,\ \forall(i,j)\}$

The Local Decomposition

Method based on Zanardi et al. construction [5] applied to gauge theories [6, 7]

Diagonalize $\mathcal{Z}_A$

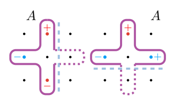
Diagonalize $\{\hat{\mathcal{C}}_\rho^{i,j}\}$

Impose $\hat{\mathcal{C}}_\rho^{i,j}\,|\psi\rangle_{\text{phy}}=0$

$$\mathcal{A}_A=\bigoplus_k\mathcal{O}_A^k\otimes\mathbb{I}_{\bar{A}}^k$$

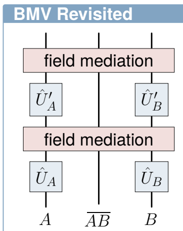
$$\bar{\mathcal{A}}_A=\bigoplus_k\mathbb{I}_A^k\otimes\mathcal{O}_{\bar{A}}^k$$

Edge Terms



- **Operational Decomposition:** $\mathcal{H}_{\text{phy}}=\bigoplus_k\mathcal{H}_A^k\otimes\mathcal{H}_{\bar{A}}^k$
- **Split Decomposition:** $\mathcal{H}_{\text{phy}}=\bigoplus_s[|s\rangle^m\otimes(\bigoplus_k\mathcal{H}_{k,A}^s\otimes\mathcal{H}_{k,\bar{A}}^s)]$

BMV Revisited



- LO: $\hat{U}_A,\,\hat{U}'_A\in\mathcal{A}_A,\quad \hat{U}_B,\,\hat{U}'_B\in\mathcal{A}_B.$
- C: field evolution and relaxation

$$\mathcal{H}_{\text{phy}}=\bigoplus_k\boxed{\mathcal{H}_A^k\otimes\mathcal{H}_B^k}\otimes\mathcal{H}_{\overline{AB}}^k$$

Measure sector-wise entanglement

↓

Non-classical field

Conclusions

- **Operational Decomposition**: sector-wise tensor products enable **LOCC without a global local factorization**.
- **Split Decomposition**: explicit mechanism for **spatial superpositions** and **matter–field–induced entanglement**.
- **Next steps**: delocalized sources; finite-resolution continuum; links to edge modes/QRFs; towards linearized gravity.

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Aleksandra Gočanin

Spacetime Tomography: Probing Spacetime Geometry with Entangled Particles

Peres time correlations as operational probes of spacetime curvature

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Abstract

Building on the operational framework developed by Perche [Phys. Rev. D 106, 025018 (2022)], we study two localized nonrelativistic quantum particles propagating along timelike geodesics in curved spacetime background. For each particle, we consider the Peres sojourn time accumulated in a spatial region transported along the reference geodesic and calculate the covariance of Peres times for separable and entangled two-particle states, making a comparison between flat and curved spacetime background. We then reformulate the protocol as a Bell-type experiment based on local, time-integrated measurements, and show that the violation of the CHSH inequality arises specifically when entangled states and gravity are involved together, the CHSH parameter acquiring a curvature-dependent correction. As an explicit example, we analyze AdS spacetime, where the local tidal term produces a confining quadratic potential and hence a positive correction to the CHSH parameter. A protocol calibrated to saturate the classical bound in flat spacetime is therefore driven above it in AdS spacetime for entangled states. Our results demonstrate how spacetime geometry can influence observable quantum correlations in a fully operational setting.

Motivation

- A central question in modern fundamental physics is how quantum mechanics and gravity can be brought into a common operational framework. In particular, one would like to understand whether spacetime geometry can leave observable signatures in genuinely quantum correlations.
- Recent developments in relativistic quantum information have shifted attention toward localized, operationally defined quantum systems such as wave packets, detectors, clocks, and quantum reference frames, which provide a natural arena for probing curved spacetime.
- In this work we study two localized, non-interacting particles propagating along timelike geodesics in curved spacetime and use the **Peres sojourn time** as our basic local observable.
- Our starting question is whether entanglement can leave an operational imprint even when the one-particle reduced dynamics is identical to that of a corresponding classical mixture.
- We show that the distinction appears at the level of **two-particle temporal correlations**, and that these correlations are also sensitive to the background geometry.
- This leads naturally to a Bell-type reformulation, where curvature modifies the effective local dynamics and turns the CHSH parameter into an **operational curvature witness**.

Framework: localized particles in curved spacetime

For a sufficiently localized wave packet, one introduces Fermi normal coordinates (τ, x^i) around a reference timelike worldline Γ. The effective one-particle Hamiltonian is

$$\hat{H} = m\mathbb{I} + \frac{\hat{p}^2}{2m} + ma_i(\tau)\hat{x}^i + \frac{m}{2}R_{0i0j}(\tau)\hat{x}^i\hat{x}^j + \cdots$$

where curvature enters through the tidal term R_{0i0j}.
For a transported spatial region $\mathcal{R} \subset \Sigma_\tau$, the Peres-time operator is

$$\hat{T}_{\mathcal{R}} = \int d\tau\, \hat{\Pi}_{\mathcal{R}}(\tau), \qquad \hat{\Pi}_{\mathcal{R}}(\tau) = \int_{\mathcal{R}} d\Sigma\, |x,\tau\rangle\langle x,\tau| .$$

Its expectation value gives the sojourn time,

$$T_{\mathrm{P}} = \langle\psi_{\mathrm{in}}|\hat{T}_{\mathcal{R}}|\psi_{\mathrm{in}}\rangle = \int d\tau \int_{\mathcal{R}} d\Sigma\, |\psi(x,\tau)|^2.$$

Two-particle state and covariance

We consider the maximally entangled two-branch state

$$|\Psi_{\mathrm{ent}}\rangle = \frac{1}{\sqrt{2}}\Big(|\psi_0\rangle^{(A)} \otimes |\psi_0\rangle^{(B)} + |\psi_1\rangle^{(A)} \otimes |\psi_1\rangle^{(B)}\Big).$$

Define the branch-resolved single-particle Peres matrix elements

$$T^{(k)}_{\alpha\beta} := \Big\langle\psi^{(k)}_\alpha\Big|\hat{T}_{\mathcal{R}_k}\Big|\psi^{(k)}_\beta\Big\rangle, \qquad \alpha,\beta \in \{0,1\}, \quad k = A, B.$$

Then the one-particle expectations are

$$\langle T^{(k)}_P\rangle = \big\langle\Psi_{\mathrm{ent}}\big|\hat{T}_{\mathcal{R}_k} \otimes \mathbb{I}\big|\Psi_{\mathrm{ent}}\big\rangle = \frac{1}{2}\Big(T^{(k)}_{00} + T^{(k)}_{11}\Big),$$

For the joint expectation value we obtain

$$\begin{aligned}\langle T^{(A)}_P T^{(B)}_P\rangle &= \big\langle\Psi_{\mathrm{ent}}\big|\hat{T}_{\mathcal{R}_A} \otimes \hat{T}_{\mathcal{R}_B}\big|\Psi_{\mathrm{ent}}\big\rangle \\ &= \frac{1}{2}\Big(T^{(A)}_{00}T^{(B)}_{00} + T^{(A)}_{11}T^{(B)}_{11}\Big) + \Re\Big(T^{(A)}_{01}T^{(B)}_{01}\Big).\end{aligned}$$

Hence the covariance is

$$\begin{aligned}\mathrm{Cov}\Big(T^{(A)}_P, T^{(B)}_P\Big) &= \frac{1}{4}\Big(T^{(A)}_{00} - T^{(A)}_{11}\Big)\Big(T^{(B)}_{00} - T^{(B)}_{11}\Big) \\ &\quad + \Re\Big(T^{(A)}_{01}T^{(B)}_{01}\Big).\end{aligned}$$

Key point: the off-diagonal contribution $\Re\Big(T^{(A)}_{01}T^{(B)}_{01}\Big)$ is what distinguishes the entangled state from the corresponding classical mixture with identical one-particle marginals.

Science Fund of the Republic of Serbia

Why choose a subregion?

If the spatial integration covered the full support, orthogonality of ψ_0 and ψ_1 would suppress off-diagonal terms. By choosing a detector region such as a half-line, the restricted overlaps become nonzero and the covariance becomes sensitive to entanglement.

This choice also has a clear operational meaning: the Peres-time observable is associated with a *local detector region* inside the laboratory, rather than with the entire support of the wavefunction. In this way, the measurement becomes genuinely selective and probes how much of each branch passes through the chosen spatial sector. The subregion therefore plays a double role: it prevents the off-diagonal contribution from vanishing trivially, and it provides a realistic local measurement setting from which temporal correlations can be constructed.

Explicit example: free fall in AdS_2

Near a timelike geodesic in anti–de Sitter spacetime,

$$R_{0i0j} = \ell^{-2}\delta_{ij}, \qquad \omega = \ell^{-1}.$$

The effective Hamiltonian becomes

$$\hat{H} = m\mathbb{I} + \frac{\hat{p}^2}{2m} + \frac{1}{2}m\omega^2x^2.$$

Choose the first two oscillator eigenstates $\psi_0(x)$ and $\psi_1(x)$, and the Peres region

$$\mathcal{R} = \{x > 0\}.$$

Then

$$T_{00} = T_{11} = \frac{\tau_*}{2}, \qquad T_{01} = \frac{1}{\sqrt{2\pi}}\frac{1 - e^{-i\omega\tau_*}}{i\omega}.$$

Hence

$$\hat{T}_{\mathcal{R}} = \begin{pmatrix}\tau_*/2 & T_{01} \\ T^*_{01} & \tau_*/2\end{pmatrix}.$$

Main covariance result

For the entangled state,

$$\mathrm{Cov}_{\mathrm{ent}}\big(T^{(A)}_P, T^{(B)}_P\big) = \Re\Big(T^{(A)}_{01}T^{(B)}_{01}\Big),$$

while for the corresponding classical mixture,

$$\mathrm{Cov}_{\mathrm{mix}}\big(T^{(A)}_P, T^{(B)}_P\big) = 0.$$

In the symmetric AdS case,

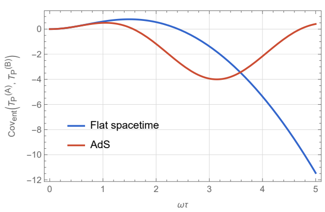
$$\mathrm{Cov}^{\mathrm{AdS}}_{\mathrm{ent}} = \frac{2}{\pi\omega^2}\sin^2\Big(\frac{\omega\tau_*}{2}\Big)\cos(\omega\tau_*).$$

This result makes the central distinction transparent. Although the entangled state and the corresponding classical mixture have the same one-particle marginals, they differ at the level of joint temporal correlations. In particular, the entangled case retains the interference contribution encoded in the off-diagonal matrix elements $T^{(A)}_{01}$ and $T^{(B)}_{01}$, whereas this contribution is absent for the classical mixture. The Peres-time covariance therefore acts as an operational witness of coherence between the two branches.

Covariance: flat spacetime vs. AdS

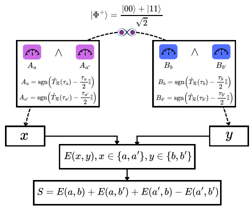
For the same initial branch wavefunctions, the Peres-time covariance exhibits qualitatively different behaviour in flat spacetime and in anti–de Sitter spacetime. In AdS, the local tidal term produces an effective harmonic potential, leading to bounded oscillatory behaviour of the off-diagonal Peres amplitude and hence of the covariance. In flat spacetime, by contrast, the same initial states evolve under free-particle dynamics, and the covariance develops a different time dependence associated with dispersive spreading.

This comparison shows that the covariance is sensitive not only to entanglement, but also to the geometry of the background spacetime. It therefore provides an operationally defined curvature-sensitive observable.



Blue: flat spacetime. Red: AdS. The qualitative difference arises from the distinct background-induced evolution.

Bell-protocol schematic



Bell-type protocol realized through local Peres-time measurements.

Alice and Bob each choose an interrogation time in their laboratory, which fixes the corresponding effective qubit observable. The measured correlations are then combined into the CHSH parameter, whose value is shifted by the background curvature.

Bell-type protocol

We recast the Peres-time setup as a Bell experiment between two distant observers, Alice and Bob, each holding one particle of an entangled pair. In each laboratory, the particle is probed only locally through a Peres-time measurement associated with a chosen detector region and a chosen interrogation time.

The key idea is that, when the dynamics is restricted to the two-mode subspace

$$\mathrm{span}\{|0\rangle, |1\rangle\},$$

the Peres-time operator takes the form of an effective qubit observable. Its off-diagonal matrix element can be written as

$$T_{01}(\tau) = |T_{01}(\tau)|e^{-i\phi(\tau)},$$

so that after binarization the corresponding local observable becomes

$$\hat{O}(\tau) \simeq \cos\phi(\tau)\,\sigma_x + \sin\phi(\tau)\,\sigma_y.$$

Therefore, each choice of interrogation time τ defines a measurement direction on the equator of the Bloch sphere. Alice chooses between two local settings a, a', and Bob between b, b', exactly as in the standard CHSH protocol, but here the settings are implemented dynamically through local time-integrated measurements rather than by rotating spin analyzers.

Curvature-induced shift

Curvature does not modify the Bell inequality itself, and the local protocol is kept the same as in flat spacetime. What changes is the effective local dynamics. In each laboratory, the tidal term shifts the oscillator frequency,

$$\omega_k^2 = \omega_0^2 + R^{(k)}_{0x0x}, \qquad \delta\phi_k \simeq \frac{\tau_k}{4\omega_0}R^{(k)}_{0x0x},$$

so the same interrogation times correspond to different effective measurement axes.

$$\delta S \simeq \frac{1}{4\omega_0}\Big[(\tau_a + \tau_{a'})R^{(A)}_{0x0x} - (\tau_b + \tau_{b'})R^{(B)}_{0x0x}\Big].$$

Hence a protocol calibrated to give $S_{\mathrm{flat}} = 2$ in flat spacetime is shifted to

$$S_{\mathrm{curv}} = 2 + \delta S.$$

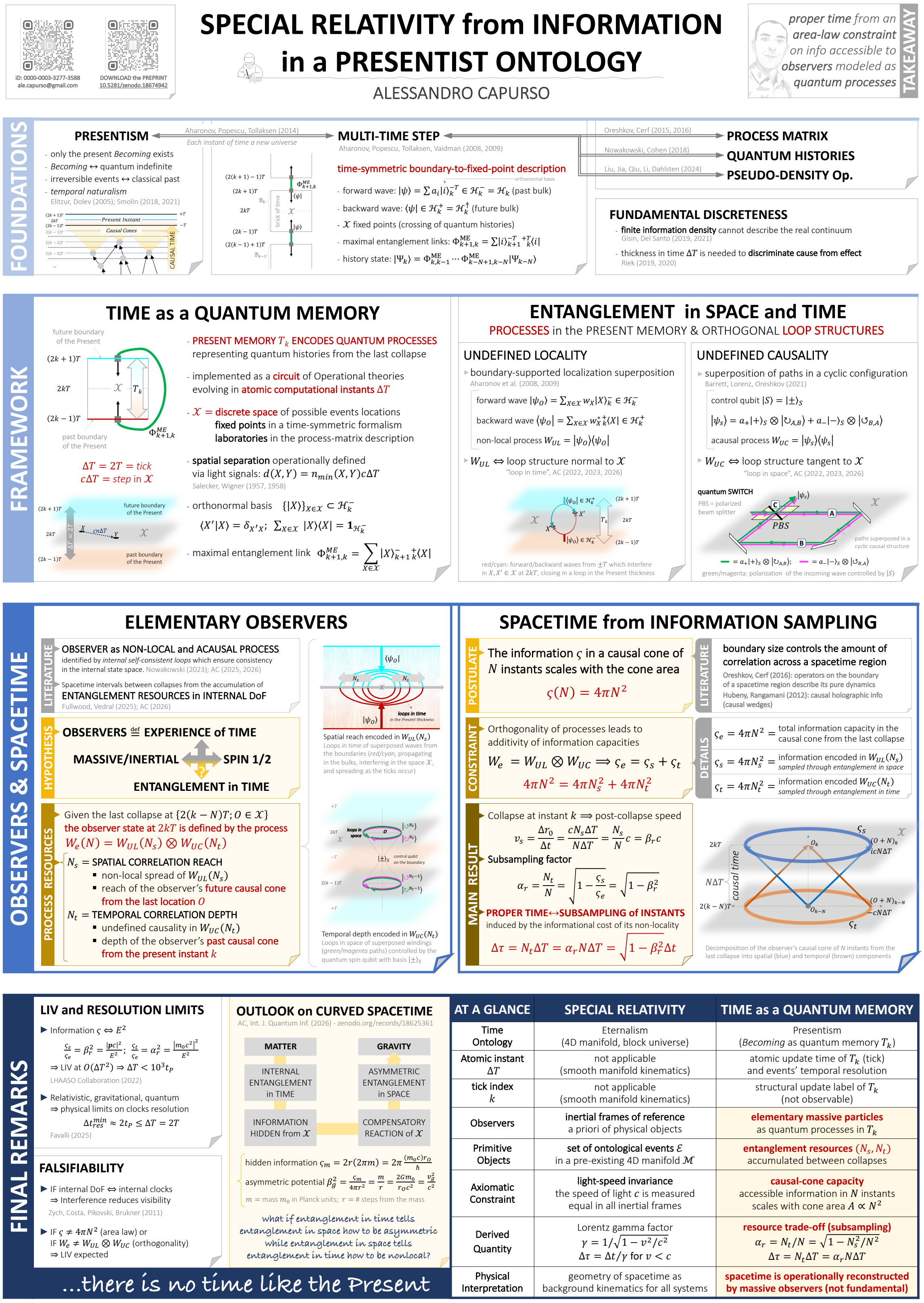
For positive tidal curvature, this yields a positive correction and can drive the CHSH parameter above the classical bound. The Bell parameter therefore acts as an **operational witness of background curvature**.

Take-home messages

1. Peres-time covariance distinguishes entanglement from a classical mixture with identical one-particle marginals.
2. The relevant signature appears in joint temporal correlations, not in reduced one-particle dynamics.
3. The covariance is also sensitive to the background spacetime and can serve as a curvature-sensitive observable.
4. In the Bell reformulation, local interrogation times play the role of effective measurement settings.
5. Curvature shifts these settings dynamically, so the CHSH parameter becomes an operational curvature witness.

Alessandro Capurso

Time as a Quantum Memory, Causality as a Resource: The Minkowski Metric from an Information-Theoretic Axiom



Andreas Leitherer

Paradox-free classical non-causality and unambiguous non-locality without entanglement are equivalent

Classical non-causality and nonlocality without entanglement are equivalent

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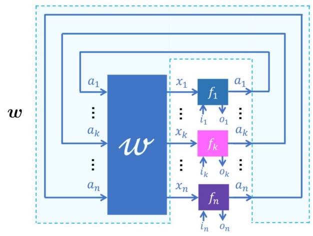
Abstract

Closed timelike curves (CTCs) challenge our conception of causality by allowing information to loop back into its own past. Process functions [1] are information-theoretic CTCs that generalize deterministic classical communication. Certain process functions can violate causal inequalities, signifying non-causality, while being logically consistent: time-travel paradoxes are avoided without restricting local operations. We show that process functions correspond to *unambiguous* product bases in which every local state belongs to a unique local basis [2] (significantly extending [3] to arbitrary dimensions and number of parties). Non-causality of process functions is exactly mirrored by *quantum nonlocality without entanglement*, i.e., the impossibility of perfectly distinguishing separable states using local operations and classical communication [4]. We represent this equivalence using hypergraphs in a theory-independent way [5,6,7]. All process functions - including non-causal ones - emerge in a specific composite hypergraph which surprisingly, makes no assumption about causal order while invoking the no-signaling principle. This work suggests an interpretation of classical non-causality as a non-standard way of wiring theory-independent measurements.

Process functions [1]

- Collection of non-self-signaling functions

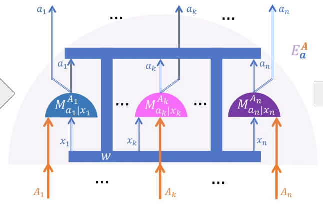
$$\boldsymbol{w} = (w_k)_k, \quad w_k : \mathcal{A}_{\setminus k} \to \mathcal{X}_k$$



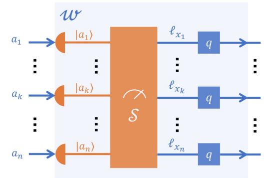
- $\boldsymbol{w}$ is logically consistent iff *p(**o**|**i**)* define valid probabilities or [2]:
  iff $\boldsymbol{w}$ is valid for all fixed local operations that are constant and non-erasing
- There are non-causal $\boldsymbol{w}$ [1]:

$$(x, y, z) = (w_A(b,c), w_B(a,c), w_C(a,b)) \qquad \begin{aligned} x &= c\ (b \oplus 1) \\ y &= a\ (c \oplus 1) \\ z &= b\ (a \oplus 1) \end{aligned}$$

Joint projective measurement from local operations with process function [2]:



Classical channel underlying $\boldsymbol{w}$ from projective measurement [2]:



Unambiguous product bases [2]

- *Complete* product basis of *n*-partite Hilbert space:

$$\mathcal{S} = \{|\psi^j\rangle = |\psi_1^j\rangle \otimes ... \otimes |\psi_n^j\rangle\}_{j=1}^{|S|}$$

- For each party, define local sets and partition them into local bases

$$\mathcal{S}^{(k)} = \{|\psi_k^j\rangle\}, \quad \mathcal{S}^{(k)} = \bigsqcup_{x_k=0}^{|\mathcal{X}_k|-1} \mathcal{S}_{x_k}^{(k)}$$

- x_k is measurement setting, position of local state is outcome a_k
- *Unambiguity*: two local vectors are orthogonal iff they belong to the same local basis $\mathcal{S}_{x_k}^{(k)}$
- Basis exhibits quantum nonlocality without entanglement if each local set has more than one local basis:

$$\{\,|000\rangle, |{+}01\rangle, |01{+}\rangle, |01{-}\rangle, |1+0\rangle, |{-}01\rangle, |1-0\rangle, |111\rangle\}$$ *SHIFT* basis [3]

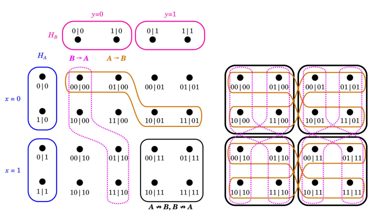
$$\mathcal{S}^{(k)} = \{|0\rangle, |1\rangle, |+\rangle, |-\rangle\},$$

$$\mathcal{S}_{x_k=0}^{(k)} = \{|0\rangle, |1\rangle\}, \; S_{x_k=1}^{(k)} = \{|+\rangle, |-\rangle\}$$

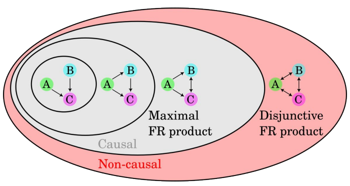
$$|000\rangle \to (000|000) \quad |{+}01\rangle \to (001|100)$$

Non-causality from no-signaling [7]

$$\sum_{w \in e} p(v,w) = \sum_{w \in e'} p(v,w) \quad \forall v \in V(H_A), e, e' \in E(H_B)$$



- Hypergraph $H = (V, E)$ defines *contextuality scenario* [5,6]
- Probabilistic model $p : V \to [0,1]$ is theory-independent
- Composite contextuality scenarios: p satisfies no-signaling only if allow for communication
- Foulis-Randall (FR) product achieves this for N=2 parties; hierarchy of products for N>2

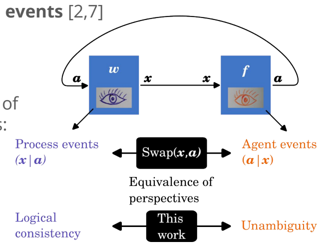


- *Maximal* FR product: hyperedges are all causal process functions
- *Disjunctive* FR product: hyperedges are all process functions
- Graphical criterion for non-causality: $\boldsymbol{w}$ in disjunctive but not maximal FR product
- Causal order was not predefined but emerges from product rule

Theory-independent character of non-locality without entanglement and non-causality. Non-causal process functions as non-standard wirings.

Agency and events [2,7]

Equivalence of perspectives: who is the agent?



What is an event?

$$(\boldsymbol{a}|\boldsymbol{x}) = (a_1, ..., a_n | x_1, ..., x_n)$$

Specific event labeling captures logical consistency

Unambiguity:	Complete basis:
Local states require **fine-grained labeling** (both setting and output)	**Global states** require **coarse-grained labeling** (only output)
Identification of **local events** is **setting-dependent**	Identification of **global events** is **setting-independent**

Non-causality: no local event can be identified in setting-independent manner: $\nexists\, k : (a_k | x_k = x_0)_{a_k}$

Related to Grinbaum [8]

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Anna Horváth

Quantum effects of strong gravity in the Kaluza-Klein theory

Quantum effects of strong gravity in the Kaluza–Klein theory

Anna Horváth[1,2], A. Wojnar[3], G.G. Barnaföldi[1]: *The effects of strong gravity on the dispersion relation of massive particles in the Kaluza–Klein theory*, (2025) arXiv: 2510.16631

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Can strong gravity have a quantum effect on particles?
Could it cause them to decay?
We investigate observational possibilities of such a phenomenon in general relativity, and in the extra-dimensional Kaluza–Klein spacetime.

Extended Uncertainty Principle (EUP)

In order to treat particles within a quantum mechanical framework, a separation of time and space is needed. This can be done using the ADM formalism, with which the uncertainty relation can be written as

$$\sigma_{\mathrm{p}}\rho \gtrsim \pi\hbar\left[1-\frac{\rho^2\mathcal{R}|_{x_0}}{12\pi^2}+\xi\frac{\rho^4}{\lambda_C^2}\bar{\nabla}_j N_i\bar{\nabla}^j N^i\Big|_{x_0}\right] .$$

L. Petruzziello, F. Wagner, PRD 103, 104061 (2021)

Observations

The effect considered is tiny, however, there could still be a possibility for observation. These include:

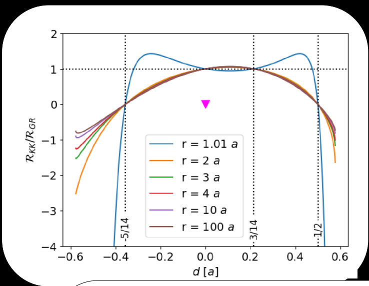
- modifications to neutron star equation of state
- position of photon ring
- Hawking radiation

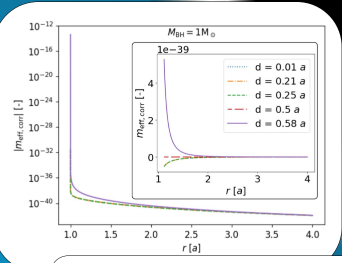
Effective mass and particle decay

From the modified uncertainty relation, one can write a modified dispersion relation

$$p^\mu p_\mu = \overbrace{k^\mu k_\mu}^{m^2} - \frac{\hbar^2\mathcal{R}}{6} + \hbar^2\xi\frac{\rho^2}{2\lambda_{\mathrm{C}}^2}\bar{\nabla}_j N_i\bar{\nabla}^j N^i = m_{\mathrm{eff}}^2$$

which results in an effective mass. When the modification is a large negative number, the effective mass becomes imaginary. Then the quantum mechanical evolution describes the decay of the state.





Kaluza–Klein spacetime

Kaluza–Klein is a 5D theory with one additional compactified spatial dimension. However, one can perform dimensional reduction, such that the 5th direction is integrated out. Thus it effectively becomes a scalar-tensor theory of gravity with action

$$S=\int\frac{\sqrt{-g}}{16\pi\mathbb{G}\mathrm{e}^{-\sigma}}\left[R-\frac{\kappa^2}{4}\mathrm{e}^{2\sigma}F_{\mu\nu}F^{\mu\nu}-2\mathrm{e}^{-\sigma}\Box\mathrm{e}^{\sigma}\right]\mathrm{d}^4x ,$$

written in the Jordan frame. Not taking the vector potential into account, solving the field equations for a spherically symmetric spacetime results in the metric

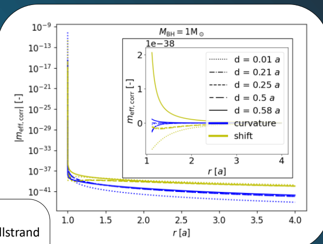
$$\mathrm{d}s^2=-\left(1-\frac{a}{r}\right)^{\frac{b-d}{a}}\mathrm{d}t^2+\left(1-\frac{a}{r}\right)^{-\frac{b+d}{a}}\mathrm{d}r^2+r^2\left(1-\frac{a}{r}\right)^{1-\frac{b+d}{a}}\mathrm{d}\Omega^2 ,$$

with a constraint for the parameters *a*, *b* and *d*.

R. Coquereaux, G. Esposito-Farese, in *Annales de l'IHP Physique Théorique,* Vol. 52 (1990) pp. 113-150.

Observers – static

The observer of the effect is the Eulerian observer of the ADM slicing. In Schwarzschild coordinates this corresponds to the hovering observer: static and accelerating. In this case the shift vector vanishes, so only the phase space curvature contributes.



Observers – flat spacetime

Modifications to the uncertainty and dispersion relations do not only appear when the spacetime is curved. Even in flat spacetime written in Rindler coordinates

$$\mathrm{d}s^2=-(\alpha x)^2\mathrm{d}t^2+\mathrm{d}x^2+\mathrm{d}y^2+\mathrm{d}z^2 ,$$

where the observer is an accelerating one, the curvature of the phase space is non-zero, showing a connection to the Unruh effect.

Observers – free-falling

To study what a free-falling, raindrop observer sees, one needs to perform a coordinate transformation to Painlevé-Gullstrand coordinates, with which the Kaluza-Klein metric becomes

$$\mathrm{d}s^2=-\left(1-\frac{a}{r}\right)^{\frac{b-d}{a}}\mathrm{d}T^2++2\left[\left(1-\frac{a}{r}\right)^{-\frac{2d}{a}}-\left(1-\frac{a}{r}\right)^{\frac{b-3d}{a}}\right]^{\frac{1}{2}}\mathrm{d}T\mathrm{d}r++\left(1-\frac{a}{r}\right)^{-\frac{2d}{a}}\mathrm{d}r^2+r^2\left(1-\frac{a}{r}\right)^{1-\frac{b+d}{a}}\mathrm{d}\Omega^2 .$$

In this case, both the phase space curvature and the shift vector contribute.

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Can observers "agree to disagree" in quantum mechanics?

Based on arXiv:2603.23595

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Abstract

Agreement theorems are no-go results about rational disagreement: if two agents start from a common prior and their posterior beliefs are common knowledge, they cannot assign different probabilities to the same event. Standard treatments of the result have the agents reason about an underlying state of the world, which has led some to ask whether the result can extend to quantum or postquantum phenomena, where such a description may no longer be appropriate. We derive an operational version of Aumann's agreement theorem without assuming an objective state of the world and instead focusing only on what is observed. This allows us to establish the theorem's validity in quantum theory and even in situations with indefinite causal order or involving hypothetical postquantum phenomena. We comment on seemingly contradictory results in the literature and point to the one place where the theorem might fail: Wigner's friend-type situations.

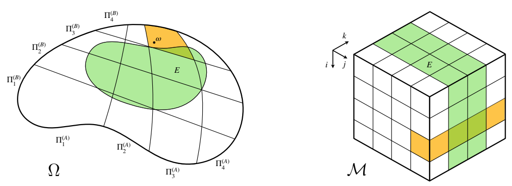
Classical Agreement Theorem

Consider a finite state space Ω, a common prior p on Ω, and a true state of the world $\omega^* \in \Omega$. Alice and Bob obtain coarse-grained information via partitions of Ω; after observing their respective information sets, they update by ordinary Bayesian conditioning to form posteriors for an event of interest $E \subseteq \Omega$.

Common knowledge of posteriors. Alice and Bob have common knowledge when Alice knows Bob's posterior, Bob knows Alice's posterior, Alice knows that Bob knows hers, and so on indefinitely (and vice versa). Specifically, fix two numbers $q_A, q_B \in [0,1]$. Let A_0 be the set of states where Alice's posterior equals q_A, and B_0 the set of states where Bob's posterior equals q_B. Define A_1 as the subset of A_0 where Alice can deduce (from her information set) that Bob must be in B_0; define B_1 analogously. Continue iterating: A_{n+1} is the subset of A_n where Alice can deduce that Bob is in B_n, and B_{n+1} is the subset of B_n where Bob can deduce that Alice is in A_n. The posteriors are common knowledge at ω^* when ω^* survives all rounds of this reasoning.

Aumann's theorem (1976). If Alice's posterior is q_A and Bob's posterior is q_B, and this is common knowledge at ω^*, then $q_A = q_B$: two Bayesian agents with a common prior cannot "agree to disagree".

Figure



Does This Hold in Quantum Theory?

What happens when Alice and Bob's measurements do not merely reveal information about a pre-existing state of the world, but instead update the state itself? Different claims have been made in the literature:

- Contreras–Tejada *et al.* (*Nat. Commun.*, 2021): agreement holds for quantum non-signalling boxes, but **postquantum boxes can support disagreement** about counterfactual outcomes.
- Brandenburger *et al.* (*Phil. Trans. R. Soc. A*, 2024): using signed probabilities, they show that there can be **common certainty of disagreement** in quantum mechanics.
- García Díaz *et al.* (arXiv:2511.21258): there can be **common certainty of disagreement in noncommuting scenarios**.

These results appear to point in different directions because they do not adopt the same notions of event, posterior, or common knowledge. The operational theorem clarifies that, once these notions are formulated purely in terms of jointly observable outcomes, agreement is recovered.

Operational Agreement Theorem

We formulate agreement in an **operational** way: instead of referring to underlying states of the world, our formulation involves only **measurement outcomes**. Let $\mathcal{M} = \mathcal{I} \times \mathcal{J} \times \mathcal{K}$, where i is Alice's observation, j is Bob's observation, and k labels the event of interest, with common prior $p(i,j,k)$. After observing only their own outcomes, Alice and Bob update by Bayesian reasoning. For an event $E \subset \mathcal{K}$, define

$$A_0 = \{i : p_A(E \mid i) = q_A\}, \qquad B_0 = \{j : p_B(E \mid j) = q_B\},$$

and iterate

$$A_{n+1} = \{i \in A_n : p_A(B_n \mid i) = 1\}, \quad B_{n+1} = \{j \in B_n : p_B(A_n \mid j) = 1\}.$$

The posteriors are common knowledge at (i,j) when $(i,j) \in \bigcap_{n \geq 0}(A_n \times B_n)$.

Operational agreement theorem. If a theory provides such a joint distribution for the relevant observations, then common knowledge of posteriors implies

$$q_A = q_B.$$

So agreement **does not require a hidden classical state space, but only a consistent joint probability assignment** for the outcomes entering the agents' reasoning.

Applications

Quantum mechanics: Let $\mathcal{H} = \mathbb{C}^4$. Alice measures in the basis $\{|a_i\rangle\}$, and Bob measures in the basis $\{|b_j\rangle\}$, with

$$|b_0\rangle = c_\theta\,|a_0\rangle + s_\theta\,|a_1\rangle\,, \;\; |b_1\rangle = -s_\theta\,|a_0\rangle + c_\theta\,|a_1\rangle\,,$$
$$|b_2\rangle = c_\phi\,|a_2\rangle + s_\phi\,|a_3\rangle\,, \;\; |b_3\rangle = -s_\phi\,|a_2\rangle + c_\phi\,|a_3\rangle\,,$$

and event measurement $E_0 = |e_0\rangle\langle e_0|$, $E_1 = \mathbb{I} - |e_0\rangle\langle e_0|$ with

$$|e_0\rangle = \sqrt{q}\,|b_0\rangle + \sqrt{q}\,|b_1\rangle + \sqrt{r}\,|b_2\rangle + \sqrt{1-2q-r}\,|b_3\rangle\,.$$

Then

$$q_A(i) = \big(q, q, c_\phi^2 r + s_\phi^2(1-2q-r), s_\phi^2 r + c_\phi^2(1-2q-r)\big),$$

$$q_B(j) = \big(q, q, r, 1-2q-r\big).$$

When common knowledge holds (e.g. Alice sees $i = 0$), the posteriors coincide; when it does not hold (e.g. $i = 2$), they may differ.

Indefinite causal order: the theorem holds even in scenarios with indefinite causal order, as long as they know the process matrix.

Postquantum: the same logic holds in GPT-like frameworks whenever a joint $p(i,j,k)$ exists, independently of the theory.

Wigner's friend? A likely boundary is Wigner's friend-type situations, where the consistency of a shared outcome description across observers can break down. If no single joint probability assignment is available for all relevant outcome statements, the assumptions of the operational agreement theorem fail, and disagreement can persist.

Caroline Lima

On sufficient conditions for holographic scattering.

On sufficient conditions for holographic scattering

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Introduction, Definitions, and Setup

The **Ryu-Takayanagi (RT) formula** $S_A = \frac{\text{Area}(\gamma_A)}{4G_N}$ relates the entanglement entropy between a boundary region A and its complement to minimal codim-2 homologous surfaces in the bulk. It kicked off a vast literature that has revealed how quantum entanglement on the boundary theory is key to the emergence of the bulk geometry in the semiclassical limit.

The **entanglement wedge** of a boundary region A, $\mathcal{E}[A]$, is the causal development of the region enclosed between A and the RT surface γ_A. If there exists more than one minimal geodesic, γ_A is chosen to make the entanglement wedge as large as possible. Example:

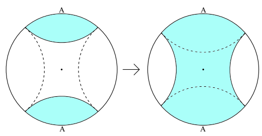


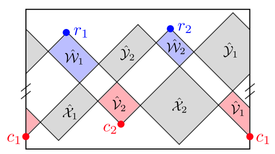
Figure 1: Figure extracted from [1]

Setup: Take input points c_1, c_2 and output points r_1, r_2 on the boundary of an asymptotically AdS_3 spacetime, such that c_1 and c_2 are in the past of both r_1 and r_2. Define the boundary regions:

$$\hat{\mathcal{V}}_i \equiv \hat{J}^+(c_i) \cap \hat{J}^-(r_1) \cap \hat{J}^-(r_2), \; \hat{\mathcal{W}}_i \equiv \hat{J}^+(c_1) \cap \hat{J}^+(c_2) \cap \hat{J}^-(r_i).$$

Boundary scattering region: $\hat{\mathcal{S}} \equiv \hat{J}^+(c_1) \cap \hat{J}^+(c_2) \cap \hat{J}^-(r_1) \cap \hat{J}^-(r_2)$.

Bulk scattering region: $\mathcal{S} \equiv J^+(c_1) \cap J^+(c_2) \cap J^-(r_1) \cap J^-(r_2)$.



Holographic Scattering and the CWT

Holographic scattering or **bulk-only scattering** is the phenomenon of concomitantly having $\hat{\mathcal{S}} = \emptyset$ and $\mathcal{S} \neq \emptyset$. In order not to contradict the holographic principle, the boundary must be able to reproduce the outcomes of such processes, which is accomplished by using the entanglement of the boundary joint quantum state [2].

Thm. [Connected Wedge Theorem (CWT)]: Pick points c_1, c_2, r_1, r_2 on the boundary of an asymptotically AdS spacetime with a holographic dual, such that $\hat{\mathcal{S}} = \emptyset$ and $\mathcal{S} \neq \emptyset$. Then $\mathcal{E}[\hat{\mathcal{V}}_1 \cup \hat{\mathcal{V}}_2]$ is connected [3, 4].

The CWT is an **if and only if statement in pure AdS$_3$**. In other asymptotically AdS_3, **the converse does not hold**. That is, **a connected entanglement wedge does not imply that scattering is possible through the bulk**.

Towards an if and only if statement

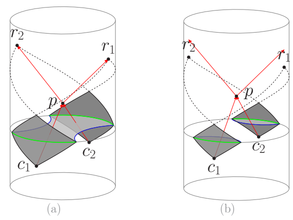


Figure 2: Figure reproduced from [2].

How can we modify the CWT to get a theorem with a converse?

Option 1: Add extra assumptions on top of connectedness of the entanglement wedge. In another work, we have shown that for vacuum geometries, if a certain relation between non-minimal Ryu-Takayanagi surfaces is satisfied, then the converse holds. See 1st QR code [5].

Option 2: Replace the scattering bulk region $\mathcal{S}$ with some bulk region that contains $\mathcal{S}$ (but it is in general larger). For pure AdS_3, this region must reduce to $\mathcal{S}$ exactly, as in that case the theorem is already an iff statement. More below and on the 2nd QR code:

Our Proposal: modified bulk region

Based on the fact that the operator algebras associated to local subregions in large N field theories display the phenomenon of superadditivity, Leutheusser and Liu conjecture a Generalized Connected Wedge (GCW) proposal by defining the bulk subregion [6]: $\mathcal{S}_E \equiv \mathcal{E}[\hat{\mathcal{V}}_1 \cup \hat{\mathcal{V}}_2] \cap \mathcal{E}[\hat{\mathcal{W}}_1 \cup \hat{\mathcal{W}}_2]$, which is guaranteed to contain the usual scattering region $\mathcal{S}$. Their conjecture was that $\mathcal{E}[\hat{\mathcal{V}}_1 \cup \hat{\mathcal{V}}_2]$ is connected if and only if $\mathcal{S}_E \neq \emptyset$.

We established a modified version of the GCW by introducing an alternative bulk region $\tilde{\mathcal{S}}_E$ which subtracts the individual entanglement wedges:

$$\tilde{\mathcal{S}}_E \equiv (\mathcal{E}[\hat{\mathcal{V}}_1 \cup \hat{\mathcal{V}}_2] \backslash (\mathcal{E}[\hat{\mathcal{V}}_1] \cup \mathcal{E}[\hat{\mathcal{V}}_2])) \cap (\mathcal{E}[\hat{\mathcal{W}}_1 \cup \hat{\mathcal{W}}_2] \backslash (\mathcal{E}[\hat{\mathcal{W}}_1] \cup \mathcal{E}[\hat{\mathcal{W}}_2]))$$

We proved: $\mathcal{E}[\hat{\mathcal{V}}_1 \cup \hat{\mathcal{V}}_2]$ and $\mathcal{E}[\hat{\mathcal{W}}_1 \cup \hat{\mathcal{W}}_2]$ are connected $\Leftrightarrow \tilde{\mathcal{S}}_E \neq \emptyset$.

Quantum Information Interpretation

While the non-emptiness of the usual region $\mathcal{S}$ maps to a boundary non-local quantum computation task, $\mathcal{S}_E \neq \emptyset$ relates naturally to subsystem quantum error correction with complementary recovery. If Alice has access to $\hat{\mathcal{V}}_1 \cup \hat{\mathcal{V}}_2$ and Bob to $\hat{\mathcal{W}}_1 \cup \hat{\mathcal{W}}_2$, $\mathcal{S}_E \neq \emptyset$ ensures there exists a choice of operations allowing Alice to safely encode and Bob to successfully decode a message, completely protected against an adversary (Eve) acting on the complementary CFT regions.

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David Amaro-Alcalá

Group functions in filtered randomized benchmarking for passive bosonic devices

Kostant Relation in Filtered Randomized Benchmarking

for Passive Bosonic Devices
Immanants, SU(m) characters, and feasible optics

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arXiv: 2511.00842

PRA accepted: doi.org/10.1103/xmlq-vssg

1. INTRODUCTION

Goal. Benchmark passive bosonic interferometers.

Bottleneck. The original filter needs projectors, Clebsch–Gordan data, permanents, Fock inputs, and number-resolving detectors.

This work. Use immanants and SU(m) characters as filters.

Payoff. Keep single-exponential decay while simplifying post-processing and experimental requirements.

2. BACKGROUND

Representation. For n indistinguishable photons in m modes, the channel splits into irreps, and benchmarking estimates $p_\mu(\mathcal{E})$.

$$\Gamma(U) = \mathcal{U} \otimes \overline{\mathcal{U}} \cong \bigoplus_\mu \mu$$

$$F(\mathcal{E}) = d_\lambda^{-2} \sum_\mu d_\mu p_\mu(\mathcal{E})$$

Bridge. Kostant maps zero-weight traces to immanants; whole-irrep traces give SU(m) characters. These become our filters.

3. APPROACH

1. Data. Sample sequences $\mathbf{U}_{g,s}$ and record $d^{(g,s)}(\widetilde{\mathbf{U}}_{g,s})$.

2. Filters. Build two filters from the same benchmarking data:

$$f_{\mathrm{imm},\mu} = \mathrm{Imm}_\mu, \qquad f_{\chi,\mu} = \chi_\mu$$

3. Average. Form the filtered signal:

$$\Phi_g^{(f)} = \langle f_\mu^{(g,s)} d^{(g,s)} \rangle_s \propto p_\mu^{g-1}$$

4. Fit. One decay per irrep yields the p_μ values and reconstructs $F(\mathcal{E})$.

Key simplification. Same data, but no explicit projectors or Clebsch–Gordan coefficients.

4. RESULTS

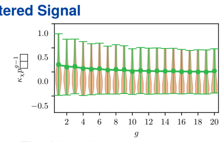
Filtering Procedure

Using the sampled data matrix D and filter matrix $F_\mu^{(f)}$,

$$\Phi_\mu^{(f)} = \sum_s (F_\mu^{(f)} \odot D)_{s,g} \propto p_\mu^{g-1}.$$

Fitting $\{g, \Phi_\mu^{(f)}\}$ yields p_μ.

Filtered Signal



The filtered quantity shows an approximately exponential trend as the circuit depth increases.

Gain/Loss Extension

The analysis extends to

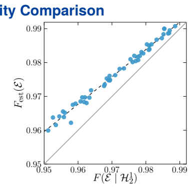
$$(\mathcal{H}_e)_m^n = \bigoplus_{n \geq 0} \mathcal{H}_m^n.$$

For a weak coherent input,

$$|\alpha\rangle \approx |0,0\rangle + \alpha|1,0\rangle + O(|\alpha|^2).$$

Intensity measurement still supports the filtered estimate.

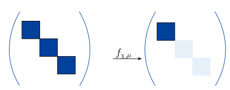
Fidelity Comparison



The estimate tracks the restricted single-photon fidelity closely; the observed overestimate comes from extra irreps in the extended space.

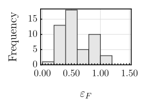
Character Filter

$$\mathrm{Var}[\chi_\mu(U)] = 1.$$



Polynomial-time evaluation; no Clebsch–Gordan coefficients.

Error Histogram



The percentage error is concentrated near zero; the paper reports about 2% error for fidelities above 0.95.

5. DISCUSSION AND CONCLUSION

Takeaway. The reformulation preserves the benchmarking decay while making the filter simpler and easier to evaluate.

Best practical option. Character filters avoid Clebsch–Gordan data, have constant variance, and are polynomial-time to evaluate.

Outlook. Extending the method to active bosonic transformations is the next natural step.

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6. DATA AND CODE AVAILABILITY

`doi.org/10.5281/zenodo.19493886`

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Dominik Šafránek

Work and entropy of mixing in isolated quantum systems

Work and entropy of mixing in isolated quantum systems

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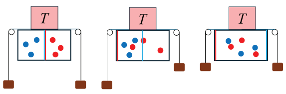
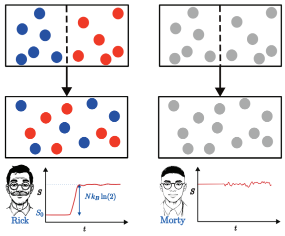
Non-quantum motivation

- CO_2 is released into the atmosphere
- Microplastics or hormones released in water supply
- Desalination → all cost energy





Usual story

Initial $S_0 = k_B \ln \Omega_1 + k_B \ln \Omega_2$

Final $S = k_B \ln \Omega$

Difference $\Delta S = S - S_0 = \cdots = k_B N(-x_1 \ln x_1 - x_2 \ln x_2)$

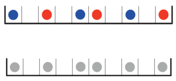
Symmetric $= k_B N \ln 2$

Indistinguishable $1/N_1!, \quad 1/N_2!$

Cost of unmixing $W = -\Delta F = T\Delta S.$

Gibbs paradox

- Rick associates higher entropy (despite having more knowledge)
- Morty associates lower (because of his ignorance)
- Morty learns of the existence of two particle types → sudden entropy increase!
  - Can this knowledge allow Morty to extract more energy?

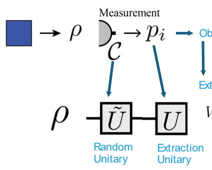


$$S_R = 2 \ln \binom{10}{3}$$

$$S_M = \ln \binom{10}{6}$$

$$\Delta S \approx 6 \ln 2$$

Work Extracted



Observational entropy

$$S_{\mathcal{C}} = S_{\mathcal{C}}^{\mathrm{Sh}} + S_{\mathcal{C}}^{B} = -\sum_i p_i \ln p_i + \sum_i p_i \ln V_i$$

Extracted work estimate

$$W_{\mathcal{C}}^{\infty}(\hat\rho) = \mathrm{tr}[\hat H(\hat\rho - \hat\rho_{\beta'})]$$

$$S_{\mathrm{vN}}(\hat\rho_{\beta'}) = S_{\mathcal{C}}$$

Extracted energy difference

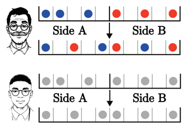
- Observational entropy difference

$$\Delta S \equiv S_{\mathcal{C}_{\mathrm{M}}} - S_{\mathcal{C}_{\mathrm{R}}}$$

- Work extraction difference

$$\Delta W \equiv W_{\mathcal{C}_{\mathrm{R}}}^{\infty} - W_{\mathcal{C}_{\mathrm{M}}}^{\infty} = \mathrm{tr}[H(\rho_{\beta_{\mathrm{M}}} - \rho_{\beta_{\mathrm{R}}})]$$

$$S_{\mathrm{vN}}(\rho_{\beta_M}) = S_{\mathcal{C}_{\mathrm{M}}} \qquad S_{\mathrm{vN}}(\rho_{\beta_R}) = S_{\mathcal{C}_{\mathrm{R}}}$$



Different levels of ignorance

	Can measure	Knows
• Rick	blue left, blue right, red left, red right	total, total blue
• Morty 1	total left, total right	total, total blue
• Morty 2	total left, total right	total, blue and red exist
• Morty 3 (ignorant)	total left, total right	total

Different coarse-grainings $\mathcal{C}_R, \mathcal{C}_{M_1}, \mathcal{C}_{M_2}, \mathcal{C}_{M_3}$

Parameters of the model N, N_+

Energy difference in small ΔS

Observational temperature (of the higher entropy S)

Observational entropy difference

$$\Delta S = N \ln 2$$

$$\Delta W = kT\Delta S\left(1 - \frac{1}{2}\frac{\Delta S}{\tilde C_E} + \frac{1}{6}\left(1 - \frac{\partial \ln \tilde C_E}{\partial T}T\right)\left(\frac{\Delta S}{\tilde C_E}\right)^2 + \cdots\right)$$

Dimensionless heat capacity $\tilde C_E = C_E/k$ Heat capacity $C_E = \partial E/\partial T$

Related to the variance in energy of the thermal state $v = -\frac{d\langle E\rangle}{d\beta} = C_E k T^2$

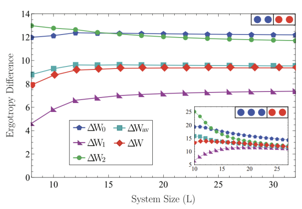
How about Morty 3

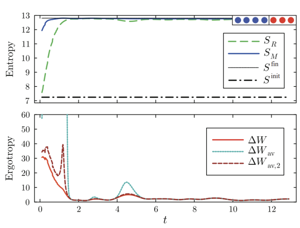
- Lower entropy than Morty 2
- But the same energy extracted!
- We proved this by studying the map
- Assumptions:
  - Consistency
    - Measurements
    - Density matrix
    - Unitary operations
  - Superselection rules
    - Superpositions $|+\rangle + |-\rangle$ can't exist and neither can be created by a
    - Unitary
    - Hamiltonian
    - Measurement
  - Color-independence of the Hamiltonian
    - $|++0\rangle, |+-0\rangle$ have the same energy

Conclusion

- Mixing is a fundamental problem
- Usual story includes a thermal bath and thermal equilibrium
- Isolated quantum systems don't have either
  - Still, we can extract work
  - Any initial state, any measurement
  - Classically, this works the same
  - Difference in extracted work $\Delta W = kT\Delta S$
- Cost of unmixing?
- Different types of Morty 3: all extract the same! → Consistent, no paradox → But details about Morty 3 matter
  - Even the one that does not know of the existence of particle types
  - Knowing of existence is not enough: one needs to be able to distinguish them by a measurement

Simulations





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https://observationalentropy.com

Dragoljub Gočanin

Rotating frames from quantum deformed spacetime

Rotating frames from quantum-deformed spacetime

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Abstract: It is often believed that the effects of quantum gravity at certain scales can be captured by making the spacetime coordinates fail to commute. One way to introduce this noncommutativity is by using the twist deformation. We can consider matter fields propagating on spacetime with twist deformation and obtain corrections to the classical equations of motion. Interestingly, assuming the noncommutativity parameter is small, the resulting equations can be written as classical (undeformed) equations with a specific effective spacetime metric. We will consider axially symmetric Melvin's solution of Einstein-Maxwell's equations with a special kind of Killing twist deformation and show that the resulting effective spacetime metric corresponds to a reference frame rotating with a constant angular velocity proportional to the field's charge and the parameter of noncommutativity, regardless of whether the matter field is a scalar boson or a Dirac fermion. This fact can be used to obtain a quantum gravity-induced version of the Sagnac effect.

One way to introduce the notion of a quantum spacetime is to promote coordinates to operators that do not commute, namely

$$[\hat{x}^\mu, \hat{x}^\nu] = i\theta(\hat{x})$$

In this work, we use the twist deformation formalism explained in [1] to analyze the effects of NC geometry. Important ingredients of this formalism are:

- A solution of Einstein-Maxwell's equations with some isometries.
- A set of commuting Killing vector fields that are used to construct a twist

$$\mathcal{F} = \exp\left(-\frac{i}{2}\theta^{IJ} X_I \otimes X_J\right)$$

which introduces a NC star-product

$$(f \star g)(x) = \mu\left(\mathcal{F}^{-1}(f \otimes g)\right)$$
$$= f(x)g(x) + \frac{i}{2}\theta^{IJ} X_I[f] X_J[g] + \mathcal{O}(\theta^2)$$

We will consider a particular solution of Einstein-Maxwell's equations, with and without a cosmological constant, known as the Melvin's universe [2]. The metric of this spacetime, for vanishing cosmological constant, is given by

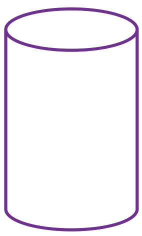
$$ds^2 = \left(1 + \frac{E^2}{4}\rho^2\right)^2 (-dt^2 + d\rho^2 + dz^2) + \frac{\rho^2}{\left(1 + \frac{E^2}{4}\rho^2\right)^2} d\phi^2$$

Electromagnetic gauge field is given by

$$A = Ez\mathrm{d}t \quad \Rightarrow \quad F = E\mathrm{d}z \wedge \mathrm{d}t$$

This spacetime has cylindrical symmetry, as a consequence of having two Killing vector fields:

$$X_1 = \partial_z$$
$$X_2 = \partial_\varphi$$



Similarly, we can construct a version of the Melvin's universe with a positive cosmological constant Λ , with the line element and gauge field given by [3]

$$ds^2 = -dt^2 + d\rho^2 + dz^2 + \sigma^2 \sin^2(\sqrt{2\Lambda}\rho) d\phi^2$$

$$A = \sqrt{\Lambda} z\mathrm{d}t \quad \Rightarrow \quad F = \sqrt{\Lambda}\mathrm{d}z \wedge \mathrm{d}t$$

We consider a charged scalar field minimally coupled to the Melvin's background, with a twist deformation resulting from two Killing vector fields that we introduced. Assuming that the noncommutativity parameter is small, we use the Seiberg-Witten map to obtain modified dynamical equation for the scalar field differing from the classical version by the term

$$-\frac{\bar{k} qE\left(\frac{\partial^2 f}{\partial t\, \partial \phi} - iqEz\frac{\partial f}{\partial \phi}\right)}{\left(1 + \frac{E^2\rho^2}{4}\right)^2}$$

The resulting equation can be written as an undeformed classical equation for the *effective* metric

$$G_{\mu\nu} = \begin{pmatrix} -\left(1+\frac{E^2\rho^2}{4}\right)^2 & 0 & 0 & \frac{\bar{k} qE\rho^2}{2\left(1+\frac{E^2\rho^2}{4}\right)^2} \\ 0 & \left(1+\frac{E^2\rho^2}{4}\right)^2 & 0 & 0 \\ 0 & 0 & \left(1+\frac{E^2\rho^2}{4}\right)^2 & 0 \\ \frac{\bar{k} qE\rho^2}{2\left(1+\frac{E^2\rho^2}{4}\right)^2} & 0 & 0 & \frac{\rho^2}{\left(1+\frac{E^2\rho^2}{4}\right)^2} \end{pmatrix}$$

In the case of positive cosmological constant, similar procedure can be done, with the resulting effective metric of the form

$$G_{\mu\nu} = \begin{pmatrix} -1 & 0 & 0 & 2\bar{k}\sqrt{\Lambda} q\sigma^2 \sin^2\left(\sqrt{2\Lambda}\rho\right) \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 2\bar{k}\sqrt{\Lambda} q\sigma^2 \sin^2\left(\sqrt{2\Lambda}\rho\right) & 0 & 0 & \sigma^2 \sin^2\left(\sqrt{2\Lambda}\rho\right) \end{pmatrix}.$$

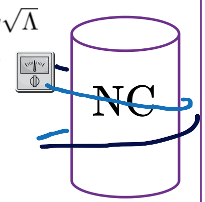
The conclusion is that the effect of our twist is to make the scalar fields "feel" like being in a noninertial rotating reference frame. The same can be shown to be true for Dirac fermions.

NC angular velocity of the frame is obtained to be

$$\Omega_{\mathrm{NC}} = -\frac{1}{2} qE\bar{k} \qquad \Omega_{\mathrm{NC}} = -\tfrac{1}{2}\bar{k} q\sqrt{\Lambda}$$

With a stationary cylindrical interferometer in NC Melvin's spacetime we can test the Sagnac effect resulting from this quantum-gravity-induced rotation felt by charged particles.

$$\Delta\phi = \frac{2mq\Phi_E \bar{k}}{\hbar^2}$$



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Elias Rothlin

Quantum Reference Frames for Quantum Error Correction with lattice QED

ETH zürich

WOST Exploring Quantum Without SpaceTime

Error Correction in Lattice Quantum Electrodynamics with Quantum Reference Frames

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IQOQI Vienna

FWF Austrian Science Fund

Motivation

- Is gauge symmetry merely a redundancy in our description, or does it carry a deeper information-theoretic significance?
- Quantum error-correcting codes (QECCs) use redundancy to protect information against noise. Previous works have established a bridge between gauge systems, stabilizer QECCs and quantum reference frames (QRFs) [1, 2].
- As an example, we study lattice quantum electrodynamics (QED) as a QECC, construct QRFs and determine correctable error sets.
- Using Gauss' law as a resource for quantum error correction has applications for fault-tolerant simulations of lattice gauge theories [3].

Gauge systems as generalized stabilizer codes

- **Gauge systems:** Kinematical space contains physical states ($\mathcal{H}_{\text{phys}} \subset \mathcal{H}_{\text{kin}}$) which are invariant under gauge transformations, $U(g)|\psi\rangle_{\text{phys}} = |\psi\rangle_{\text{phys}}$ for all $g \in G$.
- **Stabilizer codes:** Code states are encoded in a physical system ($\mathcal{H}_{\text{code}} \subset \mathcal{H}_{\text{physical}}$), invariant under the stabilizer group $S|\psi\rangle_{\text{code}} = |\psi\rangle_{\text{code}}$ for all $S \in \mathcal{S}$.
- **Identifications:**

Gauge system		Stabilizer code
$\mathcal{H}_{\text{phys}}$	$\leftrightarrow$	$\mathcal{H}_{\text{code}}$
$\mathcal{H}_{\text{kin}}$	$\leftrightarrow$	$\mathcal{H}_{\text{physical}}$
$U(g), g \in G$	$\leftrightarrow$	$S \in \mathcal{S}$

Based on the decomposition $\mathcal{H}_{\text{kin}} = \bigoplus_{\boldsymbol{q}} \mathcal{H}_{\boldsymbol{q}}$, a charge (syndrome) measurement can be used to correct errors. If $\{A_{\boldsymbol{q}}\}_{\boldsymbol{q}}$ are unitary on $\mathcal{H}_{\text{kin}}$ and satisfy $A_{\boldsymbol{q}}(\mathcal{H}_{\text{phys}}) \subset \mathcal{H}_{\boldsymbol{q}}$, the channel (see Fig. 1)

$$\sum\nolimits_{\boldsymbol{q}} A_{\boldsymbol{q}}^\dagger \Pi_{\boldsymbol{q}} (\cdot) \Pi_{\boldsymbol{q}} A_{\boldsymbol{q}}$$

corrects any error E for which $A_{\boldsymbol{q}}^\dagger \Pi_{\boldsymbol{q}} E|_{\text{phys}} \propto I_{\text{phys}}$.

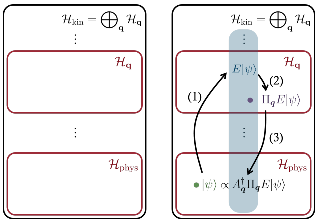

Figure 1. Left: $\mathcal{H}_{\text{kin}}$ decomposes into a direct sum of charge sectors $\mathcal{H}_{\boldsymbol{q}}$, where $\mathcal{H}_{\text{phys}}$ corresponds to the trivial representation. Right: (1) An error E maps a physical state $|\psi\rangle \in \mathcal{H}_{\text{phys}}$ to the error state $E|\psi\rangle$ which may have support spread across all charge sectors. (2) A measurement collapses the state to one of the charge sectors, $\Pi_{\boldsymbol{q}} E|\psi\rangle \in \mathcal{H}_{\boldsymbol{q}}$. (3) Applying an operator $A_{\boldsymbol{q}}^\dagger$ recovers the original state if $A_{\boldsymbol{q}}^\dagger \Pi_{\boldsymbol{q}} E|\psi\rangle \propto |\psi\rangle$.

QRFs and correctable errors

Perspective-neutral framework of QRFs [4, 5]: $\mathcal{H}_{\text{kin}} = \mathcal{H}_R \otimes \mathcal{H}_S$, and the QRF R parametrizes the action of G with covariant orientation states $|\phi(g)\rangle_R$ which resolve identity.

- **Correctable gauge-fixing operators:** For G compact, the gauge-fixing operators $\mathcal{P}_R^g \propto |\phi(g)\rangle\langle\phi(g)|_R \otimes I_S$ for orthogonal orientation states are correctable, $\Pi_{\text{phys}} \mathcal{P}_R^g \mathcal{P}_R^h \Pi_{\text{phys}} = \delta(g,h) \Pi_{\text{phys}}$.
- **Recoveries via charge-sector measurements:** For G Abelian and R an ideal QRF (all orientation states orthogonal), the gauge-fixing operators yield $\{A_{\boldsymbol{q}}\}_{\boldsymbol{q}}$ as in Fig. 1:

$$A_{\boldsymbol{q}} \propto \int \mathrm{d}g \, \overline{\chi}_{\boldsymbol{q}}(g) \mathcal{P}_R^g.$$

  If non-ideal frames are ideal for subgroups $H \subset G$, this yields correctable errors through coarse-grained charge sectors.

Lattice quantum electrodynamics

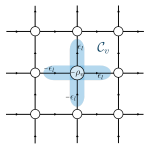

Figure 2. In Hamiltonian lattice QED, space is discretized to a lattice. On every site v, physical states satisfy $\mathcal{C}_v|\psi\rangle_{\text{phys}} = 0$.

In the Hamiltonian formulation of lattice QED, time is kept continuous ($A_0 = 0$) and space discretized to a lattice with vertices/sites $v \in \mathcal{V}$ and oriented links $l \in \mathcal{L}$.

- **Pure gauge sector:** Each link carries $L^2(\mathrm{U}(1))$, with group basis $\{|e^{i\theta}\rangle_l\}_{[0,2\pi)}$ and Fourier-transformed electric flux basis $\{|k\rangle_l\}_{k\in\mathbb{Z}}$.
  - link operator: $U_l|e^{i\theta}\rangle_l = e^{i\theta}|e^{i\theta}\rangle_l$
  - electric flux operator: $\epsilon_l|k\rangle_l = k|k\rangle_l$

  Gauss' law constraints annihilate physical states, $\mathcal{C}_v|\psi\rangle_{\text{phys}} = 0$, where

$$\mathcal{C}_v = \sum_{l\in\text{out}(v)} \epsilon_l - \sum_{l\in\text{in}(v)} \epsilon_l.$$

  Gauge transformations are of the form $\Pi_{v\in\mathcal{V}} e^{i\lambda_v \mathcal{C}_v}$.
- **Including staggered fermions:** (Anti-)Fermions are placed on even (odd) sites $v \in \mathcal{V}$, with creation and annihilation operators $\psi_v^\dagger, \psi_v$ satisfying

$$\{\psi_v, \psi_{v'}^\dagger\} = \delta_{v,v'}, \ \{\psi_v, \psi_{v'}\} = \{\psi_v^\dagger, \psi_{v'}^\dagger\} = 0.$$

  Each site is spanned by the number basis $|0\rangle_v, |1\rangle_v$ Gauss' law includes the staggered charge density $\rho_v = \psi_v^\dagger \psi_v - \frac{1}{2}(1 - (-1)^{|v|})$, i.e.,

$$\mathcal{C}_v = \sum_{l\in\text{out}(v)} \epsilon_l - \sum_{l\in\text{in}(v)} \epsilon_l - \rho_v.$$

QRFs for lattice QED

For lattice QED, we find two types of QRFs according to the perspective-neutral framework (see Fig. 3).

- **Spanning tree QRFs:** The links of a spanning tree (a loop-free subgraph containing every vertex) form an ideal QRF for the pure-gauge sector, with orientation states

$$|\phi(\boldsymbol{\lambda})\rangle_R = \bigotimes_{l\in R} |e^{i\lambda_l}\rangle_l.$$

  Gauge-fixing the orientation maps physical d.o.f. into non-tree links $S = \mathcal{L} \setminus R$. Each $l \in S$ labels a fundamental holonomy belonging to the loop on R closed by l.
- **Fermionic field QRF:** Including staggered fermions, the fermionic field on the sites forms a complete but non-ideal QRF $\tilde{R}$, with orientation states

$$|\boldsymbol{\lambda}\rangle_{\tilde{R}} = \bigotimes_{v\in\mathcal{V}} \tfrac{1}{\sqrt{2}}(|0\rangle_v + e^{-i\lambda_v}|1\rangle_v).$$

  Physical states can be expressed purely in terms of electric flux, and the fermions are recoverable from Gauss' law.

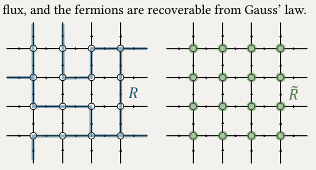

Figure 3. Left: In the pure-gauge sector, the links of a spanning tree form an ideal QRF R. Right: Including fermions, the fermionic field on the sites yields a complete but non-ideal QRF $\tilde{R}$.

Lattice QED as a QECC

Interpreted as a QECC, two types of codes arise from lattice QED:

- A **quantum rotor code** from the pure gauge sector, encoding $|\mathcal{L}| - |\mathcal{V}| + 1$ rotors ($L^2(\mathrm{U}(1))$-spaces) into $|\mathcal{L}|$ rotors with U-distance $d_U = 4$.
- A **hybrid rotor-qubit code** from lattice QED with fermions.

Correctable errors in lattice QED

We find correctable sets of gauge-violating errors for the two types of codes above. They have recoveries based charge-sector (constraint) measurements.

- **Pure gauge lattice QED:** Gauge-violating operators consist of products of link operators (open Wilson lines).
  - A spanning tree QRF R yields the correctable errors

$$\left\{ \bigotimes\nolimits_{l\in R} U_l^{m_l} \,\middle|\, \forall l \in R : m_l \in \mathbb{Z} \right\}.$$

    These are products of Wilson lines supported on the spanning tree which shift the electric flux on the corresponding links.
  - Since $d_U = 4$, a different recovery than for spanning trees can correct any single U_l^m-error, i.e., any shift in electric flux on a single link.
- **Fermionic lattice QED:** Correctable errors from the pure gauge case carry over. Additionally, $\psi_v, \psi_v^\dagger$ and combinations thereof yield gauge-violating errors.
  - The fermionic field QRF $\tilde{R}$ yields correctable errors via coarse-grained charge measurements:

$$\left\{ \bigotimes\nolimits_{v\in\mathcal{V}} A_v(\alpha_v)^{r_v} \,\middle|\, \forall v : r_v \in \{0,1\} \right\}.$$

    Here, $A_v(\alpha_v) = e^{i\alpha_v}\psi_v + e^{-i\alpha_v}\psi_v^\dagger$ implement particle number flips (with phases). E.g., fixing $\alpha_v = 0$, these are arbitrary products of Pauli-X errors.
  - Combining coarse-grained and standard charge-sector measurements:

$$\{U_l^m\}_{m\in\mathbb{Z}, l\in\mathcal{L}} \cup \{A_v(\alpha_v)\}_{v\in\mathcal{V}},$$

    i.e., any single electric flux error on a link or particle number flip on a site is correctable.

Conclusions

Gauge systems can be viewed as QECCs. We show that:

- In the Abelian case, QRFs yield sets of correctable errors via charge measurements.
- Lattice QED has ideal QRFs given by spanning trees in the pure-gauge sector, and complete but non-ideal QRFs from the fermionic field when including staggered fermions.
- Based on charge (constraint) measurements, certain electric flux error can be corrected.
- Including staggered fermions, particle number flips can be corrected through coarse-grained charge measurements. This can be combined with correcting electric flux errors.

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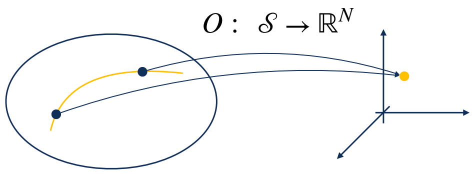
All **local information** can be **completely** encoded into **global functions.**

Observables are glocal

Emil Broukal, Andrea Di Biagio, Eugenio Bianchi, and Marios Christodoulou

Complete observables

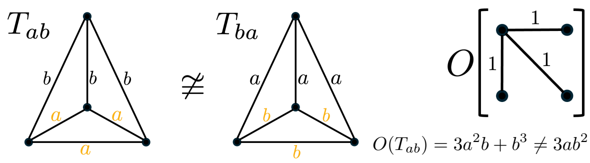


Collection of spacetimes / Space of weighted graphs

- The **gravitational field** has **local degrees of freedom.**
- **Diffeomorphism invariant** functions of the metric are expected to be **global**.
- All **local information** must be **completely** encoded into **global objects.** How precisely does this take place?

- **Setup**: Weighted graphs with permutations of node labels as discrete coordinate changes.

Local correlations encoded in background independent observables



$$O(T_{ab}) = 3a^2b + b^3 \neq 3ab^2 + a^3 = O(T_{ba})$$

- Fully **background independent** observable **can distinguish** configurations with **same global** information, but **different local** structure.
- A **glocal observable** O^M probes a **local** correlation specified by a **connected** multigraph M, by taking a **global** average of all possible appearances of this local subgraph structure. This achieves **full background independence.**

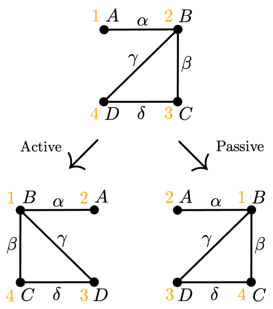
Glocal observables are complete

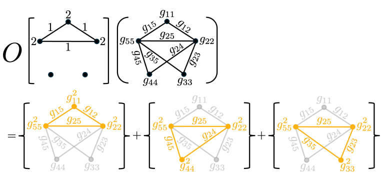
$$O^M = \frac{1}{|S_N|} \sum_{\sigma \in S_N} \left(\prod_{i \le j}^{N} x_{\sigma(i)\sigma(j)}^{m_{ij}} \right), \quad I_C^N = \left\{ O^M \mid W(M) \le \begin{pmatrix} \frac{1}{2}N(N+1) \\ 2 \end{pmatrix} \right\}$$

- **Theorem:** I_C^N **generates** the **full algebra of permutation-invariant polynomials** on weighted graphs up to N nodes.
- A **finite** number of **glocal observables** captures **complete** information about **all** kinds of **local correlations** on **any** finite weighted graph. In this sense, we can say that **all information** on a weighted graph is **glocal information.**

- **Why weighted graphs?**
Many approaches to quantum gravity use graphs as the topological basis on which they build their theory. Notable examples: **Loop Quantum Gravity**, **Causal Sets**, and **Tensor Models**.

- **Why permutations?**
Node labels are discrete **coordinates** on a graph. **Permutations** of the node labels are **coordinate changes.**





$$= g_{11}^2 g_{12} g_{22}^2 g_{25} g_{55}^2 g_{15} + g_{22}^2 g_{24} g_{44}^2 g_{45} g_{55}^2 g_{25} + g_{22}^2 g_{23} g_{33}^2 g_{35} g_{55}^2 g_{25}$$

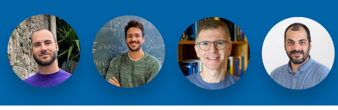
Extension to LQG:
- Unlabelled spin networks:

$$|\mathcal{O}_\Gamma\rangle = \frac{1}{\sqrt{|\mathcal{O}_\Gamma|}} \sum_{H \in \mathcal{O}_\Gamma} |H\rangle \in \mathrm{Inv}_{S_N} \mathscr{H}_N^{\mathrm{kin}}$$

$$\langle \mathcal{O}_\Gamma | \mathcal{O}_{\Gamma'} \rangle = \begin{cases} 1, \text{ if } \Gamma \cong \Gamma' \\ 0, \text{ if } \Gamma \not\cong \Gamma' \end{cases}$$

- Hilbert space spanned by unlabelled spin networks is a discrete and quantum analogue of imposing the spatial diffeomorphism constraint of general relativity.

Take a picture to download paper.



IQI Vienna

Emilien de Bank

A Fine-Grained Perspective on Higher Order Operations: Modelling Agents in Spacetime

A Fine-Grained Perspective on Higher Order Operations: Modelling Agents in Spacetime

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Abstract

Causality is a fundamental concept that takes on markedly different forms in **quantum and relativistic theories**. Yet in quantum information protocols implemented in spacetime, both notions must coexist consistently. This raises compelling questions about the physicality of indefinite causal order (ICO) processes, where the ordering of agents' operations is not fixed or acyclic.

We study a **fine-grained framework**, in which such protocols are embedded in spacetime in a way that respects relativistic causality. Our modelling of quantum measurements in spacetime informs how agents' interventions can be described within the fine-grained picture and still capture the operational correlations.

Furthermore, considering recent **experimental claims** for certifying non-classical correlations in ICO processes, we analyse how they align with a fine-grained, spacetime-consistent interpretation which can involve superpositions of arrival times [1]. This can enable a better understanding of the physical resources responsible for these non-classical features in the **fine-grained picture** which accounts for both quantum information and relativistic aspects.

Indefinite causal order

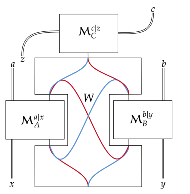

Fig. 1: The quantum switch [2].The quantum switch does not violate causal inequalities.

◇ The process matrix framework can describe circuits with **indefinite causal order** [3].

→ Generalised Born Rule

$$p(abc|xyz) = \mathrm{Tr}\left[\left(M_A^{a|x} \otimes M_B^{b|y} \otimes M_C^{c|z}\right)^T W\right].$$

◇ The canonical example is **the quantum switch**.

◇ **Causal inequalities** are introduced as analogs of Bell inequalities, under the assumptions of

→ Free Choice,

→ Closed Laboratories,

→ The presence of an acyclic causal structure in which events are localised.

Goal

◇ This work aims at providing **measurements models** for systems non-localized in spacetime,

◇ and characterises **achievable correlations** in fine-grained process realisations under different measurement model.

◇ It is used to study **loopholes** in optical experiments that implement process matrix scenarios.

◇ We introduce a resource called **no classical localisation** of states and operations, distinct from ICO, that correlations certify at the fine-grained level.

Process boxes

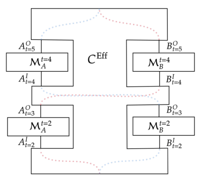

Fig. 2: Representation of the quantum switch in the process box framework.

◇ A framework that models processes in **fixed and acyclic spacetime** [4, 5, 6, 7].

◇ Spacetime is represented by a partially ordered set $\mathcal{T}$.

◇ Hilbert spaces are replaced by *Fock spaces*

$$\mathcal{F}_A^{\mathcal{T}} = \bigotimes_{m=0}^{\infty} \vee^m(\mathcal{H}_A \otimes \ell^2(\mathcal{T})).$$

◇ Assumptions are necessary to enforce a **single round scenario** between agents,

→ Acting Once

$$|\Psi\rangle^{A_k} \in Span\Big(|j_k, t_k\rangle \otimes |\Omega, \mathcal{T}_k^{I/O}\backslash\{t_k\}\rangle\Big)_{j_k, t_k \in \mathcal{T}_k^{I/O}}.$$

→ Local order

$$M_{A_k}|\psi, t\rangle|\Omega, \mathcal{T}^I\backslash\{t\}\rangle = |\phi_t, t+1\rangle|\Omega, \mathcal{T}^O\backslash\{t+1\}\rangle.$$

Result I : Process boxes and QC-QCs

- Quantum Circuit with Quantum Control of causal order (QC-QCs) are known not to violate causal inequalities
- For all process boxes there exists a QC-QC whose behaviour is equivalent (At the level of the probability distribution) [4, 5, 6, 7]. We provide

→ an alternative mapping that clarifies how time of arrival information can act as control.

→ trace preserving conditions, that are completely analogous to those satisfied by QC-QCs, for process boxes.

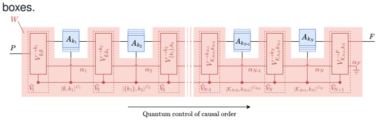

Fig. 3: Quantum Circuit with Quantum Control of causal order (QC-QC) [8].

Result II: Measurement models

◇ An agent's operation $\mathcal{M}_A$ receives a classical bit x and sends a classical bit a,

$$\mathcal{M}_{a_k|x_k}^{A_k} := (\mathcal{I}_{A_k^O} \otimes \pi_{a_k}^{t=t_F}) \circ \mathcal{M}_{x_k}^{A_k} \quad \text{with} \quad \mathcal{M}_{x_k} = \bigotimes_{t<t_F} \mathcal{M}_{x_k}^t.$$

◇ In spacetime, agents must coherently write possible outputs on a quantum memory that is **measured at the end** of the process.

◇ Projecting at the very end *erases the which way information* and gives correlations compatible with a controlled superposition of causal order.

◇ We show a correspondence between time decoherence for process boxes and control decoherence for QC-QCs through our measurement models.

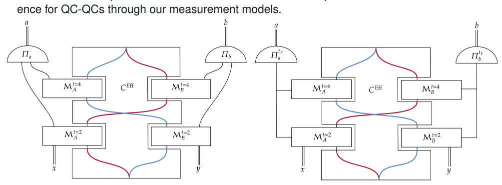

Fig. 4: The quantum switch with Alice and Bob's quantum instruments unraveled in spacetime for two different measurement models. The measurement on the left decoheres the time of arrival d.o.f. while the right one preserves it.

Discussion

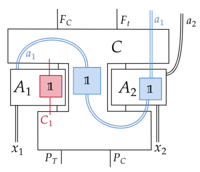

Fig. 5: Optical implementation of the quantum switch [9].

◇ The process box framework allows us to identify and clarify potential loopholes in physical implementations.

→ For instance, the outcome of the memory might travel spatially through the process.

→ The control might also be accessible to the agents.

◇ In the process box framework, certification protocols (e.g. causal witness [10] or DRF inequalities [11]) attest the *No Classical Localisation* of states and operations.

Outlook: Perspective transformation (WIP)

◇ We are developing a perspectival version of process boxes relative to agents' reference frames, and describe transformations between fine-grained perspectives.

◇ Process box protocol $(\mathcal{C}, \{\mathcal{M}_{a|x}^A, \mathcal{M}_{b|y}^B\}_{a,b,x,y})$ where Alice is localised i.e. she acts at a fixed time step.

◇ Unitary transformation to change from Alice's to Bob's perspective as in [12],

$$|J_{A\to B}\rangle\rangle = |00\rangle^{P_C F_C} \mathcal{T}^{-2}_{A^I A^O B^I B^O} + |11\rangle^{P_C F_C} \mathcal{T}^{+2}_{A^I A^O B^I B^O}.$$

◇ The transformation delocalises Alice's states and operations and localises Bob ones.

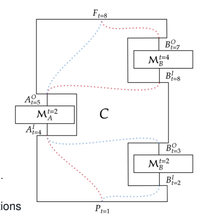

Fig. 6: The quantum switch where Alice is localised in the process box framework.

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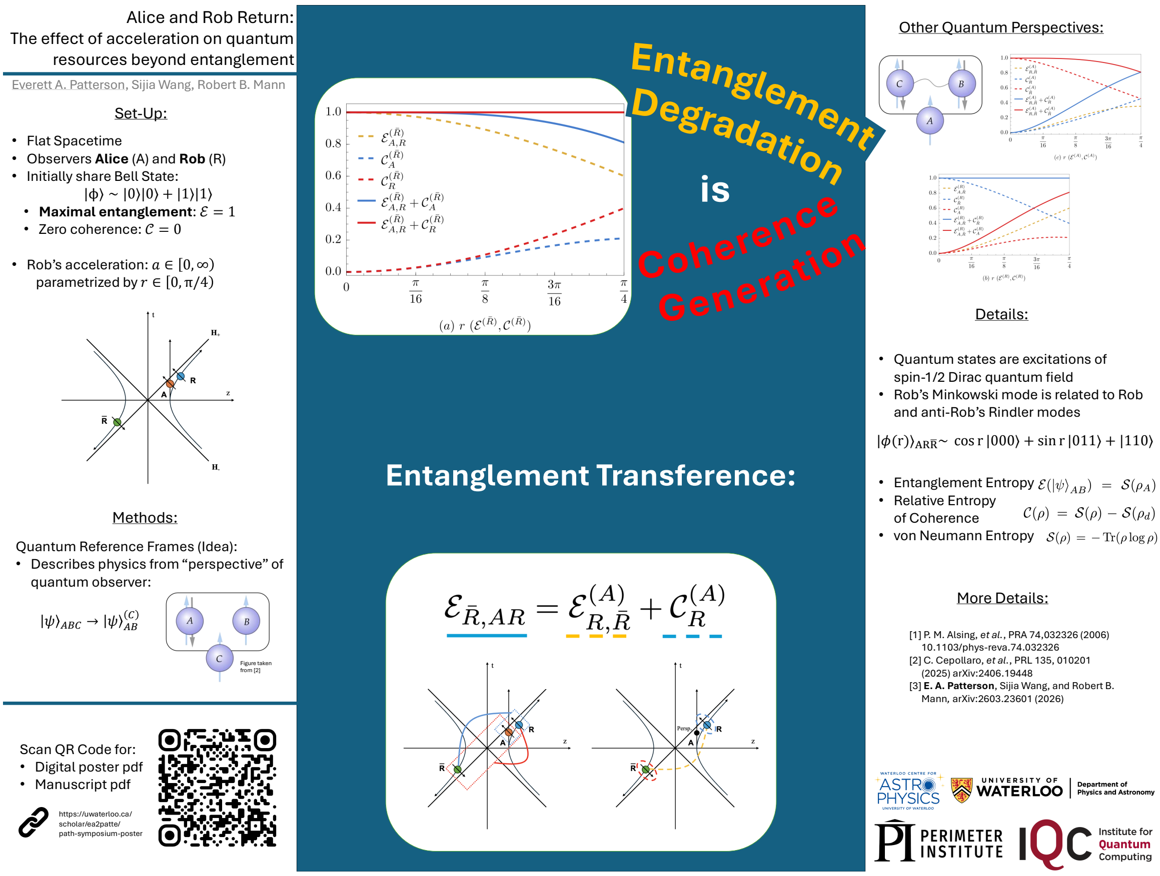
Everett A. Patterson

Perspectival Entanglement Degradation: What Alice and Rob might see using quantum reference frames



Filip Pozar

Isometry and gauge invariance on quantum spaces implies metric signature

Quantum spacetime + gauge symmetry = metric signature

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References: [1] - arXiv:2503.07176 , [2] - arXiv:0901.2710

Acknowledgements: Croatian Science Foundation Project No. IP-2025-02-8625, "Quantum aspects of gravity"

Introduction

Noncommutative geometry in physics is used to model "manifolds" with intrinsic minimal length scales. Often it is incorporated by deforming the algebra of smooth spacetime functions $C^\infty(M)$ by systematically promoting the pointwise function product $\cdot$ into a noncommutative product $\star$. This $\star$ product imposes noncommutativity of coordinate functions $[x^\mu, x^\nu]_\star = x^\mu \star x^\nu - x^\nu \star x^\mu \neq 0$ and propagates to an uncertainity relation of position measurement. As a mathematical field, noncommutative geometry studies geometric constructions on noncommutative algebraic structures (noncommutative algebras and modules over them).

In this presentation we show how deformations which impose (almost Abelian) Lie algebraic commutation relations on coordinates satisfy a unique property that $U(N)$ gauge invariance (of a natural choice of an action) is possible only on specific metric signatures for a given Lie algebraic deformation.

1 - Deformed product

Consider an almost abelian Lie group

$$\mathcal{G} = \mathbb{R} \ltimes_a \mathbb{R}^n \quad a \in \mathrm{Mat}_{n\times n}$$

whose group product is given as [1]

$$(p^M, \vec{p})\,(q^M, \vec{q}) = (p^M + q^M, \vec{p} + a(p^M)\vec{q})$$

for $p^M, q^M \in \mathbb{R}$ and $\vec{p}, \vec{q} \in \mathbb{R}^n$. Then, for $A(p^M) = a(p^M) \oplus 1_{1\times 1} \in \mathrm{Mat}_{n+1\times n+1}$ one can consider a star product on the $n+1$ dimensional spacetime as

$$(f \star g)(x) = \frac{1}{(2\pi)} \int dp^M dy^M e^{-ip^M y^M} f(x + y^M) g(A(p^M)x)$$

This deformed product imposes the commutators on coordinates as

$$[x^M, x^\mu]_\star = -i\left[\partial_{p^M} A(p^M)|_{p^M=0}\right]^\mu{}_\sigma x^\sigma \, .$$

2 - Multi-derivation structure

It turns out that partial derivatives form twisted multi-derivations [2] of the deformed algebra $C^\infty_\star(M)$ in the sense

$$\partial_\mu(f \star g) = (\partial_\mu f) \star g + A(P_M)^\rho{}_\mu(f) \star \partial_\rho g$$

where $P_M = -i\frac{\partial}{\partial x^M}$

3 - Gauge structure

In noncommutative geometry one can define a connection over a noncommutative algebra's right module M. M is a right module over an algebra A is there exists a map $\lhd \colon A \to \mathrm{Aut}(M)$ which obeys

$$(m \lhd a) \lhd b = (m) \lhd (ab)\ , \forall a, b \in A, m \in M$$

The module M models sections of vector bundles, as commutatively they form a right module over spacetime functions $C^\infty(M)$ ($(\Psi \lhd f)(x) = \Psi(x) \cdot f(x)$). On M one can define a covariant derivative in the direction of a derivation X **(let us for now forget that partial derivatives are not derivations of the algebra in our almost abelian case)** of the algebra A as any map $\nabla_X : M \to M$ satisfying the Leibniz rule

$$\nabla_X(m \lhd a) = m \lhd X(a) + \nabla_X(m) \lhd a \qquad (1)$$

where derivations X of the algebra A are linear maps $X : A \to A$ satisfying

$$X(ab) = X(a)b + aX(b)$$

Gauge transformations are defined as automorphisms of left module M, i.e., invertible linear maps ϕ which are also right A-linear

$$\phi(m \lhd a) = \phi(m) \lhd a$$

This exactly reproduces the $C^\infty(M)$ linearity of gauge transformations in the commutative case.

With that said, $\nabla^\phi := \phi \circ \nabla_X \circ \phi^{-1}$ satisfies the same Leibniz rule as ∇_X. Additionally, it satisfies $\phi\nabla_X(m) = \nabla^\phi(\phi m)$. One can define curvature of the connection as the right A-linear map for any two derivations X,Y as

$$F(X,Y)(m) = [\nabla_X, \nabla_Y](m) - \nabla_{[X,Y]}(m)\, .$$

It can be shown that $F^\phi(X,Y) = \phi \circ F(X,Y) \circ \phi^{-1}$ is the curvature of ∇^ϕ_X.
To model electromagnetism's $U(1)$ connection, we actually choose the module to equal the algebra itself and we restrict to gauge transformations given as multiplication by unitary elements g. One can show, via the Leibniz rule, that any connection is comletely determined by its action on the unit element of the algebra

$$\begin{aligned}\nabla_X(a) = \nabla_X(1 \star a) &= 1 \star X(a) + \nabla_X(1) \star a \\ &= X(a) + i\, A_X \star a\end{aligned}$$

where A_X can be thought of the contraction of the vector field X and gauge potential A. Similarly, curvature's components come from acting on the identity $F_{\mu\nu} \equiv F(\partial_\mu, \partial_\nu)(1)$ and we have

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - i\,[A_\mu, A_\nu]_\star$$

Using this curvature, we define an action principle

$$S = \int d^4x F_{\mu\nu} \star F^\dagger_{\rho\sigma} \star g^{\mu\rho} \star g^{\nu\sigma}$$

Results

A complication that we encountered is that the Leibniz rule (1) is well defined only for derivations X, but, as we have announced in Section 2, partial derivatives are actually twisted multi-derivations, not derivations. Thus, the concept of the connection needs to be adapted. (1) is impossible to consistently satisfy, but, the following Leibniz rule

$$\nabla_\mu(a \star f) = A(\lambda P_M)^\rho{}_\mu(a) \star \partial_\rho(f) + \nabla_\mu(a) \star f \quad (2)$$

is satisfied by connections parametrized by $A_\mu = \nabla_\mu(1)$ in the following usual way

$$\nabla_\mu(a) = \partial_\mu(a) + iA_\mu \star a\, . \qquad (3)$$

Additionally, for a unitary automorphism $\phi_g(a) = g \star a$ the concept of gauge transformations needs to be deformed as well. The correct choice amounts to

$$\nabla^g_\mu(a) = A(\lambda P_M)^\rho{}_\mu(g) \star \nabla_\rho(g^\dagger \star a) \qquad (4)$$

which affects the transformation rules of the potential

$$A^g_\mu = \nabla^g_\mu(1) = A(P_M)^\rho{}_\mu(g) \star \left(\partial_\rho g^\dagger + A_\mu \star g^\dagger\right)$$

and of the curvature

$$F^g_{\mu\nu} = g \star F_{\alpha\beta} \star A^{-1}(\lambda P_M)^\alpha{}_\mu A^{-1}(\lambda P_M)^\beta{}_\nu(g^\dagger)$$

The condition for gauge invariance $S^g = S$ gives finally a condition on $A(\lambda P_M)$ and $g^{\mu\nu}$ as a relation which has hold for all unitary elements $h \in C^\infty_\star(M)$

$$g^{\mu\rho} g^{\nu\sigma} (A^{-1})^\xi{}_\rho (A^{-1})^\chi{}_\sigma(h^\dagger) \star A^\alpha{}_\mu A^\beta{}_\nu(h) = g^{\xi\alpha} g^{\chi\beta}\, . \qquad (5)$$

ρ and hyperbolic ρ Minkowski spacetimes:

$$A(P_0) = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \cos(P_0) & \sin(P_0) & 0 \\ 0 & -\sin(P_0) & \cos(P_0) & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}$$

$$A(P_3) = \begin{pmatrix} \cosh(P_3) & \sinh(P_3) & 0 & 0 \\ \sinh(P_3) & \cosh(P_3) & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}$$

solve (5) for

$$g(\rho)^{\mu\nu} = {}_{\mathrm{diag}}(-1, |b|, |b|, 1)\ , \quad b \in \mathbb{R}$$

$$g(H\rho)^{\mu\nu} = {}_{\mathrm{diag}}(-|b|, |b|, 1, 1)\ , \quad b \in \mathbb{R}$$

Triangular space

$$A(P_3) = \begin{pmatrix} 1 & 0 & 0 & 0 \\ P_3 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}$$

is solved by Galilean metric $\mathrm{diag}(0, 1, 1, 1)$.

Time of arrival on a ring and relativistic quantum clocks

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Abstract

We study the time-of-arrival problem for relativistic particles constrained to move on a ring, formulating the problem entirely within Quantum Field Theory (QFT). We employ the Quantum Temporal Probabilities (QTP) method to derive a class of Positive-Operator-Valued Measures (POVMs) for time-of-arrival observables directly from QFT.

Time-of-arrival probabilities

For a relativistic particle with $\omega_m^2 = \mu^2 + \left(\frac{m}{r}\right)^2$ moving in a ring the time-of-arrival probability is

$$P_c(t,\varphi) = \frac{B}{2\pi r} \sum_{m,m'} \rho_{ps}(m,m') \langle m | \hat{L} | m' \rangle \sqrt{v_m v_{m'}} e^{i(m-m')\varphi - i(\omega_m - \omega_{m'})t},$$

- B: normalization constant.
- $v_m = m/(\omega_m r)$: *linear velocity* of the particle.
- $\langle m | \hat{L} | m' \rangle$: detection event localization operator.
- $\rho_{ps}(m,m')$: post-selected density matrix.

Rotating observers

For a detector rotating with angular velocity Ω_D the background noise increases as $\Omega_D r \to 1$.

$$\eta = \frac{P_0(\Omega_D)}{P_0(0)} = \frac{\log\left(1 - e^{-a(1-\Omega_D r)}\right)}{\log(1 - e^{-a})}$$

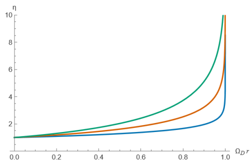


Figure 1: Noise ratio η as a fucntion of $\Omega_D r$.

The change of the background noise is a manifestation of the *circular Unruh effect*.

Quantum clocks and quantum recurrence

Time-of-arrival probabilities define a periodic clock variable on a fundamentally quantum system.

1. A large number of particles prepared in a state sharply localized in position and momentum passing through a detector behave as a quantum clock. Each peak in the probability density $W(t) = \int_0^t ds P_c(s,\varphi)$ of the time of arrival defines a tick of the clock.
2. Because of the discrete spectrum of the Hamiltonian, the probability amplitude is subject to the *quantum recurrence theorem*. We define two relevant time scales:
   - the *quantum time scale* T_q, where the spreading becomes significant and ticks become superimposed.
   - the *recurrence time scale* T_{rec}, where the wavepacket initial form is partially restored.

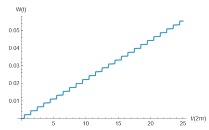
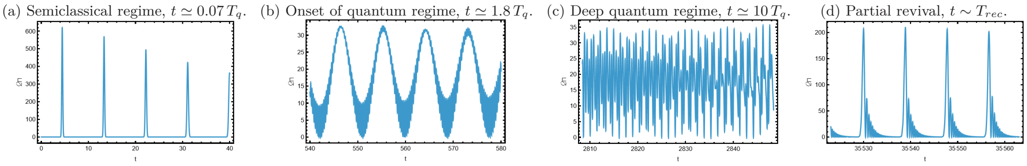


Figure 2: The cumulative probability density $W(t)$ for massless particles.

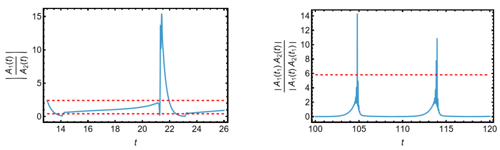
(a) Semiclassical regime, $t \simeq 0.07\, T_q$. (b) Onset of quantum regime, $t \simeq 1.8\, T_q$. (c) Deep quantum regime, $t \simeq 10\, T_q$. (d) Partial revival, $t \sim T_{rec}$.



Multi-time measurements

Measurement independence is a classical locality criterion, asserting that the measurements of two quantities are independent.

In this model, when considering multi-time measurements, classical measurement independence may be violated for entangled states when the joint probability amplitudes take certain values,



(a) Entangled particles in one ring. (b) Entangled particles in different rings.

where $\mathcal{A}_i(t,\varphi) = \sum_m \psi_i(m) \sqrt{v_m} e^{im\varphi - i\omega_m t}$.

Results

- Relativistic quantum clock based on time-of-arrival measurements in a ring.
- Entirely QFT treatment.
- Tool for probing quantum spacetime effects.
- Setting to explore quantum effects in clock readings.

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James Robinson

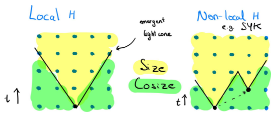
A formal refinement for operator size of qudits

A formal refinement for operator size theory of qudits

James Scott Robinson *supervised by* Dr. Ben Freivogel *for the completion of a master's thesis*

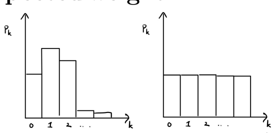
- **We study many-body teleportation** [1–4] as an experimental probe of the holographic principle
- **ETH:** *larger operators* are on average *worse* than *small operators* at distinguishing neighbouring eigenstates [5]
- **Holography:** operator size in the boundary is dual to the *momentum* (magnitude) of a particle in the bulk [6]

Growth under Heisenberg evolution



An initial perturbation (black) delocalizes over a one dimensional lattice.

Operator size is expected weight



- $\hat{\mathcal{O}} = \frac{1}{d}\sum_E c_E E/\|\mathcal{O}\|$, $\quad p_E = \frac{1}{d}|\langle E|\hat{\mathcal{O}}\rangle|^2$
- **Weight** $w(E)$, e.g. $w(X \otimes Y \otimes 1) = 2$
- **Weight distribution** p_k sum all p_E of equal $w(E) =: k$

Perfect size winding

Condition

$$c_E = e^{i\phi}|c_E| \quad \phi \propto \mathcal{S}(E)$$

SYK model

Gaussian random couplings

$$H = \sum_{ijkl} J_{ijkl}\, \psi_i\psi_j\psi_k\psi_l$$

Jordan-Wigner transform

Represent fermion on qubits

$$\mathcal{S}[\psi_k] \neq \mathcal{S}[\varphi_{\mathrm{JW}}(\psi_k)] = \mathcal{S}\left[(\otimes^{k-1}Z)\otimes X\right]$$

Chaos

Use $p = (q^2-1)/q^2$ with q qudit dimension

$$\mathbb{E}\,\mathcal{F}_t(B) = \langle\hat{\mathcal{O}}|\delta|\hat{\mathcal{O}}\rangle = 1 - \hat{\mathcal{S}}(\mathcal{O})/p$$

Operator correlator

$$G(\mathcal{O},\Phi) = \langle\hat{\mathcal{O}}|\Phi[\hat{\mathcal{O}}]\rangle$$

where Φ is a *quantum channel.* **Result:**

$$\hat{\mathcal{S}}(\mathcal{O}) = p - G(\mathcal{O},\Phi^{\times})$$

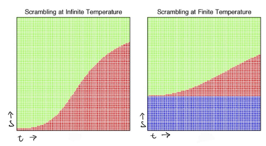
where $\delta\,|\mathcal{O}\rangle = |\Phi^{\times}[\mathcal{O}]\rangle$ and $\Phi^{\times} \propto \sum_k \mathcal{X}_k$.

Developed method

Relate **size** k **averages** $\mathbb{M}_k$ to **simple averages** $\mathbb{W}$:

$$\mathbb{M}_k(\mathcal{O}) \propto B\mathcal{O}B^{\dagger} \quad \mathbb{W}_K(\mathcal{O}) \propto \mathrm{Tr}_{k_j}(\mathcal{O})$$

Thermally renormalized size



Result:

$$\Delta_\beta\mathcal{S}_k \propto \sum_B \left[\mathcal{S}_\beta(B) - \mathcal{S}_\beta(\mathbb{1})\right] = k\delta(\sqrt{\rho_\beta})$$

We **need** a many body teleportation **experiment with fermions** because **operator size theory does not translate well**[a] to qubits

Additional insights

- The **Lyapunov exponent** λ determines the magnitude of an **operator correlator** $G(\mathcal{O},\Phi^{\times})$ over time
- The **thermally renormalized unit of size** for both fermions and qubits **is characterized by the scrambling distance**
- **Co-size** and **scrambling distance** become **equivalent** concepts as the **qudit dimension becomes large**

[a] using the **Jordan-Wigner transformation**

Further work

- Compute the *partition function* for size (related to higher order OTOCs)
- Operator size ≈ susceptibility to depolarization
  $\stackrel{?}{\Longrightarrow}$ **operator growth** ≈ (operator) **entanglement growth**?
- Operator growth of RMT Hamiltonians under depolarizing noise

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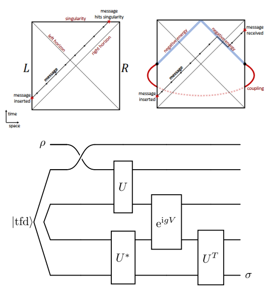
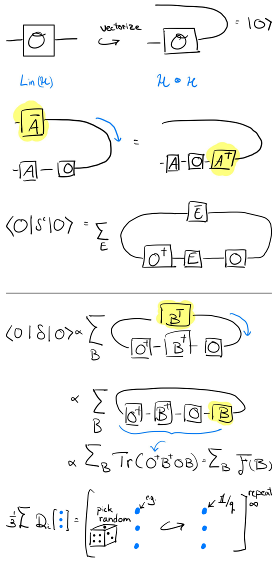
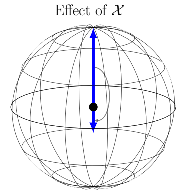
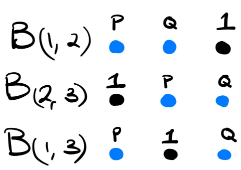




Effect of $\mathcal{X}$



$\mathbb{W}_1(\mathcal{O})$
$= \frac{1}{3}(\mathcal{O} + X^{\dagger}\mathcal{O}X + Y^{\dagger}\mathcal{O}Y + Z^{\dagger}\mathcal{O}Z) - \frac{1}{3}\mathcal{O}$



Measuring quantum time dilation with delocalised atoms

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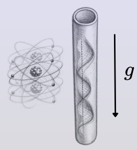
Motivation: clock superposition

- Classical theory: proper time is associated with a spacetime trajectory (worldline).
- Quantum theory: a clock can follow a superposition of trajectories.
- What is the time measured by a superposed clock? Is it a superposition of proper times?
- Are there any observable consequences of such a superposition?

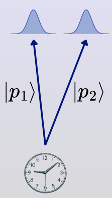
Simplifying assumptions

- One-dimensional photon propagation (atom coupled to a waveguide).
- Constant gravitational field (Rindler frame).
- Neglect the atom's kinetic energy (very massive atom; no recoil, no kinematic time dilation).



Quantum time dilation

- Consider two wave packets (momentum- or position-basis).
- Compare clocks in:
  - quantum superposition, or
  - classical (statistical) mixture.
- Do they measure the same time?
- No! A superposed clock experiences a quantum correction to the measured time dilation [1-4]!



Results: superposition vs mixture

Single wave packet:

$$\psi_i(z) = (2\pi\Delta^2)^{-1/4} e^{-(z-z_i)^2/2\Delta^2}$$

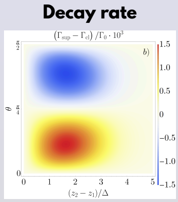
Quantum superposition:

$$\psi_{\mathrm{sup}}(z) = \mathcal{N}\left[\cos\theta\psi_1(z) + e^{i\varphi}\sin\theta\psi_2(z)\right]$$

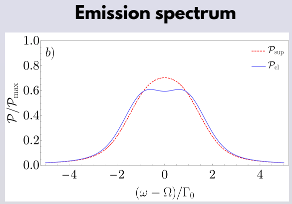
Classical mixture:

$$P_{\mathrm{cl}}(z) = (\cos\theta)^2\psi_1(z) + (\sin\theta)^2\psi_2(z)$$

Decay rate

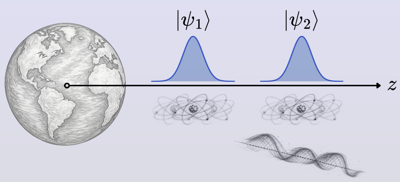


Emission spectrum



- Quantum time dilation affects both:
  - total decay rate, and
  - spectrum of emitted photons.
- The overlap between the wave packets is crucial (interference!).
- The effects might be observable in near-future experiments!

Our problem: spontaneous emission from a delocalized atom



- Two wave packets localized at different heights in a gravitational field.
- Gravitational time dilation → position-dependent spectrum of the atom.
- What is the spectrum of a superposed atom?
- How does quantum time dilation manifest in the spontaneous emission process?

Our work

Paczos, J., Dębski, K., Grochowski, P.T., Smith, A.R.H., Dragan, A., "Quantum time dilation in a gravitational field", Quantum 8, 1338 (2024)

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Quantum 8, 1338 (2024)
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Jorge Escandón-Monardes

Multiparameter estimation with a photonic quantum switch

Multiparameter estimation with a photonic quantum switch

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1. Introduction

Several experiments have demonstrated the advantages that indefinite causal order offers for quantum information. For instance, the quantum switch, which is the most prominent example of an indefinite causal order process, has been shown to enhance the precision of some metrological tasks compared to fixed order strategies [1].

In this work, we aim to experimentally demonstrate the advantages of indefinite causal order for multiparameter estimation in a photonic quantum switch. Our setup, based on Ref. [2], uses multicore optical fibers technology to coherently control the order of three quantum operations, two of them being noisy channels with variable noise strength.

Our setup is expected to estimate parameters even in noisy regimes where the consecutive application of the operations in a fixed order would make it unattainable. Additionally, we assess the Fisher information matrix for different configurations of the setup and different amounts of noise, showing that the best configuration of the quantum switch depends on a priori information and weighing of the parameters. Our results will highlight the pertinence of indefinite causal order for quantum information under noisy conditions.

2. Preliminaries

- **Three qubit operations:**

$$\hat{U}(\theta) = \exp(-i\theta\hat{\sigma}_x) \quad , \quad \theta \in [0, \pi/2] \quad ; \tag{1}$$

$$\mathcal{E}^{(A)}(\rho) = p_A\rho + (1-p_A)\frac{I}{2} \quad , \quad p_A \in [0,1] \quad ; \tag{2}$$

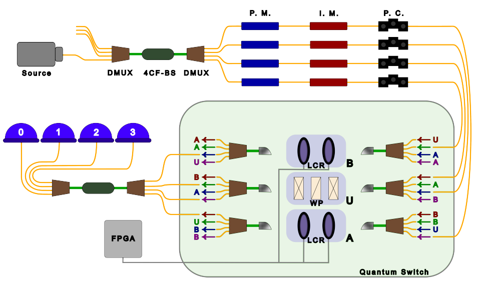
$$\mathcal{E}^{(B)}(\rho) = p_B\rho + (1-p_B)\frac{I}{2} \quad , \quad p_B \in [0,1] \quad . \tag{3}$$

- **Six possible orders:**

(0): $\mathcal{E}^{(A)} \circ \mathcal{E}^{(B)} \circ U$; (1): $U \circ \mathcal{E}^{(A)} \circ \mathcal{E}^{(B)}$; (2): $\mathcal{E}^{(B)} \circ U \circ \mathcal{E}^{(A)}$;
(3): $U \circ \mathcal{E}^{(B)} \circ \mathcal{E}^{(A)}$; (4): $\mathcal{E}^{(B)} \circ \mathcal{E}^{(A)} \circ U$; (5): $\mathcal{E}^{(A)} \circ U \circ \mathcal{E}^{(B)}$.

- **Task:**
To estimate (θ, p_A, p_B) when the operations are applied consecutively.

3. Experimental Setup



The setup uses the polarization degree of freedom of a photon as target and the path degree of freedom as control. The photon is coupled to one core of a four-core optical fiber and a multicore-fiber beam splitter prepares a state of superposition of four paths. Single-mode fibers route the photon through the operations in a coherent superposition of different orders.

Configurations: *(orders combined in the experiment)*
(0,1,3,4) (in the figure) and **(2,3,4,5)**

4. Expected results

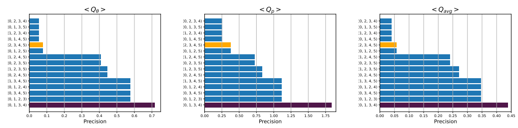
Figure of merit: According to the Cramér-Rao bound, $cov(\vec{x}) \geq F^{-1}/n$, where F is the Fisher information matrix and n is the size of the ensemble. We define the following measure for the achievable precision of estimating the i-th parameter in $\vec{x}$:

$$Q_i \equiv \begin{cases} \frac{1}{(F^{-1})_{ii}} & , \quad \text{if } F \text{ is invertible} \\ 0 & , \quad \text{otherwise} \end{cases} \quad . \tag{4}$$

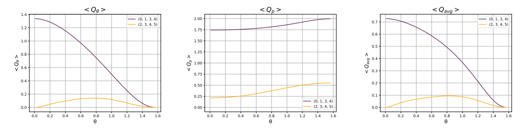
Similarly, for the average precision of estimating N parameters, we define

$$Q_{avg} \equiv \frac{N}{\mathrm{Tr}(F^{-1})} \quad , \tag{5}$$

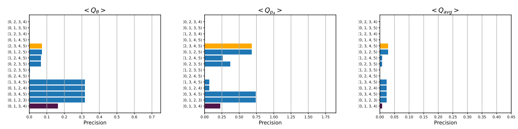
SCENARIO I: ($p_A = p_B = p$, with no prior information)



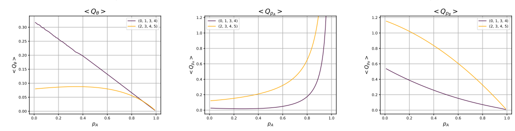
SCENARIO II: ($p_A = p_B = p$, with prior information on θ)



SCENARIO III: (independent parameters, with no prior information)



SCENARIO IV: (independent parameters, with prior information on p_A)



5. Conclusions

- We aim to demonstrate in a photonic setup that indefinite causal order can be leveraged to solve multiparameter estimation tasks.
- The performance of our experiment is independent of the initial polarization of the photons, as in other photonic realizations of the quantum switch.
- The quantum switch allows the simultaneous estimation of multiple parameters even in scenarios where a definite order would make it unattainable, e.g. in presence of fully depolarising channels.
- The quantum switch is a tunable device whose performance depends on its configuration. The optimal configuration for a given task should be chosen considering prior information or relevance of the involved parameters.

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Boundary-Controlled Phase Transitions in Non-Euclidean Tensor-Network Models of Social Influence

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RECOVERY AND RESILIENCE PLAN · Funded by the European Union NextGenerationEU

Abstract

We study equilibrium models of social influence on non-Euclidean lattices using tensor-network methods. Each agent carries discrete "features" (opinions, cultural traits), yielding a multi-component Potts/clock spin model. We apply the corner transfer matrix renormalization group (CTMRG) and the higher-order tensor renormalization group (HOTRG) to compute thermodynamic quantities on Euclidean square lattices, regular hyperbolic tilings, tree-like limits, and a fractal lattice. On the square lattice, the two-feature binary-trait model exhibits two continuous transitions; increasing traits induces a first-order transition. On hyperbolic lattices, boundary bias continuously shifts the transition temperature, reflecting the dominance of boundary degrees of freedom — analogous to boundary/bulk correspondence in holographic settings.

Motivation

- The Axelrod model [1] describes how cultural influence among interacting agents leads to either global consensus or fragmentation. We reformulate it as a **thermodynamic spin model** amenable to tensor-network methods, enabling treatment in the **thermodynamic limit**.
- **Non-Euclidean lattices**: real social networks are far from flat grids — their "small-world" properties are better captured by **hyperbolic** and **fractal** geometries.
- **Workshop connection**: on negatively curved lattices, the boundary dominates the bulk, mirroring the **holographic principle** (AdS/CFT boundary–bulk correspondence). Boundary agents can continuously control collective behavior — a concrete realization of how "observers at the boundary" shape the interior.

Model

Each agent (lattice site) carries f cultural features, each with $q^{(\alpha)}$ possible traits. We label models by the trait vector $\{q^{(1)}, \dots, q^{(f)}\}$ (e.g., $\{2,2\}$ for $f{=}2$ binary features; total $\prod_\alpha q^{(\alpha)}$ states per site).

Lattice notation (p,q): a regular tiling by p-gons with q meeting at each vertex. Examples: $(4,4)$ = Euclidean square lattice; $(5,4)$ and $(7,4)$ = hyperbolic.

Hamiltonian [2, 3]: $\mathcal{H} = \sum_{\langle ij\rangle} \mathcal{H}_{ij} + \sum_{i\in\partial} B_i$, where the two-site interaction and boundary field are

$$\mathcal{H}_{ij} = -J \sum_{\alpha=1}^{f} J_{ij}^{(\alpha)} \cos\left(\theta_i^{(\alpha)} - \theta_j^{(\alpha)}\right), \quad J_{ij}^{(\alpha)} = \prod_{\beta\neq\alpha} \delta\left(\sigma_i^{(\beta)}, \sigma_j^{(\beta)}\right), \quad B_i = -h \sum_{\beta=1}^{f} \cos\theta_i^{(\beta)} \tag{1}$$

Here $\theta_i^{(\alpha)} = 2\pi\sigma_i^{(\alpha)}/q^{(\alpha)}$, the sum $\langle ij\rangle$ runs over nearest neighbors, and ∂ denotes boundary sites. The conditional coupling $J_{ij}^{(\alpha)}$ activates the clock interaction on feature α only when all other features match (Potts condition). T plays the role of social noise; h is the boundary field (f_b).

Bulk free energy: $B = F - F_b$, where F is the total free energy per site and F_b the boundary contribution (removed via CTMRG). B is the key observable for locating phase transitions; normalized as B/q for comparison.

Lattice Geometries

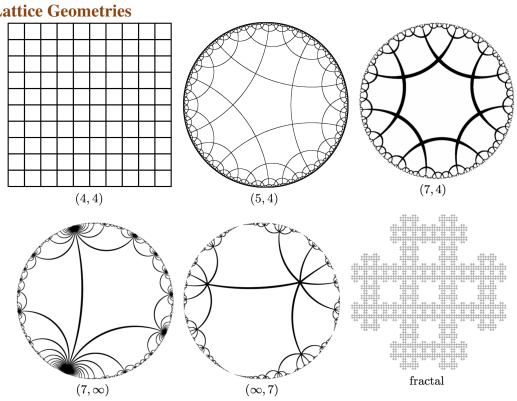

Figure 1: Top row: Euclidean $(4,4)$ square lattice, hyperbolic $(5,4)$ and $(7,4)$ lattices (Poincaré disk). **Bottom row:** tree-like limits $(7,\infty)$ and $(\infty,7)$, and the fractal lattice [7].

Methods

- **CTMRG** [4, 6]: works directly in the thermodynamic limit on regular lattices. Computes free energy, magnetization, and entanglement entropy. Boundary conditions controlled via boundary tensors: fixed (FBC, $f_b{=}1$) or open (OBC, $f_b{\to}0$).
- **HOTRG** [5]: used for fractal (non-translationally invariant) lattices; the local tensor is iteratively coarse-grained with bond dimension D.

Results: Euclidean $(4,4)$ Lattice

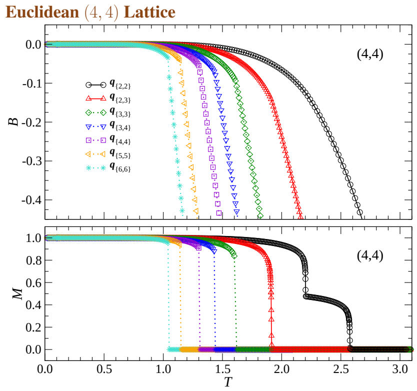

Figure 2: B/q (top) and M (bottom) vs. T on the $(4,4)$ square lattice for various $\{q^{(1)}, q^{(2)}\}$. $\{2,2\}$: **two continuous transitions** (solid lines); $q \geq 3$: single **first-order transition** (dotted/dashed).

Results: Hyperbolic $(5,4)$ Lattice

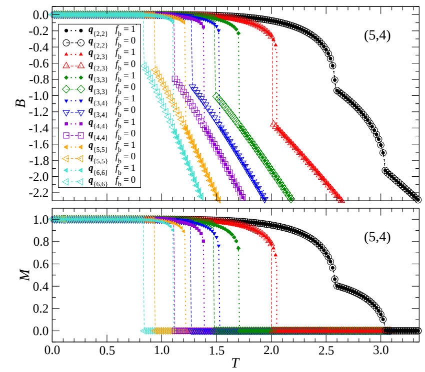

Figure 3: Same observables on the **hyperbolic** $(5,4)$ **lattice**. FBC (fixed boundary, $f_b{=}1$) and OBC (open boundary, $f_b{\to}0$) shown as separate curves. No free-energy crossing between branches (unlike Euclidean case).

Key Result: Boundary Control of Phase Transitions

On hyperbolic lattices, a **phase-coexistence region** $[T_{\rm low}, T_{\rm high}]$ exists [8], within which the phase-transition temperature $T_c(f_b)$ can be **continuously tuned** by varying the boundary field f_b. This does **not** happen on flat (Euclidean) lattices, where boundary effects are subdominant.

Physical interpretation: On negatively curved lattices, boundary sites grow exponentially — boundary "agents" control where within $[T_{\rm low}, T_{\rm high}]$ the transition occurs, analogous to how **boundary conditions control bulk physics** in holographic / AdS settings.

Geometry Dependence

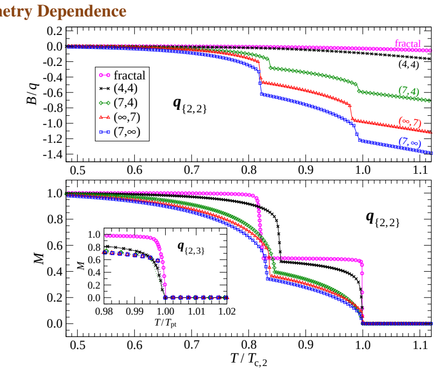

Figure 4: B/q and M vs. $T/T_{c,2}$ for the $\{2,2\}$ model across five geometries: fractal, $(4,4)$, $(7,4)$, $(\infty,7)$, $(7,\infty)$. Inset: $\{2,3\}$ magnetization. Phase-transition character varies with lattice curvature.

Conclusions & Outlook

1. On the Euclidean $(4,4)$ lattice: two continuous transitions for $\{2,2\}$; first-order for $q \geq 3$.
2. On the hyperbolic $(5,4)$ lattice: **boundary fields continuously tune** T_c within a phase-coexistence region $[T_{\rm low}, T_{\rm high}]$ [8] — a discrete analogue of holographic boundary/bulk correspondence.
3. Geometry shapes criticality: different curvatures yield distinct phase-transition profiles across Euclidean, hyperbolic, fractal, and tree-like lattices.
4. These tensor-network results serve as toy models for boundary/bulk correspondence, directly relevant to observers and causality in quantum gravity.

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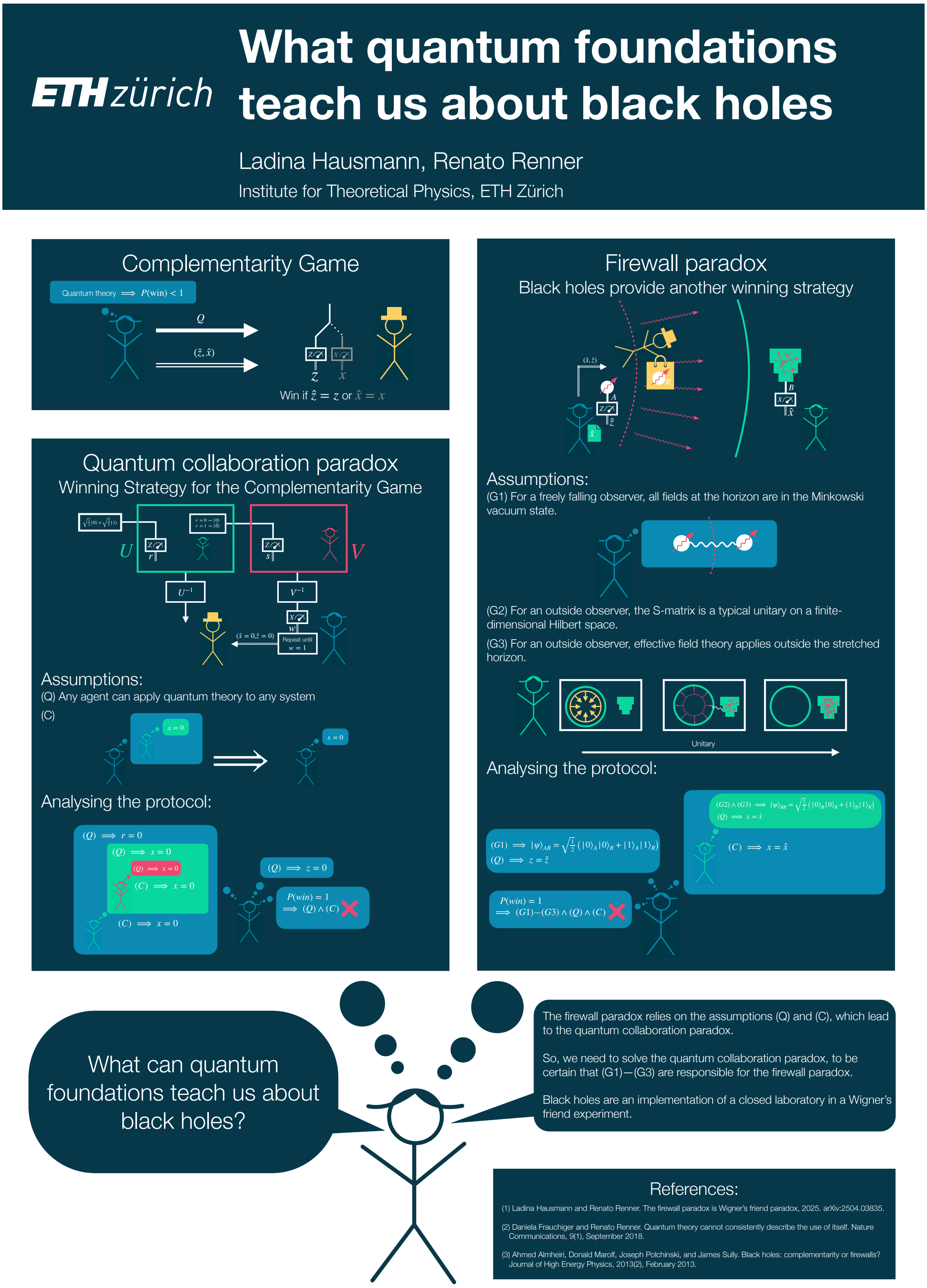
Ladina Hausmann

What Wigner's friend teaches us about black holes



Marek Liška

Semiclassical spacetime thermodynamics

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Introduction

It has been shown that the condition for vanishing of the variation of the total von Neumann entropy in a small sphere in curved spacetime is equivalent to the (traceless) semiclassical Einstein equations [1]. Our works [2] explicitly incorporates the backreaction of quantum fields in this paradigm. We focus on the 2D case where the effect of backreaction on von Neumann entropy is fully encoded in Wald entropy corresponding to the effective conformal anomaly action [3]. In order to address this case, we propose a new way to derive the full semiclassical dynamics of scalar-tensor gravity theories from equilibrium conditions. Our approach may be further extended to incorporate quantum reference frames, opening a way to study genuine quantum gravitational effects in the simple 2D setting.

Stretched light cones and their thermodynamic properties

- at any point P in a generic 2D spacetime we consider a local Minkowski coordinate system (t, x)
- we consider a (flat) future-directed light cone with the tip in P
- stretched light cone is a timelike surface Σ located just outside the light cone
- around P it has spatial radius α much smaller than the local curvature length scale, $\alpha \ll \sqrt{1/|R|}$
- corresponds to world lines of uniformly accelerating observers $u^{\mu} = \frac{1}{\alpha}(x, t)$ with acceleration $1/\alpha$
- in 2D, u^{μ} is an approximate Killing vector
- we invoke Einstein equivalence principle to construct a local Minkowski vacuum
- the observers see this vacuum as a thermal state at Unruh temperature $T_{\mathrm{U}} = \hbar/(2\pi\alpha)$
- valid at time scales $t \leq \epsilon \ll \alpha$

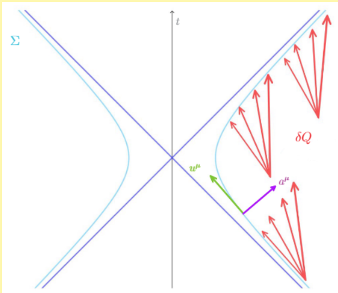

Figure 1: A sketch of a stretched light cone (image credit: Celia López Pineros)

- the stretched light cone has a well defined temperature
- the light cone itself is further a causal horizon with entanglement entropy
- due to their closeness, we can associate entanglement entropy with the stretched light cone
- the total generalised entropy of the stretched light cone consist of its entanglement entropy and of von Neumann entropy of the matter enclosed in it [4]
- the stretched light cone geometry only changes with t^2
- it is then nearly stationary at small time scales
- in 2D [1], entanglement entropy is then given by the dynamical Wald entropy prescription [5]
- in equilibrium change of entropy times Unruh temperature equals the work performed on the stretched light cone
- this relation is equivalent to the equations of motion of the semiclassical JT gravity

physical process equilibrium condition

$$\Delta S_0 + \Delta S_{\mathrm{P}} + \Delta S_{\mathrm{M}} = \Delta W / T_{\mathrm{U}}$$

Leading order entropy

- 2D gravity can be understood as effective dynamics of the horizon of a higher dimensional extremal black hole
- entanglement entropy of a test field in this extremal black hole background scales with area
- upon compactification, the area-scaled entropy becomes $S_0 = \eta\phi$, where ϕ is a scalar degree of freedom and η a constant
- we assume this behaviour of horizon entropy is universal and independent of its higher dimensional motivation
- in this way, we introduce the leading order contribution to 2D entanglement entropy without referencing gravitational dynamics

Conformal anomaly contribution

- conformal anomaly is a feature of dynamics of test conformal quantum field theory in a curved spacetime
- in 2D, it can be fully captured in the Polyakov action

$$I_{\mathrm{P}} = (-c/24\pi) \int \left(\chi R + \nabla_{\lambda}\chi\nabla^{\lambda}\chi\right) \sqrt{-\mathfrak{g}} \mathrm{d}^2 x$$

- c denotes the central charge of the conformal field theory ($c \gg 1$)
- the corresponding entanglement entropy reads

$$S_{\mathrm{P}} = (c/6)\,\chi + (c/24\pi) \int_{\Sigma} \nabla^{\mu}\chi\nabla^{\nu}\chi\alpha u_{\mu} \mathrm{d}\Sigma_{\nu}$$

- also obtained without referencing gravitational dynamics

evaluation between times $t = 0$ and $t = \epsilon$

$$0 = t^{\mu}t^{\nu}\Bigg[\left(\nabla_{\mu}\nabla_{\nu} - g_{\mu\nu}\frac{\Box}{2}\right)\left(\eta\phi - \frac{c}{6}\chi\right) + \frac{c}{6}\nabla_{\mu}\chi\nabla_{\nu}\chi - \frac{c}{12}g_{\mu\nu}\nabla_{\lambda}\chi\nabla^{\lambda}\chi + \frac{2\pi}{\hbar}T_{\mu\nu} + \eta\left(\frac{V(\phi)}{2} + W_0\right)g_{\mu\nu}\Bigg]$$

time direction t^{μ} is arbitrary, $G \equiv 1/(4\hbar\eta)$

$$0 = \left(\nabla_{\mu}\nabla_{\nu} - g_{\mu\nu}\frac{\Box}{2}\right)\left(\phi - \frac{2G\hbar c}{3}\chi\right) + \frac{2G\hbar c}{3}\nabla_{\mu}\chi\nabla_{\nu}\chi - \frac{G\hbar}{3}g_{\mu\nu}\nabla_{\lambda}\chi\nabla^{\lambda}\chi + 8\pi G T_{\mu\nu} + \left(\frac{V(\phi)}{2} + W_0\right)g_{\mu\nu}$$

divergenceless condition fixes W_0

$$\left(\nabla_{\mu}\nabla_{\nu} - g_{\mu\nu}\Box\right)\left(\phi - \frac{2G\hbar c}{3}\chi\right) + \frac{2G\hbar c}{3}\nabla_{\mu}\chi\nabla_{\nu}\chi + V(\phi)\,g_{\mu\nu} = 8\pi G T_{\mu\nu}$$

We recover semiclassical JT gravity with an arbitrary dilaton potential

Matter entropy

- matter crossing the boundary of the stretched light cone carries entropy
- at $t = 0$, von Neumann entropy of the conformal quantum fields inside the stretched light cone obeys

$$S_{\mathrm{M}} = (2\pi/\hbar) \int x T_{tt} dx$$

(requires the field state to be a small perturbation of Minkowski vacuum)

- for nearly stationary situations, it holds

$$\Delta S_{\mathrm{M}} = (2\pi/\hbar) \int_{\Sigma} t T_{tt} \mathrm{d}\Sigma$$

- compatible with the second law [5]
- reduces to the Clausius entropy flux in the classical limit [6]

Work

- the standard thermodynamic derivations effectively recover only the traceless equations of motion [7]
- cannot incorporate the effect of a potential of the scalar field $V(\phi)$
- $V(\phi)$ may be seen as an external condition imposed on the stretched light cone by the heat bath
- we treat the stretched light cone as an open system with $\nabla_{\nu}T_{\mu}{}^{\nu} = F_{\mu}$
- consequently, integral of $F_{\mu}u^{\mu}$ becomes a work term in the first law
- we consider work as consisting of two contributions, one proportional to $V(\phi)$, the other, W_0, unfixed

To do list

generalise to 4D

- same logic, sketched in our paper
- R^2 term in conformal anomaly cannot be localised
- conformal anomaly effective action does not seem to fully describe semiclassical backreaction
- combine with nonlocal (e.g. logarithmic) corrections to entropy?

rephrase in terms of von Neumann algebras

- allows more formal definition of generalised entropy
- can account for the effect of the clock dynamics
- fairly straightforward in 2D

apply the quantum reference frame formalism

- conceptually difficult
- basic idea: represent Unruh observers in terms of QRFs
- might allow QRF perspective on gravitational dynamics
- in principle can go beyond the semiclassical regime

Further reading

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Higher-order quantum processes respecting closed labs in a spacetime have quantum controlled causal order

Higher-order quantum processes respecting closed labs in a spacetime have quantum controlled causal order

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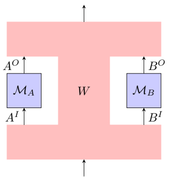
Process matrices provide a framework to model quantum information processing protocols in the absence of a well-defined acyclic causal order, defined without reference to a background spacetime.
Main question: Which indefinite causal order processes can be physically realised through experiments in a meaningful way in a classical acyclic spacetime?
Main contribution: any protocol realisable in an acyclic background spacetime that respects the closed labs assumption of the process matrix framework necessarily behaves as a QC-QC.

Frameworks

Process matrix framework [1]



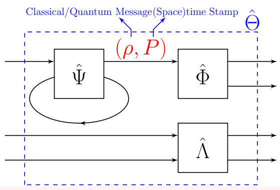
- **Local agents** apply CP maps in their labs
- Environment modelled by **process matrix** *W*
- **Causally non-separable process matrices** cannot be interpreted as mixtures fixed order processes and
- Can violate **causal inequalities**. Cf. Bell inequalities.

Quantum circuits with quantum control of causal order (QC-QC) [2]

- Subset of process matrices interpretable as **generalized quantum circuits**
- Clear physical interpretation
- Causal order can be controlled via a quantum system and/or dynamically by agent actions

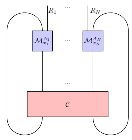
Causal box framework [3]

- Models protocols satisfying relativistic causality in fixed spacetime
- Quantum messages in superpositions of different spacetime locations
- **Composition**: closed under parallel, sequential composition and feedback loops
- **Causality**: Outputs only depend on inputs in the causal past



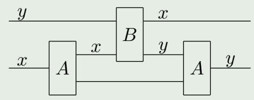
Results

Modelling multi-agent protocols in spacetime



A **causal box protocol** consists of
- a causal box $\mathcal{C}$ (= environment)
- sets of causal boxes $\{\mathcal{M}^{A_k}_{x_k}\}$ for each agent (modelling allowed operations of each agent)

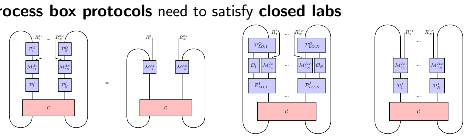
Trivial violation of causal inequalities



Maximal trivial causal inequality violation with a causal box protocol
Game: Alice to guess Bob's input bit y and Bob Alice's input bit x
Protocol: Alice at $t = 1$ sends x to Bob, Bob at $t = 2$ sends y to Alice
Lesson: Arbitrary causal box protocols are not processes

What makes a protocol in spacetime a process?

Process box protocols need to satisfy **closed labs**



Acting once: Agents can only act once
Local order: If an agent receives a message at spacetime point t, they output at $\mathcal{O}_k(t) \succ t$ where $\mathcal{O}_k$ is injective. In particular, agents must receive a message before they send one

Mapping process box protocols to QC-QCs

The behaviour of any process box protocol can be reproduced by a QC-QC. For every process box protocol $\mathfrak{P} = \{\mathcal{C}, \{\mathcal{M}^{A_k}_{x_k}\}_{k,x_k}\}$ there exists a QC-QC $\mathfrak{Q}$ together with operations $\{\mathcal{M}'^{A_k}_{x_k}\}_{k,x_k}$ such that $\forall \{x_k\}_k$

$$\mathcal{C} \circ \bigotimes_k \mathcal{M}^{A_k}_{x_k} = \mathfrak{Q} \circ \bigotimes_k \mathcal{M}'^{A_k}_{x_k}$$

Conclusions and outlook

- Only QC-QCs are physical in fixed, acyclic spacetime
- Recover QC-QCs from spacetime principle
- Genuine violations of causal inequalities are not possible in a fixed, acyclic spacetime
- Cyclic or quantum spacetimes? Some of our techniques may be useful as we disentangle information-theoretic and spacetime notions of causality
- Can we weaken our closed labs assumptions?

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Max Joseph Fahn

Relational observables and quantum master equations for bosonic matter coupled to linearised gravity

Relational Observables and Causality in Gravitational Decoherence Scenarios

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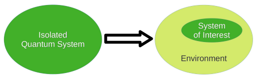
Open Quantum Systems and Decoherence

- Investigate effective dynamics of a matter system of interest $\hat{\rho}_S$ interacting with an environment of gravitational waves (GWs) based on GR and QFT/QM



- Total Hamiltonian: $\hat{H}_{tot} = \hat{H}_S \otimes \hat{\mathbb{1}}_{\mathcal{E}} + \hat{\mathbb{1}}_S \otimes \hat{H}_{\mathcal{E}} + \hat{H}_{int}$
- Master equation for effective evolution of matter system:

$$\frac{\partial}{\partial t}\hat{\rho}_S(t) = -i[\hat{H}_S + \hat{H}_{add}, \hat{\rho}_S(t)] + \mathcal{D}[\hat{\rho}_S(t)]$$

$\implies$ Goal: Derive red terms

- New processes compared to an isolated quantum system, encoded in red terms:
  - **Energy shifts/renormalisation** of the energies of $\hat{\rho}_S$
  - **Dissipation** (energy flux from system into environment)
  - **Decoherence** (information flux from system into environment) $\implies$ diagonalisation ("classicalisation") of $\hat{\rho}_S$ in a certain basis

Relational Observables and Quantum Master Equations

with Kristina Giesel and Michael Kobler (FAU Erlangen-Nürnberg)
In: Class. Quant. Grav. 40.9 (2023), p. 094002. arXiv:2206.06397
with Kristina Giesel and Roman Kemper (FAU Erlangen-Nürnberg)
arXiv:2602.07622

- **Aim:** Derive master equation from classical action of scalar field/electromagnetic field coupled to linearised GR using **relational formalism** [1, 2] and **Ashtekar variables** in Fock quantisation
- Degrees of freedom (canonically conjugate pairs):

linearised gravity: $(\delta\mathcal{E}^a_i, \delta\mathcal{A}^i_a)$ scalar field: (ϕ, π) photon field: (E^a, A_a)

- Application of the relational formalism:
  - Total system contains n (linearised) constraints $\{C_I\}$: Hamiltonian constraint, spatial diffeomorphism constraint, Gauß constraint(s)
  - Construct n reference fields $\{T^I\}$ from degrees of freedom ($\implies$ "geometrical clocks") s.t. $\{C_I, T^J\} = \delta^J_I$ and $\{T^I, T^J\} = 0$
  - Each reference field comes with a clock parameter τ^I; picking specific values for these corresponds to gauge fixing; here we leave them general
  - Apply the observable map (up to desired perturbative order) to all degrees of freedom F

$$\mathcal{O}_{F,\{(T^I-\tau^I)\}} = \exp\left(\kappa \int_{\mathbb{R}^3} d^3y \, \left(T^I(\vec{y}) - \tau^I(\vec{y})\right) \cdot \{C_I(\vec{y}), .\}\right) F$$

and the (vacuum) dual map

$$\mathcal{O}^{\text{dual}}_{F,\{C_I\}} = \exp\left(-\kappa \int_{\mathbb{R}^3} d^3y \, C_I(\vec{y}) \cdot \left\{\left(T^I(\vec{y}) - \tau^I(\vec{y})\right), .\right\}\right) F$$

$\implies$ physical degrees of freedom (relational observables): STT modes for linearised gravity + scalar field / transverse modes of the photon field
$\implies$ Fock quantisation only of the physical degrees of freedom

- Trace out environment (in thermal state with temperature Θ)
$\implies$ master equation for small coupling κ (similarities to [3–5]):

$$\frac{\partial}{\partial t}\hat{\rho}_S(t) = -\,i[\hat{H}_S + \kappa\,\hat{U} + \kappa\,\hat{H}_{LS}, \hat{\rho}_S(t)] + \frac{\kappa}{2} \sum\!\!\!\!\!\!\int_{a,b} \sum_{\epsilon} R_{ab}(t)\left(\hat{j}^b_\epsilon\,\hat{\rho}_S(t)\,(\hat{j}^a_\epsilon)^\dagger - \frac{1}{2}\left\{(\hat{j}^a_\epsilon)^\dagger\,\hat{j}^b_\epsilon, \hat{\rho}_S(t)\right\}\right)$$

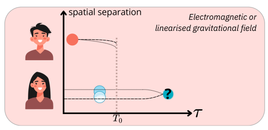
with self-interaction term $\hat{U}$, Lamb-Shift-like Hamiltonian $\hat{H}_{LS}$, time-dependent coefficients $R_{ab}(t)$ and "current" operators $\hat{j}^b_\epsilon$

- **Results:**
  - Implementation of relational formalism in open quantum systems in perturbation theory; master equation depends on choice of reference fields
  - Single-particle projection for the scalar field case [6]

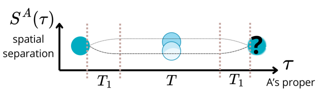
Communication across Quantised Black Hole Horizons?

with Alessandro Pesci (INFN Bologna)
In: Phys.Rev.D 112 (2025) 12, p. L121502 arXiv:2507.16911
In: Phys.Rev.D 112 (2025) 12, p. 124036 arXiv:2507.18709

- **Basic question:**
  - Alice (A) has a system in a spatial superposition. After a certain time interval, she recombines the branches and checks for coherence.
  - Bob (B), at some distance, can entangle a particle with A's superposition and can decide to measure at T_0 before A recombines.



  - B's measurement might collapse A's superposition $\implies$ A finds decoherence upon recombination $\implies$ Possible communication from B to A [7]
  - What changes if B is placed behind a horizon? [8]
- Setup: Alice stationary outside a black hole, opens and closes spatial superposition of an electrically charged/massive particle described by $\hat{\rho}_S(t)$



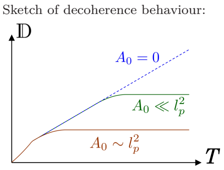
Here: Focus for simplicity on electromagnetic case;
Assumption: $T \gg T_1$

- Density matrix of A's particle when closing the superposition:
$\hat{\rho}_S(t) = \hat{\rho}_S(0) \cdot e^{-\frac{i}{\hbar}(H_i - H_j)t - \mathbb{D}}$ with decoherence factor (see [9]):

$$\mathbb{D} = \frac{q^2}{2}\int_{-\infty}^{\infty} d\tau\, d\tau'\; S^A(\tau)\; S^B(\tau')\langle\{\hat{E}_A(\tau, X), \hat{E}_B(\tau', X)\}\rangle_\Omega$$

with $\hat{E}_A(\tau, X)$ electric field quantised on a black hole spacetime (e.g. Schwarzschild) evaluated in A's frame, Ω Unruh/Hartle-Hawking vacuum

- Main result for horizons with classical horizon geometry (see e.g. [8–10]): Very low frequency radiation ("entangling photons") pierces the horizon and leads to decoherence $\mathbb{D} \propto T$
$\implies$ horizons with classical geometries destroy spatial superpositions
$\implies$ A's superposition decoheres, no matter what B does behind the horizon
$\implies$ no communication from B to A possible
- **Horizons with quantum geometries**: Assumption of a minimal quantum of area A_0 (e.g. of the order of the Planck length squared, see [11–13]) arising for instance from a minimal length in spacetime ([14, 15])
$\implies$ minimal frequency required to interact with the horizon
- Implementation in Minkowski spacetime with Alice following a Rindler trajectory as near horizon approximation to a Schwarzschild black hole
- Sketch of decoherence behaviour:



  - Modified behaviour: Linear increase in T saturates after a certain time
  - Saturation value $\propto \frac{1}{A_0}$
  - Negligibly small if $A_0 \sim l_p^2$ $\implies$ decoherence strongly suppressed

- Communication from B to A possible?
Horizon quantum geometry prevents entangling photons to cross horizon
$\implies$ B cannot entangle his particle with A's superposition
$\implies$ his measurement does not influence A's superposition
$\implies$ no communication from B to A possible
- **Results:**
  - Modified decoherence behaviour compared to classical horizon geometries
  - Horizons with quantum geometries decohere superpositions only negligibly if $A_0 \sim l_p^2$, as the low frequency radiation cannot interact with the horizon
  - No communication from B to A possible in this gedankenexperiment through classical and quantised horizons $\implies$ Causality preserved

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Mathematical Metrology and Information Processing

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Abstract

The poster presents a unification of classical and quantum mechanics in terms of mathematical metrology. In that regard, the uncertainty relation of position and momentum corresponds to one of classical time and velocity. An intermediary between theories concerns signal processing wherein the frequency operator $v = \hbar\, {}^{d}/_{dt}$ is induced by the Fourier transform

$$\hat{\psi}(v) = \int e^{-vt/\hbar}\,\psi(t)dt$$

considering $\hbar = {}^{h}/_{2\pi i}$ to be an imaginary constant. The derivation by quantum time is defined in order to satisfy a kinetical requirement that gives rise to the Schrödinger equation. Having inducted statistical superoperators which act onto densities, one gets to a measurement process that is presented by the equation of a free particle. The time superoperator should provide a change in representation from the reversible to an irreversible dynamics.

Uncertainty relation

Heisenberg algebra: $[t, v] = \hbar \Leftrightarrow [{}^{d}/_{dt}, {}^{d}/_{dv}] = {}^{1}/_{\hbar}$

Quantum mechanics: position t & momentum v

Signal processing: position (time) t & frequency v

Classical mechanics: time t & velocity v

Galilean transformation: $s{+}{+}: s \mapsto s+1$

$$t{+}{+}: t \mapsto t+1$$
$$v{+}{+}: s \mapsto s+t$$

CM	QM
time	position
velocity	momentum
1	mass

Heisenberg group: $t{+}{+}; v{+}{+}; t{-}{-}; v{-}{-} = s{+}{+}$

Group commutator: $e^{d/dt}\, e^{d/dv}\, e^{-d/dt}\, e^{-d/dv} = e^{1/\hbar}$

Classical distribution: $\psi = e^{s/\hbar} \Leftrightarrow |\psi|^2 = 1$

Quantum distribution: $|\psi|^2 = \psi\bar{\psi} = e^{s-\bar{s}/\hbar} = e^{2i\mathrm{Im}(s)/\hbar}$

Quantum states are related to a complexification of the distance.

Statistical mechanics

Quantum ensembles

Density operator: $\rho \geq 0,\ \mathrm{tr}\,\rho = 1$

Derivative by time: $\frac{d\rho\cdot}{d\tau} = \frac{\partial\rho}{\partial\tau}\cdot + \rho\frac{d\cdot}{d\tau}$

$$\Leftrightarrow \frac{\partial}{\partial\tau} = \left[\frac{d}{d\tau},\cdot\right]$$

Schrödinger equation: $\frac{\partial t}{\partial\tau} = v$

$$\Leftrightarrow \hbar\frac{d}{d\tau} = \underbrace{\frac{v^2}{2} + E(t)}_{h}$$

Evaluation derivative: $\underbrace{\hbar\frac{d\psi}{d\tau} - {}^{v^2\psi}/_2}_{\hbar\frac{\delta\psi}{\delta\tau}} = E\psi$

Evaluation operator: $\psi(\cdot)_\tau = e^{-\frac{v^2\tau}{2\hbar}}\,\psi$

$$= \underbrace{\frac{1}{\sqrt{2\pi\hbar\tau}}e^{\frac{t^2}{2\hbar\tau}}}_{\Phi_p} * \psi$$

Derivative redefinition: $\frac{\delta\psi}{\delta\tau} = \frac{\psi(\cdot)_\tau - \psi(\cdot)_{\tau_0}}{\tau-\tau_0}$

$$\psi = e^{s/\hbar} \Rightarrow \underbrace{\frac{ds}{d\tau} - \frac{d^2s}{2\hbar dt^2}}_{\frac{\delta s}{\delta\tau}} = E + {}^{(\frac{ds}{\hbar dt})^2}/_2$$

Wave potential: ${}^{-\left|\frac{d\psi}{dt}\right|^2}/_2 = {}^{(\frac{ds}{\hbar dt})^2}/_2$

Linear frequency: $k = \frac{v}{\mathrm{h}} = \frac{d}{2\pi i dt}$

Coordinates and derivatives

Derivation by operator:

$$\hbar\frac{\partial}{\partial t} = \left[\hbar\frac{d}{dt},\cdot\right] = [v,\cdot]$$
$$\hbar\frac{\partial}{\partial v} = \left[\hbar\frac{d}{dv},\cdot\right] = [-t,\cdot]$$

Comutativity: $\left[\frac{\partial}{\partial t}, \frac{\partial}{\partial v}\right]\cdot = \left[\left[\frac{d}{dt}, \frac{d}{dv}\right],\cdot\right] = 0$

$$[T, V] = 0$$

Canonical coordinate:

$$T = \{t,\cdot\} = {}^{t\cdot+\cdot t}/_2$$
$$V = \{v,\cdot\} = {}^{v\cdot+\cdot v}/_2$$

Derivation by superoperator:

$$\left[\frac{\partial}{\partial t}, T\right] = 1 \Rightarrow \frac{\partial}{\partial T} = \left[\frac{\partial}{\partial t},\cdot\right]$$
$$\left[\frac{\partial}{\partial v}, V\right] = 1 \Rightarrow \frac{\partial}{\partial V} = \left[\frac{\partial}{\partial v},\cdot\right]$$

Hamiltonian: $H = \{h,\cdot\} = {}^{h\cdot+\cdot h}/_2$

$$= \frac{V^2}{2} + E(T) + \mathrm{o}(\hbar)$$

Liouvillian: $L = [h,\cdot] = h\cdot - \cdot h$

$$= \hbar\underbrace{\left(\frac{\partial H}{\partial V}\frac{\partial}{\partial t} - \frac{\partial}{\partial v}\frac{\partial H}{\partial T}\right)}_{\langle H\rangle} + \mathrm{o}(\hbar^2)$$

Liouville equation: $\frac{\partial\rho}{\partial\tau} = \langle H\rangle\rho + \mathrm{o}(\hbar)$

Reversibility to irreversibility

Derivative by time: $\frac{\partial}{\partial\boldsymbol{\tau}} = \left[\frac{\partial}{\partial\tau},\cdot\right] = [{}^{L}/_{\hbar},\cdot]$

Hamilton equation: $\frac{\partial T}{\partial\boldsymbol{\tau}} = \frac{\partial H}{\partial V} \qquad \frac{\partial V}{\partial\boldsymbol{\tau}} = -\frac{\partial H}{\partial T}$

Free particle: $h = \hbar\frac{d}{d\tau} = \frac{v^2}{2}$

Hamiltonian: $H = \{h,\cdot\} = \frac{V^2}{2} + \frac{\hbar^2}{8}\frac{\partial^2}{\partial v^2}$

Liouvillian: $L = \hbar\frac{\partial}{\partial\tau} = \left[\hbar\frac{d}{d\tau},\cdot\right] = \hbar V\frac{\partial}{\partial t}$

Time superoperator: $\boldsymbol{\tau} = \frac{T}{V}$

Uncertainty relation: $[L, \boldsymbol{\tau}] = \left[\hbar V\frac{\partial}{\partial t}, \frac{T}{V}\right] = \hbar$

Change in representation: $\Lambda = \lambda(\boldsymbol{\tau})$

Liouville equation: $\hbar\frac{\partial\rho}{\partial\tau} = L\rho$

General solution: $\rho = e^{\tau L/\hbar}\rho_0$

Lie group: $U^\tau = e^{\tau L/\hbar}$

Markov semigroup: $W^\tau = \Lambda U^\tau \Lambda^{-1}$

Irreversible process: $\rho \geq 0 \not\Rightarrow W^{-1}\rho \geq 0$

Increase of the Rényi entropy: $-\log\|W^\tau\cdot\|^2$

Preservation of positivity: $\rho \geq 0 \Rightarrow \Lambda\rho \geq 0$

Preservation of equilibrium: $\Lambda 1 = 1$

Preservation of trace: $\mathrm{tr}\,\Lambda\rho = \mathrm{tr}\,\rho$

Preservation of information: Λ^{-1} on a dense domain

Measurement process

Euclidean algorithm: $\frac{a}{b} = \cfrac{1}{n_1 + \cfrac{1}{n_2 + \cfrac{1}{\ddots}}}$

Automorphism: $?: \cfrac{1}{n_1 + \cfrac{1}{n_2 + \cfrac{1}{\ddots}}} \mapsto \underbrace{\frac{1}{2^{n_1-1}} - \frac{1}{2^{n_1+n_2-1}} + \cdots}_{\underbrace{0.0\ldots0}_{n_1}\underbrace{1\ldots1}_{n_2}\underbrace{0\ldots}_{\ldots}}$

Frobenius-Perron group: $U^\vartheta = e^{\vartheta\log 2\cdot vt/\hbar}$

$$\frac{d\Phi_p}{d\vartheta\log 2} = {}^{vt}/_{\hbar}\,\Phi_p = -2\tau\frac{d\Phi_p}{d\tau} = \frac{d\Phi_p}{-\frac{1}{2}d\log\tau}$$

Group reparametrization: $\tau = 2^{-2\vartheta}$

$$\frac{d\psi(\cdot)_\tau}{d\tau} = \psi * \frac{d\Phi_p}{d\tau} = \psi * \frac{d^2\Phi_p}{2\hbar dt^2} = \frac{v^2}{2\hbar}\psi(\cdot)_\tau$$

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Prequantum Physics and the Spin-Geometry Theorem

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Abstract

Standard QM presupposes a well-defined classical "context" (e.g. reference frames, devices, repeatable preparations) out of which the parameters of the Born rule are written. This presupposition breaks down in « low-resource » regimes of physics. We revisit Penrose's spin-networks programme understanding as a finite-resource formulation of QM. Built with no background space, Hilbert space, or continuum of parameters assumed, this framework carries a natural rational probability rule, encoding a degradable notion of predictability. We introduce the ***δ-classicality*** *condition that characterises regimes of effectively unbounded resources, in which repeated operations yield stable probabilities and an operational separation between "system" and "classical context" emerges. In this regime we prove the* ***Spin-Geometry Theorem****: for δ-classical spin-networks, the probabilities generated by the ½-unit exchange define angles satisfying all constraints of 3D Euclidean geometry, and the probability reduces to the Born rule. Quantum mechanics thus emerges as the large-resource limit of a more primitive, discrete theory.*

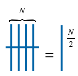
1. Classical contexts, probabilities and physicality

- An **operational notion of existence:** one can say that something exists *in science* if it can in principle be realised and probed.
  —> Everything exists relative to the « context » necessary to realise or probe it.
- No reference frame/measurement apparatus, no property: a physical quantity (e.g. expressed in terms of a physical unit) one cannot possibly measure, not even in principle, cannot be said to exist.
  —> **Physicality and measurability require a context**.
- Quantum states are containers of probabilities, so they can be legitimately used or defined, only insofar as the probabilities they summarise can be defined.
  The context necessary to possibly *realise* a probability is any situation where one can repeat the preparation-measurement a large amount of times in an i.i.d. way.
- But what if one *cannot* prepare-and-measure N times identically ? What if the preparation-measurement context in which we do QM degrades every time we use it, or worse what if it is not well-defined in the first place ?
  —> Probabilities (as usually understood) **cannot be legitimately used**, and neither can quantum states. **QM breaks** in such scenarios.
- In a low-resource scenario, one can no longer distinguish a context from an object, a measurement from a preparation, a time duration from a length from a mass.

2. Spin networks and the logic of relations

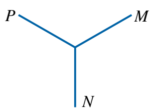
Two objects and a relation: |

Two (unordered) sets of N objects and N relations:



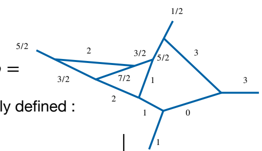
Given two sets of N and M objects, how many objects can a third set P hold ?

$$\boxed{\begin{array}{c} M+N+P \in 2\mathbb{N} \\ |M-N| \leq P \leq M+N \end{array}}$$



(recoupling rules for spin!)

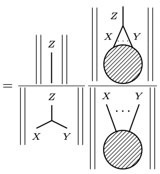
Build trivalent graphs half-integer labels: $\omega =$



A purely combinatorial norm can be naturally defined :

$$\|\omega\| := \prod_j \frac{1}{(2j)!} \prod_v A(v)^2 \left| \sum_c (-1)^{\#\text{ crossings in } c} (-2)^{l_c} \right|$$

Out of which a *rational* probability rule based on *counting* can be defined:

$P\,(XY \to Z) =$ 

—> Crucially, **no physical content nor Hilbert space** has been introduced here.

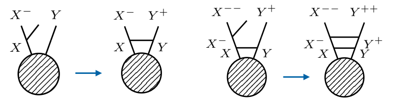
3. Stable probabilities

- Consider an « experiment » inside a SN with n outcomes. This experiment is said to be δ**-classical** if:

$$\left| \mathrm{Prob}(\sigma_n^{(2)} \,|\, \sigma_m^{(1)}) - \mathrm{Prob}(\sigma_n^{(1)}) \right| < \delta \qquad \forall m, n, \qquad \text{with } \delta \ll 1$$

—> Rebuilding the experiment in *the same* SN yields the same probabilities as the first time.

Example:



- One can repeat identically the « experiment », building the context for the probability rule to be realisable.

4. The Spin-Geometry theorem

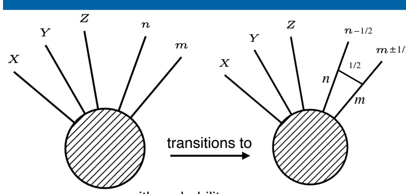
Theorem (Spin-Geometry): *Let ω be a spin-network with $N \geq 3$ δ-classical open ends with respect to the 1/2-unit exchange operation.*
Then, for each pair (X, Y) of these open ends, the probability associated with the 1/2-unit exchange from X to Y takes the form:

$$\mathrm{Prob}_{XY}(X^{\pm}Y^{\pm}) = \frac{1}{2}(1 + \cos\theta_{XY}) + O(\varepsilon),$$

where the θ_{XY} satisfy all the constraints of angles between directions in three-dimensional Euclidean space, and $\varepsilon \to 0$ as $\delta \to 0$.
(Proof in our upcoming paper)

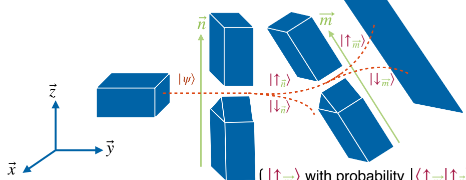
- Probabilities are now repeatable + « Bulk » of the SN decouples from the experiment, only parameters built by it stay in the probabilities.
  —> Production of the context that justifies the use of probabilities
  —> The same context provides the parameters the quantum states are expressed in.

5. Example



with probability :
$P(nm \to (n-1/2)(m \pm 1/2))$

Spin-Geometry theorem



spin-1/2 particle is found in $\begin{cases} |\uparrow_{\vec{m}}\rangle \text{ with probability } |\langle\uparrow_{\vec{m}}|\uparrow_{\vec{n}}\rangle|^2 \\ |\downarrow_{\vec{m}}\rangle \text{ with probability } |\langle\downarrow_{\vec{m}}|\uparrow_{\vec{n}}\rangle|^2 \end{cases}$

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Subsystems (in)dependence in GIE proposals

Subsystems (in)dependence in GIE proposals

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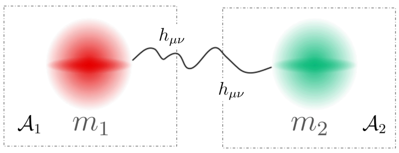
I. Recent proposals suggest that detecting entanglement between two spatially superposed masses would establish the quantum nature of gravity. However, these gravitationally induced entanglement (GIE) experiments rely on assumptions about subsystem independence [1,2].



II. We sharpen the theoretical underpinnings of such proposals by examining them through the lens of algebraic quantum field theory, distinguishing distinct operational and algebraic notions of independence [3,4].

Assumption 1 (Operational: Independent preparation).

$$\omega|_{\mathcal{A}_1} = \omega_1, \qquad \omega|_{\mathcal{A}_2} = \omega_2 .$$

Assumption 2 (Operational: Uncorrelated initial preparation).

$$\omega(A_1A_2) = \omega_1(A_1)\,\omega_2(A_2), \qquad \forall\, A_i \in \mathcal{A}_i$$

Assumption 3 (Operational: Independent measurability).

$$[\mathcal{A}_1, \mathcal{A}_2] = 0$$

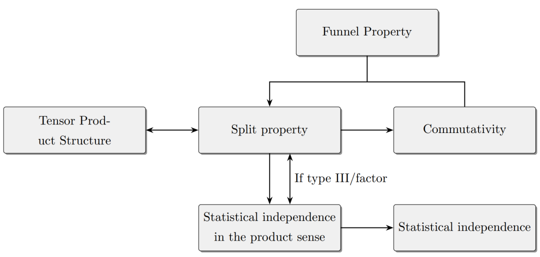
Assumption 4 (Algebraic: Finite-dimensionality and completeness).

$$\mathcal{H} \simeq \mathcal{H}_1 \otimes \mathcal{H}_2$$

III. We argue that state and measurement independence of subsystems, essential to the experimental logic, is nontrivial in the presence of gauge constraints and gravitational dressing. Using gravitationally dressed fields, we recall that commutation relations between spacelike separated observables are nontrivial, undermining strict Hilbert space factorization [5,6,7].

$$[\dot{\Phi}(x), \Phi(x')] = -\frac{i\kappa^2\hbar}{32\pi c^4 |x - x'|}\dot{\phi}(x)\dot{\phi}(x')$$

$$[\Phi(x), \Phi(x')] = -\frac{i\kappa^2\hbar}{64\pi c^4}\left[\dot{\phi}(x)\partial_i\phi(x') + \partial_i\phi(x)\dot{\phi}(x')\right]\frac{x^i - x'^i}{|x - x'|}$$



IV. We further explore the implications for entanglement witnesses, investigating the Tsirelson bound when subsystem algebras fail to commute, and showing that the Tsirelson bound persists for a suitably symmetrized CHSH observable even though the operational status of such "joint" observables becomes delicate when commensurability fails.

Proposition. Let A_1, A_2, B_1, B_2 be self-adjoint observables satisfying $A_i^2 = B_j^2 = I$. Without assuming commensurability (i.e. without requiring $[A_i, B_j] = 0$), the symmetrized CHSH operator

$$E_\circ \;=\; A_1 \circ B_1 + A_1 \circ B_2 + A_2 \circ B_1 - A_2 \circ B_2 \;=\; \frac{E + E'}{2}$$

satisfies the bound

$$E_\circ \;\leq\; 2\sqrt{2}\, I\,,$$

where $E = A_1B_1 + A_1B_2 + A_2B_1 - A_2B_2$ and $E' = B_1A_1 + B_2A_1 + B_1A_2 - B_2A_2$.

V. Finally, we derive estimates for dressing-induced microcausality violations, which suggest a complementary avenue to current proposals: in principle, bounding dressing-induced microcausality violations themselves as a probe of the quantum nature of gravity.

$$\phi(x) = \frac{\lambda}{\sqrt{2m}}\, e^{-imc^2t/\hbar}\, \psi(x) \quad \dot{\phi}(x) = -\frac{imc^2}{\hbar}\, \phi(x)$$

$$\frac{\left|[\Phi_C(x), \Phi_C(x')]\right|}{|\phi(x)\phi(x')|} \sim \frac{Gm}{c^2L} \;\approx\; 10^{-36}$$

$$\frac{\left|[\dot{\Phi}_C(x), \Phi_C(x')]\right|}{|\phi(x)\phi(x')|} \sim \frac{Gm}{c^2L}\,\frac{mc^2}{\hbar} \;\approx\; 10^{-36}\,\frac{mc^2}{\hbar} \;\approx\; 10\,\mathrm{s}^{-1}$$

$$\frac{\left|[\dot{\Phi}_C(x), \Phi_C(x')]\right|}{|\phi(x)\phi(x')|} \sim \frac{Gm}{c^2L}\,\omega_{\mathrm{kin}} \sim 4 \times 10^{-44}\,\mathrm{s}^{-1}$$

$$\omega_{\mathrm{kin}} \sim \hbar/(2mL^2) \simeq 5 \times 10^{-9}\,\mathrm{s}^{-1}$$

VI. The standard quantum-information narrative behind GIE-style inference is not conceptually automatic once one insists on gauge-invariant observables for gravity. The relevant observables must be gravitationally dressed, and the resulting gauge-invariant subalgebras need not commute even at spacelike separation. This obstructs implementing exact commensurability and blocks the usual route from microcausality to split inclusions, tensor-product factorization, and statistical independence.

However, the present results still support and qualify the viewpoint on GIE: they support it by showing that, for realistic parameters, the induced violations of microcausality are tiny enough that the quantum-information modelling can remain an excellent approximation; and they qualify it by stressing that strong ontological conclusions rely on formal properties that are not preserved once gauge-invariant gravitational observables are taken seriously.

The fact that factorization itself is undermined suggests that gravity may affect not only how we test quantumness, but also how we should think about subsystems and measurement in regimes where diffeomorphism constraints are inescapable.

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Nikolaos Mitrakos

When does entanglement through gravity imply gravitons?

Does entanglement through gravity imply gravitons?

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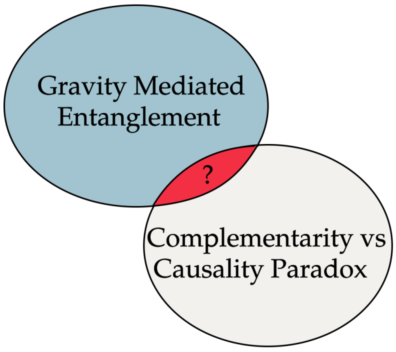
The claim under study

In a thought experiment involving entanglement through gravity, a paradox can arise if quantum fluctuations are ignored. **Causality** and **complementarity** are in conflict [1], which has led to:

The received view

1. **Newtonian entanglement** and consistency with **causality** imply the existence of **gravitons**.
2. There is a regime where Newtonian entanglement can be seen as entanglement through gravitons.



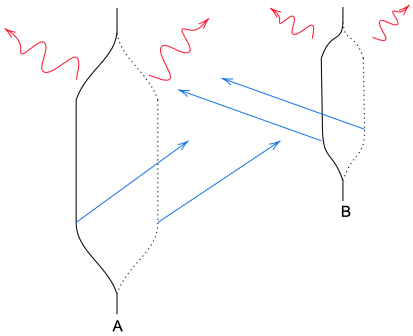


Figure 1. Scalar model of mediated entanglement. The colored arrows correspond to the colored terms in $\hat{\rho}_{AB}$.

The paradox revisited

The paradox involves a one way signalling setup, Fig. 2. The two sides of the paradox correspond to two ways of ignoring quantum fluctuations:

- Set $\Gamma_A = \Gamma_B = 0$: **Complementarity** is violated, B acquires which-path information while A remains coherent.
- Stationary phase approximation: **Causality** is violated, A loses coherence due to B, even though B is in spacelike separation.

For both approximations the which-path information propagates at the speed of light.

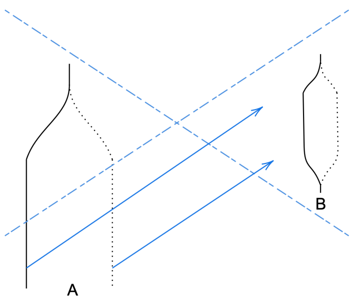


Figure 2. Setup of the paradox.

Scalar model of mediated entanglement

Model linearized gravitational interaction with a scalar field mediator:

- A and B are the two masses,
- each follows a **superposition of trajectories** $\sim |\mathrm{L}\rangle + |\mathrm{R}\rangle$,
- each mass couples only to the scalar field linearly.

After splitting and recombination, the final state of the two masses

$$\hat{\rho}_{AB} = \frac{1}{4}\begin{pmatrix} 1 & e^{-\Gamma_B}\, e^{i\varphi_L^A} & e^{-\Gamma_A}\, e^{i\varphi_L^B} & e^{-(\Gamma_A+\Gamma_B+\Gamma_{AB})}\, e^{\frac{i}{2}\sum_s(\varphi_s^B+\varphi_s^A)} \\ & 1 & e^{-(\Gamma_A+\Gamma_B-\Gamma_{AB})}\, e^{\frac{i}{2}\sum_s(\varphi_s^B-\varphi_s^A)} & e^{-\Gamma_A}\, e^{i\varphi_R^B} \\ & & 1 & e^{-\Gamma_B}\, e^{i\varphi_R^A} \\ & & & 1 \end{pmatrix}$$

Entangling phases:

- from **retarded potential** of each mass at the position of the other,
- reproduce GME in 'Newtonian' limit [2].

Decoherence terms:

- from **quantum fluctuations** of the mediating field,
- local radiation production $\Gamma_{A/B}$ or correlations harvested from field vacuum Γ_{AB}.

Related through **uncertainty principle** [3, 4]:

$$\Gamma_A\Gamma_B \geq \left((\varphi_R^A - \varphi_R^B) - (\varphi_L^A - \varphi_L^B)\right)^2$$

Our conclusion

1. Entanglement through retarded potentials and consistency with causality imply the existence of gravitons.
2. Only retarded potentials lead to entanglement in this setup, gravitons lead to decoherence.

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Scattering in a Holographic Quantum Field Theory

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Scattering in a Holographic QFT (hQFT)

Paul Schneidewind Telge, Oliver Friedrich

Holographic Scaling of Degrees of Freedom

According to the holographic principle, the number of physical degrees of freedom within a spatial region scales (at most) with its boundary area rather than its volume, as implied by the entropy bound derived from black hole thermodynamics[1,2]:

$$S(R) \leq \frac{\partial R}{4}.$$

This seems to contradict typical QFT, where the amount of field operators scales with volume. In Fourier space this means that the amount of Fourier modes and ladder operators inside a momentum shell scale like[3]:

number of degrees of freedom

$$N_s \propto k_s^2 \Delta_s.$$

shell width

For holographic scaling of degrees of freedom this should be[3,4]:

a "holographic" number of degrees of freedom

$$n_s \propto k_s \Delta_s.$$

How do we get N_s degrees of freedom from n_s degrees of freedom if $N_s \geq n_s$!?

Overlapping Quantum Field Theory

The holographic principle suggests that the N_s many *apparent* degrees of freedom are actually a function of only n_s many *true* degrees of freedom. We make the assumption that this is a linear function, given by a $N_s \times n_s$ matrix, F:

"apparent" ladder operators

$$\hat{A}_p = \sum_k F_{pk} \hat{a}_k$$

"true" ladder operators

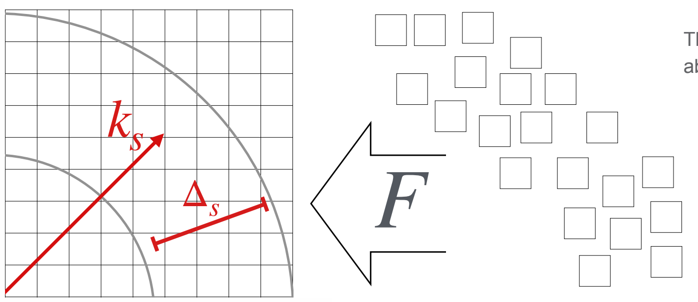
We take N_s many combinations of n_s many operators!

We split the hamiltonian in multiple shells and make the replacement $\hat{a}_p \rightarrow \hat{A}_p$ in each of them. We explore the resulting theory.

gives the correct expectation value to $A_k^\dagger |0\rangle$

$$\hat{H} = \sum_s \hat{H}_s \qquad \hat{H}_s = \frac{k_s}{\lambda_s} \sum_k \hat{A}_k^\dagger \hat{A}_k$$

We do a similar thing to anti-particles and their hamiltonian



We also make said replacement in the mode expansion of quantum fields:

Weyl spinor

$$\psi(x,t) = \frac{1}{\sqrt{V}} \sum_k \frac{u(k)}{\sqrt{k}} A_k(t) e^{-ikx} + \dots$$

Anti-particle part

Scattering in hQFT

As the underlying theory is still a QFT on a lattice (with varying couplings) it is possible to construct a path integral and explore what modifications we expect from said overlaps to the Feynman rules in scattering amplitudes. These include:

- Modified propagator, which conserves energy but not necessarily **3**-momentum:

matrix handling the spinor components

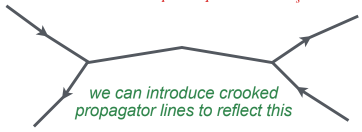
matrix that diagonalizes $FF^\dagger$

$$\int \frac{d\omega}{2\pi V} \sum_{pqi} (FR^T)_{pi} \frac{-\Sigma(q,p,\tilde{k}_i(\omega)) e^{i\omega(x^0-y^0)-iqy+ipx}}{(\tilde{k}_i(\omega))^2 + i\epsilon} (RF^T)_{iq}.$$

Comes in with momentum p …

4-momentum with ω energy and "momentum" that gets mapped to the shell of p and q it is also $\approx k_s$

…comes out with momentum q



we can introduce crooked propagator lines to reflect this

- Although momentum seems to be conserved at the vertex, because it isn't conserved in the propagator, we need to sum over it every time two propagators meet.
- External legs of diagrams come with energy $\omega = k_s$ instead of the one determined by their momentum.

Connections to Renomalization?

We can also understand the F matrices as fine graining, taking the lattice from n points to N points. Such effects can indeed be absorbed into the vertices if the matrix $F^T F$ is approximately diagonal:

required to conserve momentum at the vertex

$$\sum_q F_{pq}^T g F_{qp} \rightarrow g'$$

The effect of the additional degrees of freedom is absorbed by the coupling constant!

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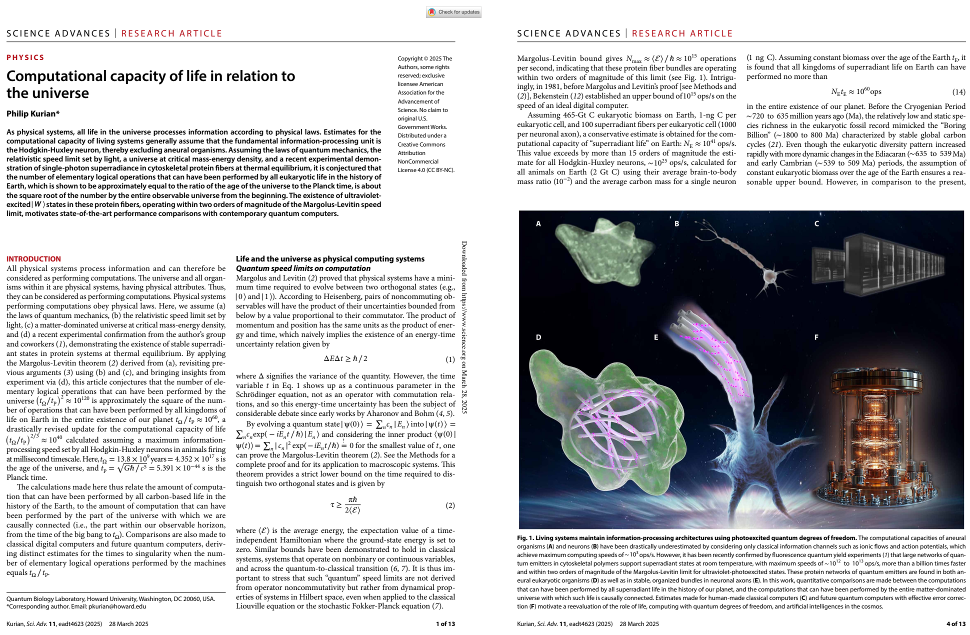
Philip Kurian

Superradiant life: The Heisenberg cut within the observer?



PHYSICS

Computational capacity of life in relation to the universe

Philip Kurian*

As physical systems, all life in the universe processes information according to physical laws. Estimates for the computational capacity of living systems generally assume that the fundamental information-processing unit is the Hodgkin-Huxley neuron, thereby excluding aneural organisms. Assuming the laws of quantum mechanics, the relativistic speed limit set by light, a universe at critical mass-energy density, and a recent experimental demonstration of single-photon superradiance in cytoskeletal protein fibers at thermal equilibrium, it is conjectured that the number of elementary logical operations that can have been performed by all eukaryotic life in the history of Earth, which is shown to be approximately equal to the ratio of the age of the universe to the Planck time, is about the square root of the number by the entire observable universe from the beginning. The existence of ultraviolet-excited $|W\rangle$ states in these protein fibers, operating within two orders of magnitude of the Margolus-Levitin speed limit, motivates state-of-the-art performance comparisons with contemporary quantum computers.

INTRODUCTION

All physical systems process information and can therefore be considered as performing computations. The universe and all organisms within it are physical systems, having physical attributes. Thus, they can be considered as performing computations. Physical systems performing computations obey physical laws. Here, we assume (a) the laws of quantum mechanics, (b) the relativistic speed limit set by light, (c) a matter-dominated universe at critical mass-energy density, and (d) a recent experimental confirmation from the author's group and coworkers (*1*), demonstrating the existence of stable superradiant states in protein systems at thermal equilibrium. By applying the Margolus-Levitin theorem (*2*) derived from (a), revisiting previous arguments (*3*) using (b) and (c), and bringing insights from experiment via (d), this article conjectures that the number of elementary logical operations that can have been performed by the universe $(t_\Omega/t_P)^2 \approx 10^{120}$ is approximately the square of the number of operations that can have been performed by all kingdoms of life on Earth in the entire existence of our planet $t_\Omega/t_P \approx 10^{60}$, a drastically revised update for the computational capacity of life $(t_\Omega/t_P)^{2/3} \approx 10^{40}$ calculated assuming a maximum information-processing speed set by all Hodgkin-Huxley neurons in animals firing at millisecond timescale. Here, $t_\Omega = 13.8 \times 10^9$ years $= 4.352 \times 10^{17}$ s is the age of the universe, and $t_P = \sqrt{G\hbar/c^5} = 5.391 \times 10^{-44}$ s is the Planck time.

The calculations made here thus relate the amount of computation that can have been performed by all carbon-based life in the history of the Earth, to the amount of computation that can have been performed by the part of the universe with which we are causally connected (i.e., the part within our observable horizon, from the time of the big bang to t_Ω). Comparisons are also made to classical digital computers and future quantum computers, deriving distinct estimates for the times to singularity when the number of elementary logical operations performed by the machines equals t_Ω/t_P.

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Life and the universe as physical computing systems

Quantum speed limits on computation

Margolus and Levitin (*2*) proved that physical systems have a minimum time required to evolve between two orthogonal states (e.g., $|0\rangle$ and $|1\rangle$). According to Heisenberg, pairs of noncommuting observables will have the product of their uncertainties bounded from below by a value proportional to their commutator. The product of momentum and position has the same units as the product of energy and time, which naively implies the existence of an energy-time uncertainty relation given by

$$\Delta E \Delta t \geq \hbar / 2 \tag{1}$$

where Δ signifies the variance of the quantity. However, the time variable t in Eq. 1 shows up as a continuous parameter in the Schrödinger equation, not as an operator with commutation relations, and so this energy-time uncertainty has been the subject of considerable debate since early works by Aharonov and Bohm (*4*, *5*).

By evolving a quantum state $|\psi(0)\rangle = \sum_n c_n |E_n\rangle$ into $|\psi(t)\rangle = \sum_n c_n \exp(-iE_n t/\hbar)|E_n\rangle$ and considering the inner product $\langle\psi(0)|\psi(t)\rangle = \sum_n |c_n|^2 \exp(-iE_n t/\hbar) \stackrel{!}{=} 0$ for the smallest value of t, one can prove the Margolus-Levitin theorem (*2*). See the Methods for a complete proof and for its application to macroscopic systems. This theorem provides a strict lower bound on the time required to distinguish two orthogonal states and is given by

$$\tau \geq \frac{\pi\hbar}{2\langle\mathcal{E}\rangle} \tag{2}$$

where $\langle\mathcal{E}\rangle$ is the average energy, the expectation value of a time-independent Hamiltonian where the ground-state energy is set to zero. Similar bounds have been demonstrated to hold in classical systems, systems that operate on nonbinary or continuous variables, and across the quantum-to-classical transition (*6*, *7*). It is thus important to stress that such "quantum" speed limits are not derived from operator noncommutativity but rather from dynamical properties of systems in Hilbert space, even when applied to the classical Liouville equation or the stochastic Fokker-Planck equation (*7*).

Margolus-Levitin bound gives $N_{max} \approx \langle\mathcal{E}\rangle/\hbar \approx 10^{15}$ operations per second, indicating that these protein fiber bundles are operating within two orders of magnitude of this limit (see Fig. 1). Intriguingly, in 1981, before Margolus and Levitin's proof [see Methods and (*2*)], Bekenstein (*12*) established an upper bound of 10^{15} ops/s on the speed of an ideal digital computer.

Assuming 465-Gt C eukaryotic biomass on Earth, 1-ng C per eukaryotic cell, and 100 superradiant fibers per eukaryotic cell (1000 per neuronal axon), a conservative estimate is obtained for the computational capacity of "superradiant life" on Earth: $N_E \approx 10^{41}$ ops/s. This value exceeds by more than 15 orders of magnitude the estimate for all Hodgkin-Huxley neurons, $\sim 10^{25}$ ops/s, calculated for all animals on Earth (2 Gt C) using their average brain-to-body mass ratio (10^{-2}) and the average carbon mass for a single neuron (1 ng C). Assuming constant biomass over the age of the Earth t_E, it is found that all kingdoms of superradiant life on Earth can have performed no more than

$$N_E t_E \approx 10^{60} \text{ops} \tag{14}$$

in the entire existence of our planet. Before the Cryogenian Period ~720 to 635 million years ago (Ma), the relatively low and static species richness in the eukaryotic fossil record mimicked the "Boring Billion" (~1800 to 800 Ma) characterized by stable global carbon cycles (*21*). Even though the eukaryotic diversity pattern increased rapidly with more dynamic changes in the Ediacaran (~635 to 539 Ma) and early Cambrian (~539 to 509 Ma) periods, the assumption of constant eukaryotic biomass over the age of the Earth ensures a reasonable upper bound. However, in comparison to the present,

Fig. 1. Living systems maintain information-processing architectures using photoexcited quantum degrees of freedom. The computational capacities of aneural organisms (**A**) and neurons (**B**) have been drastically underestimated by considering only classical information channels such as ionic flows and action potentials, which achieve maximum computing speeds of $\sim 10^3$ ops/s. However, it has been recently confirmed by fluorescence quantum yield experiments (*1*) that large networks of quantum emitters in cytoskeletal polymers support superradiant states at room temperature, with maximum speeds of $\sim 10^{12}$ to 10^{13} ops/s, more than a billion times faster and within two orders of magnitude of the Margolus-Levitin limit for ultraviolet-photoexcited states. These protein networks of quantum emitters are found in both aneural eukaryotic organisms (**D**) as well as in stable, organized bundles in neuronal axons (**E**). In this work, quantitative comparisons are made between the computations that can have been performed by all superradiant life in the history of our planet, and the computations that can have been performed by the entire matter-dominated universe with which such life is causally connected. Estimates made for human-made classical computers (**C**) and future quantum computers with effective error correction (**F**) motivate a reevaluation of the role of life, computing with quantum degrees of freedom, and artificial intelligences in the cosmos.

Philip LeMaitre

Particle Detectors Witness the Non-Classicality of the Vacuum with Contextuality

Harvesting Contextuality from the Vacuum

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Basic Idea: Vacuum correlations can generate contextuality in Unruh-DeWitt (UDW) systems using suitable measurements for a single qutrit and qubit-qutrit setup, including gapless systems which cannot harvest entanglement. The qubit-qutrit system also exhibits tunable contextuality-entanglement trade-off relations, allowing for non-classicality to remain in the system if one resource disappears.

Motivating question: are there other resources besides entanglement in quantum fields?

Yes, contextuality!

Key question: can an initially non-contextual UDW system state become contextual after a local interaction with the vacuum?

Contextuality Framework

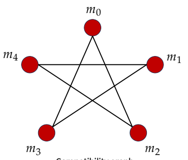
Contextuality asks whether a family of compatible local measurement data admit a single global model independent of context.

Example:

$X = \{m_0, m_1, m_2, m_3, m_4\}, \qquad O = \{-1, 1\},$

$\mathscr{C} = \{\{m_0, m_2\}, \{m_0, m_3\}, \{m_1, m_3\}, \{m_1, m_4\}, \{m_2, m_4\}\}$

	$-1,-1$	$-1,1$	$1,-1$	$1,1$
(m_0, m_2)	0	1/9	2/3	2/9
(m_0, m_3)	0	2/3	1/3	0
(m_1, m_3)	0	1/3	1/3	1/3
(m_1, m_4)	0	1/3	2/3	0
(m_2, m_4)	0	2/3	1/9	2/9



Compatibility graph

Contextuality generalizes familiar nonclassical resources such as magic [5] and nonlocal entanglement, and is equivalent to Wigner negativity as a measure of non-classicality

Above probability table is called an empirical model [1] and can be written as

$$e = \underbrace{CF(e)e^{SC}}_{\text{Contextual part}} + \underbrace{NCF(e)e^{NC}}_{\text{Non-contextual part}},$$

Measure contextuality with the contextual fraction $CF(e)$

Can phrase as linear program:

$$\boxed{\begin{array}{rl} \text{Find} & \boldsymbol{b} \in \mathbb{R}^n \\ \text{maximizing} & \mathbf{1} \cdot \boldsymbol{b} \\ \text{subject to} & \boldsymbol{M}\boldsymbol{b} \leq \boldsymbol{v}^e \\ \text{and} & \boldsymbol{b} \geq \mathbf{0}. \end{array}}$$

e as a vector: $\boldsymbol{v}^e[(C,s)] = e_C(s).$

Global assignment matrix: $\boldsymbol{M}[(C,s),g] = \begin{cases} 1 & \text{if } g|_C = s \\ 0 & \text{else.} \end{cases}$

Single Qutrit Contextuality Harvesting

$$\hat{\rho}_D(t) = \begin{pmatrix} 0 & 0 & -2\bar{W}_{11}^{++} \\ 0 & 2W_{11}^{+-} & 0 \\ -2(\bar{W}^*)_{11}^{++} & 0 & 1-2W_{11}^{+-} \end{pmatrix} + \mathcal{O}(\lambda^4)$$

$$\bar{W}_{dd'}^{pq} = \lambda_d \lambda_{d'} W_{\Delta t}(\Lambda_d^p, \Lambda_{d'}^q) \quad W_{dd'}^{pq} = \lambda_d \lambda_{d'} W(\Lambda_d^p, \Lambda_{d'}^q)$$

A single UDW system can harvest non-classicality as magic [3] or contextuality

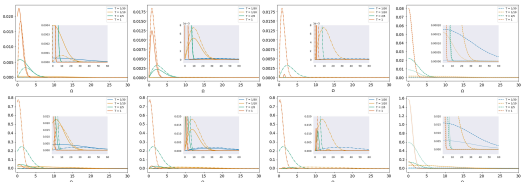


Figure 1: single-qutrit contextuality compared with harvested magic and with the genuine-harvesting diagnostic based on the Hadamard and signalling contributions.

- Sharp temporal profiles preserve genuinely quantum windows
- Longer interactions amplify signalling and can wash out harvested peak
- Smaller spatial smearing strengthens harvested contextuality
- Can tune measurements to alter harvested amount at zero gap, and contextuality can persist after harvested magic has already vanished

Can harvest contextuality and magic for gapless systems depending on the measurement scenario and smearing profile

Key Takeaways

1. **Single** qutrits harvest contextuality from vacuum correlations.
2. **Gapless** systems can harvest contextuality for suitable measurement choices.
3. Qubit–qutrit systems can harvest **both** negativity and local contextuality, with coexistence and trade-offs.
4. New genuine-harvesting criteria separate vacuum-resource extraction from field-mediated signaling.

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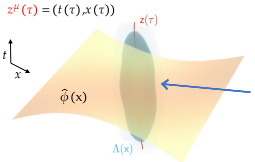
UDW Harvesting Protocol

Qudit following trajectory through spacetime

$z^\mu(\tau) = (t(\tau), x(\tau))$



Internal dynamics prescribed by free Hamiltonian

$$\hat{H}_0 = \sum_d \hat{H}_{0_d} = \sum_d \Omega_d(\hat{J}_z^{(d)} + j_d \mathbb{1}_d)$$

Interacts with massless scalar quantum field $\hat{\phi}(x)$ via Hamiltonian density

$$\hat{h}_d(x) = \lambda_d \Lambda_d(x) \hat{\mu}_d(\tau) \otimes \hat{\phi}(x)$$

Interaction smeared out over spacetime region via smearing function $\Lambda_d(x)$

Harvested contextuality measured using

$$\boxed{\Delta CF(\boldsymbol{v}^e) = CF(\boldsymbol{v}_t^e) - CF(\boldsymbol{v}_{t_0}^e).}$$

The protocol:

1. **Prepare** Non-contextual system ground state
2. **Couple** Local interaction with $\hat{\phi}(x)$
3. **Reduce** Trace out the field after time evolution
4. **Quantify** Evaluate $\Delta CF(v^e)$

Genuine Contextuality Harvesting

Can we call it "harvesting?"

- Vacuum state is highly non-classical [4] → Existing contextuality
- Resource theory of contextuality identifies free operations [6] → Use on field and UDW systems to reduce explanation of acquired contextuality to non-classicality of vacuum

Not all harvested contextuality is due to field

- Hadamard function encodes field state-dependence

Quantify: $\left|\dfrac{\Delta(\Lambda_d^+, \Lambda_{d'}^+)}{H(\Lambda_d^+, \Lambda_{d'}^+)}\right| \ll 1 \quad \text{and} \quad \Delta CF(\boldsymbol{v}^e) > 0.$

Genuine harvesting is identified when the Hadamard contribution dominates the signalling term in the relevant peak region

Qubit-Qutrit Contextuality Harvesting

Adding a qubit allows the joint system to harvest negativity while the qutrit still carries local contextuality. The two resource windows need not coincide, so they can compete, align, or separate as the detector profiles and separation are varied.

$$\hat{\rho}_D(t) = \begin{pmatrix} 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & r_{62}^* \\ 0 & 0 & W_{11}^{+-} & 0 & \sqrt{2}W_{21}^{+-} & 0 \\ 0 & 0 & 0 & 0 & 0 & r_{46} \\ 0 & 0 & \sqrt{2}W_{12}^{+-} & 0 & 2W_{22}^{+-} & 0 \\ 0 & r_{62} & 0 & r_{46}^* & 0 & r_{66} \end{pmatrix} + \mathcal{O}(\lambda^4),$$

$$r_{66} = 1 - W_{11}^{+-} - 2W_{22}^{+-},\ r_{62} = -\sqrt{2}(G^*)_{21}^{++},\ r_{46} = -2\bar{W}_{22}^{++},\quad G_{dd'}^{pq} = \lambda_d \lambda_{d'} G_F(\Lambda_d^p, \Lambda_{d'}^q)$$

Negativity: $$\mathcal{N}(\hat{\rho}_D(t)) \approx \left| \min\left(0, -\frac{1}{2}(W_{21}^{+-}W_{12}^{+-} + |r_{46}|^2)\right) + \min(0, -|r_{62}| + 2W_{22}^{+-}) \right| + \mathcal{O}(\lambda^4)$$

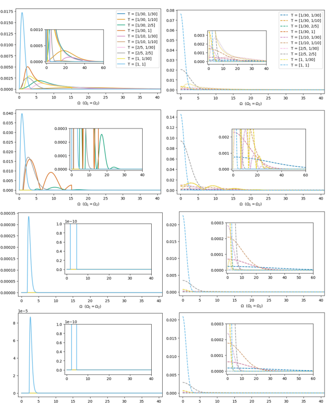


Figure 2: negativity in the qubit-qutrit system as a function of the shared gap, temporal widths, spatial width, and separation.

Harvesting Behaviour

Energy Gap Ω Both resources fade for large gaps

Spatial Width σ Smaller smearing generally amplifies harvesting

Temporal Width T Controls peak alignment and signalling contamination

Separation L Nonlocal entanglement dies sooner than local contextuality

Coexistence Window Fix the qutrit in a low-gap, high-contextuality regime and widen the qubit temporal profile to pull the negativity peak toward the same region. Both can coexist inside a controllable overlap region.

Non-classicality remains in the system even at large separation

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Pop Nicolina

Some properties of coherent states representation with applications in the Quantum Information

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COHERENT STATES FOR SOME QUANTUM OSCILLATORS REPRESENTATION OF QUANTUM INFORMATION THEORY

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ABSTRACT

In the quantum information theory are operating with qubits and N-qubits that can be express through of coherent states.
Density operator admits a representation in terms of coherent states formalism. Consequently, in this paper the notions of a qubit and density operators are cast in the framework of coherent states.
We have expressed a qubit as a coherent state, and thus a sequence of qubits is a tensor product of coherent states. For the ensembles of qubits, it could be used the density operator , in order to describe the informational content of the ensemble.
The coherent states representation may play an important role in the qantum information theory and the use of this formalism is not only theoretical, but also of some practical relevance by its applications.

INTRODUCTION

In quantum information theory (QIT) the fundamental unit of information is a *qubit* (i.e. quantum bit) or, generally, a *multiqubit* (or *N-qubit*) which is a vector in a two - dimensional (or2^N dimensional) Hilbert space. .

From the point of view of mathematics, a qubit is a linear combination of two 2X1 basis vectors from the two-dimensional Hilbert space [1], associated to a considered physical system:

$$|\Psi>=\widetilde{N}(\mu|0>+\nu|1>)=\widetilde{N}\begin{pmatrix}\mu\\ \nu\end{pmatrix} \qquad (1)$$

$$\widetilde{N}=\left(|\mu|^2+|\nu|^2\right)^{-1/2} \qquad |0>=\begin{pmatrix}1\\0\end{pmatrix}, \qquad |1>=\begin{pmatrix}0\\1\end{pmatrix}$$

A qubit containts a infinity of the states because of different values of the complex coefficients μ, ν. Often it is convenient to take:

$$|\mu|^2+|\nu|^2=1_{T_g}$$

In the next, we propose to indicate that it is possible another represantation of a qubit, namely, through the coherent states.
The normalized coherent states (CSs) of the one-dimensional harmonic oscillator (HO-1D) can be expanded on a infinite Fock - vector basis, where $|n>$ is the - boson Fock state :

$$|z>=e^{-\frac{1}{2}|z|^2}\sum_{n=0}^{\infty}\frac{z^n}{\sqrt{n!}}|n> \qquad (2)$$

where *z* is a complex variable of the coherent states.
The CSs must fulfill the identity operator decomposition (the identity resolution), as well

$$1=\int\frac{d^2z}{\pi}|z><z| \qquad (3)$$

where $d^2z=\frac{d\varphi}{2}d(|z|^2)$ is elementary area in complex z - plane.

$$|\Psi>=\int\frac{d^2z}{\pi}|z>\Psi(z^*)=\int\frac{d^2z}{\pi}|z>e^{-\frac{1}{2}|z|^2}(\mu+\nu z^*) \qquad (4)$$

A message $|M_n>$ is a tensor products of qubits [3]:

$$|M_n>=|\Psi_1>\otimes|\Psi_2>\otimes\ldots\otimes|\Psi_n>=\bigotimes_{i=1}^{n}|\Psi_i> \qquad (5)$$

$$|M_n>=\prod_{i=1}^{n}\left[\int\frac{d^2z_i}{\pi}|z_i>e^{-\frac{1}{2}|z_i|^2}(\mu_i+\nu_i z_i^*)\right] \qquad (6)$$

DENSITY OPERATOR

The mixed quantum states can not be described by the wave function or, equivalently, by an ket vector (as the pure states) so, in this case, we must use the density operator (or in some their concrete representation - the density matrix). The most widely used representation is is the energy representation whose basis vectors are the eigenvectors of the Hamiltonian: $|E_n>\equiv|n>$.

$$\rho=\sum_n w_n|n><n| \qquad (7)$$

$w_n=<n|\rho|n>$ - the probability to find the pure state $|n><n|$ in the irrespective mixture.

If we consider that the examined system is in the thermal contact with a reservoir to the equilibrium temperature $T=(\beta k_B)^{-1}$

The normalized density operator is:

$$\rho=\frac{1}{\overline{n}+1}\sum_n\left(\frac{\overline{n}}{\overline{n}+1}\right)^n|n><n| \qquad \overline{n}=\frac{1}{e^{\beta\hbar\omega}-1} \qquad (8)$$

The density operator can be also expanded on the CSs projectors $|z><z|$ as follows:

$$\rho=\int\frac{d^2z}{\pi}|z>P(|z|^2)<z| \qquad (8)$$

where $P(|z|^2)$ is so-called *P*- (quasi) distribution function (or Glauber - Sudarshan - function). This function is expressed as

$$P(|z|^2)=\frac{1}{\overline{n}}e^{-\frac{1}{\overline{n}}|z|^2} \qquad (9)$$

i.e. it is a Gaussian function of mean $|z|=0$ and variance $\overline{n}$

In this way, the corresponding density operator is the isotropic - Gaussian coherent state mixture,

$$\rho=\frac{1}{\overline{n}}\int\frac{d^2z}{\pi}e^{-\frac{1}{\overline{n}}|z|^2}|z><z| \qquad (10)$$

TRANSMISSIONS OF COHERENT STATE QUBITS

Let us we consider the following general expression of a qubit and their conjugate, as a linear combination of two coherent states, in the matrix representation:

$$|\Psi>=\widetilde{N}(\alpha,\beta)(\mu|\alpha>+\nu|\beta>)=\widetilde{N}(\alpha,\beta)\begin{pmatrix}\mu|\alpha>\\ \nu|\beta>\end{pmatrix} \qquad (11)$$

$$<\Psi|=\widetilde{N}(\alpha,\beta)(\mu^*<\alpha|+\nu^*<\beta|)=\widetilde{N}(\alpha,\beta)(\mu^*<\alpha| \quad \nu^*<\beta|) \qquad (12)$$

where α, β are two complex variables labelling the corresponding coherent states. The normalization constant is:

$$\widetilde{N}(\alpha,\beta)=\left[1+e^{-\frac{1}{2}\left(|\alpha|^2+|\beta|^2\right)}\left(\mu\nu^*e^{\alpha\beta^*}+\mu^*\nu e^{\alpha^*\beta}\right)\right]^{-\frac{1}{2}} \qquad (13)$$

Our task is to calculate the fidelity of the transmission of this qubit through a long optical fiber, in order to achieve the transmission of quantum information. The optical fibers can be characterized, from the point of view of qubit transmission, as having an exponential energy loss $\exp(-\lambda L)$ where λ is the loss coefficient of the fiber, and *L* is the transmission distance. As mentioned, high commercial fibers typically have $\lambda\approx 0.06/\mathrm{km}$ [4].

We will to realize the transmission of the pure qubit state

$$\rho_{in}=|\Psi><\Psi| \qquad (14)$$

through a thermal-noise channel and we will find the final state of the pure qubit statea after transmission.During the transmission process, information was lost to the environment and the pure qubit state will be into a mixture described by a density operator

$$\rho_{out}=\int\frac{d^2z}{\pi}q(|z|^2)D(z)\rho_{in}D^+(z) \qquad (15)$$

$D(z)=\exp(z a^+-z^*a)$ -the displacement operator

$$D^+(z)=D(-z) \quad ; \quad D(z)|0>=|z> \quad ; \qquad (16)$$

$$D(z)D(\beta)=e^{\frac{1}{2}(z\beta^*-z^*\beta)}D(z+\beta)$$

Their action on the qubits $|\Psi>$ and $<\Psi|$ are

$$D(z)|\Psi>=\widetilde{N}(\alpha,\beta)\begin{pmatrix}\mu\, e^{\frac{1}{2}(z\alpha^*-z^*\alpha)}|z+\alpha>\\ \nu\, e^{\frac{1}{2}(z\beta^*-z^*\beta)}|z+\beta>\end{pmatrix} \qquad (17)$$

$$<\Psi|D^+(z)=$$

$$=\widetilde{N}(\alpha,\beta)\left(\mu^*e^{\frac{1}{2}(-z\alpha^*+z^*\alpha)}<z+\alpha| \quad \nu^*e^{\frac{1}{2}(-z\beta^*+z^*\beta)}<z+\beta|\right)$$

$$D(z)|\Psi><\Psi|D^+(z)= \qquad (18)$$

$$=\left[\widetilde{N}(\alpha,\beta)\right]^2\begin{pmatrix}|\mu|^2|z+\alpha><z+\alpha| & \mu\nu^*|z+\alpha><z+\beta|\\ \mu^*\nu|z+\beta><z+\alpha| & |\nu|^2|z+\beta><z+\beta|\end{pmatrix}$$

The noise function $q(|z|^2)$ is of the Gaussian kind

$$q(|z|^2)=\frac{1}{\widetilde{n}_s}e^{-\frac{1}{\widetilde{n}_s}|z|^2} \qquad (19)$$

where the noise variance is: $\widetilde{n}_s=2\left[1-T\left(1-e^{-2s}\right)\right]$ (20)

$0\le T\le 1$ is the channel transmission coefficient of the noisy channel .

We observe that the transmission channel randomly displaces an input pure coherent state according to the Gaussian distribution which results in a thermal state.

Finally, the output density operator (density matrix) becomes

$$\rho_{out}=\left[\widetilde{N}(\alpha,\beta)\right]^4\frac{1}{\widetilde{n}_s}\int\frac{d^2z}{\pi}e^{-\frac{1}{\widetilde{n}_s}|z|^2}\begin{pmatrix}|\mu|^2|z+\alpha><z+\alpha| & \mu\nu^*|z+\alpha><z+\beta|\\ \mu^*\nu|z+\beta><z+\alpha| & |\nu|^2|z+\beta><z+\beta|\end{pmatrix}$$

For the pure qubit input state $\rho_{in}=|\Psi><\Psi|$ the channel fidelity F_{quant} can be derived as

$$F_{quant}=<\Psi|\rho_{out}|\Psi> \qquad (21)$$

$$F_{quant}=\left[\widetilde{N}(\alpha,\beta)\right]^2\frac{1}{\widetilde{n}_s}\int\frac{d^2z}{\pi}e^{-\frac{1}{\widetilde{n}_s}|z|^2}|<\Psi|D(z)|\Psi>|^2 \qquad (22)$$

In order to solve the integral we we will use the following general complex integral due to Perelomov [7], rather, an extension of this integral [8]:

$$\int\frac{d^2z}{\pi}e^{-\lambda|z|^2+\sigma z^*}f(z)=\frac{1}{\lambda}f\left(\frac{\sigma}{\lambda}\right) \qquad (23)$$

Finally, the quantum fidelity becomes:

$$F_{quant}=\frac{1}{\widetilde{n}_s+1}\frac{1-2|\mu|^2|\nu|^2\left(1-e^{-\frac{1}{4}\frac{\widetilde{n}_s}{\widetilde{n}_s+1}|\alpha-\beta|^2}\right)}{\left[1+e^{-\frac{1}{2}\left(|\alpha|^2+|\beta|^2\right)}\left(\mu^*\nu\, e^{\alpha^*\beta}+\mu\nu^*\, e^{\alpha\beta^*}\right)\right]^2} \qquad (24)$$

PARTICULAR CASES

a) The input state is a coherent state of the HO-1D

$$|\Psi>^{(z';0)}=|z'> \qquad (25)$$

This means that *μ=1, α=z', ν=0, β=0* and the fidelity is

$$F_{quant}^{(z';0)}=\frac{1}{\widetilde{n}_s+1} \qquad (26)$$

b) The input pure qubit state is a linear combination of Schrödinger's „cat states"

$$|\Psi>^{(-z';z')}=\frac{1}{\sqrt{1+e^{-2|z'|^2}(\mu^*\nu+\mu\nu^*)}}\begin{pmatrix}\mu|-z'>\\ \nu|z'>\end{pmatrix} \qquad (27)$$

$$F_{quant}^{(-z';z')}=\frac{1}{\widetilde{n}_s+1}\frac{1-2|\mu|^2|\nu|^2\left(1-e^{-\frac{\widetilde{n}_s}{\widetilde{n}_s+1}|z'|^2}\right)}{\left[1+e^{-2|z'|^2}(\mu^*\nu+\mu\nu^*)\right]^2} \qquad (28)$$

c) The input pure qubit state is of the following encoding method: $|\alpha>=|0>$, $|\beta>=|2z'>$

$$|\Psi>^{(0;2z')}=\frac{1}{\sqrt{1+e^{-2|z'|^2}(\mu^*\nu+\mu\nu^*)}}\begin{pmatrix}\mu|0>\\ \nu|2z'>\end{pmatrix} \qquad (29)$$

$$F_{quant}^{(0;2z')}=F_{quant}^{(-z';z')} \qquad (30)$$

CONCLUSIONS

- We have examined the relationship between the qubit and the coherent states of the one - dimensional quantum oscillator and, if we refer to the mixed quantum states - the density operator or density matrix.
- W e have calculated the standard fidelity of an transmission channel (optical fiber), in a general manner, if the input state is a pure qubit state . The fidelity is an important characteristic of quantum channels and a nontrivial quantity to study, because it evaluates how well the channel preserves the transmitted information.

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An inequality for relativistic local quantum measurements

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ÖSTERREICHISCHE AKADEMIE DER WISSENSCHAFTEN

Abstract

We investigate the trade-off between vacuum insensitivity and sensitivity to excitations in finite-size detectors, taking measurement locality as a fundamental constraint. We derive an upper bound on the detectability of vacuum excitation, given a small but nonzero probability of false positives in the vacuum state. The result is independent of the specific details of the measurement or the underlying physical mechanisms of the detector and relies only on the assumption of locality. To connect with realistic photodetection scenarios, we model the detection region as a square prism operating over a finite time window and consider a normally incident single-mode coherent state. Numerical results show that suppressing dark counts necessarily tightens the achievable click probability.

Setup: detector measurement

The detector is modeled by a POVM of the form $\{\widehat{E}_{\text{click}}, \mathbb{I} - \widehat{E}_{\text{click}}\}$, where a positive outcome of $\widehat{E}_{\text{click}}$ corresponds to a click. We denote

$$P_{\text{click}} := \langle\psi|\widehat{E}_{\text{click}}|\psi\rangle, \qquad P_{\text{dark}} := \langle\Omega|\widehat{E}_{\text{click}}|\Omega\rangle$$

as the probability that the detector registers a click in

$$|\psi\rangle := \text{arbitrary state}, \qquad |\Omega\rangle := \text{vacuum}.$$

Ideal detectors

In line with the conventional understanding of particle detectors, an idealized detector is expected to satisfy:

(i) **Vacuum insensitivity**: The detector should not respond when the system is in the vacuum state: the dark count probability should be zero:

$$P_{\text{dark}} = 0. \tag{1}$$

(ii) **Sensitivity to some state**: There exists at least one state $|\psi\rangle$ for which the detector registers a click:

$$P_{\text{click}} = 1. \tag{2}$$

Local detectors

We assume the detector operates within a fixed spacetime region

$$\mathcal{O}_{\text{det}} := \text{detector region}.$$

In accordance with the principles of AQFT, the measurement operator $\widehat{E}_{\text{click}}$ must be an element of the local algebra associated with this region:

$$\widehat{E}_{\text{click}} \in \mathfrak{A}(\mathcal{O}_{\text{det}}) := \text{local algebra in } \mathcal{O}_{\text{det}}. \tag{3}$$

No-go theorem: Equations (1), (2) and (3) are incompatible.
↓
Trade-off: The more the detector is ideal with respect to one of the two properties, (i) and (ii), the less it is with respect to the other. A detector that responds less to the vacuum tends to be less responsive to all states.
How to quantify this trade-off?

Inequality [1]

Reeh–Schlieder theorem: any state $|\psi\rangle$ can be approximated arbitrarily well by applying a local operator to the vacuum $|\Omega\rangle$. Hence,

$$\exists \widehat{A}_\zeta \in \mathfrak{A}(\mathcal{O}'_{\text{det}}) : \lim_{\zeta\to 0} \mathcal{E}_\zeta = 0, \text{ with } \mathcal{E}_\zeta := \left\| |\psi\rangle - \widehat{A}_\zeta|\Omega\rangle \right\|$$

By applying this theorem, together with the triangle inequality and the fact that $\widehat{E}_{\text{click}}$ is a POVM element, we obtain the inequality

$$\boxed{P_{\text{click}} \leq \min_{\zeta>0} \left(\mathcal{E}_\zeta + \left\|\widehat{A}_\zeta\right\| \sqrt{P_{\text{dark}}}\right)^2 .} \tag{4}$$

Example: Coherent states [2]

$$\text{Coherent state: } |\psi\rangle = \widehat{W}(f)|\Omega\rangle = \exp[i\widehat{\phi}(f)]|\Omega\rangle,$$

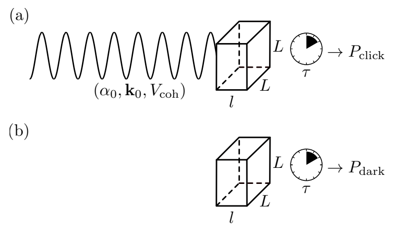
$$\widehat{A}_\zeta(f) = \int_{\mathbb{R}} d\eta\, G_\zeta(\eta - i\pi)\, \widehat{W}[f \circ J \circ \Lambda_1(-\eta)]\,,$$

$$\mathcal{E}_\zeta(f) = \sqrt{1 - \int_{\mathbb{R}} d\eta [2G_\zeta(\eta) - G_{2\zeta}(\eta)] \exp\{W_2[f, f \circ \Lambda_1(\eta)] - W_2(f,f)\}},$$

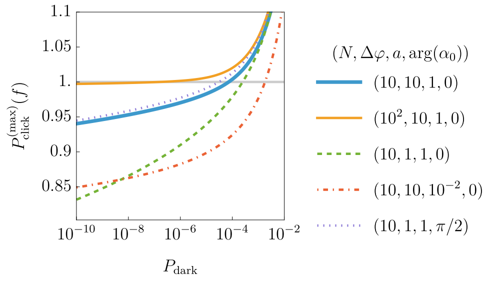
with $G_\zeta(\eta) =$ Gaussian, $J =$ reflection about t and x, $\Lambda_1(\eta) =$ boost along x, $W_2(f,f') =$ two-point Wightman function.

Photodetection scenario [3]

$\mathcal{O}_{\text{det}}$ = prismatic detector $l \times L \times L$ operating at time interval τ.
$|\psi\rangle$ = single-mode coherent state, with amplitude α_0, momentum $\mathbf{k}_0$ and coherence volume V_{coh}.
The same detector is used to measure the click probability P_{click} of $|\psi\rangle$ (panel a) and the dark-count probability P_{dark} in the vacuum (panel b):



Upper bound $P^{(\text{max})}_{\text{click}}(f)$ on the click probability $P_{\text{click}}(f)$ vs dark count P_{dark}:



$N := |\alpha_0|^2(l+\tau)(L+\tau)^2/V_{\text{coh}}$ number of photons "seen" by the detector, $\Delta\varphi := k_0(l+\tau)$ cumulative optical phase, $a := (l+\tau)/(L+\tau)$ aspect ratio.

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Samuel Fedida

Foundations of Relational Quantum Field Theory

Foundations of Relational Quantum Field Theory: Scalars

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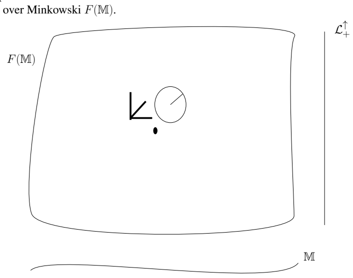
Abstract

We develop foundations for a relational approach to quantum field theory (RQFT) based on the operational quantum reference frames (QRFs) framework considered in a relativistic setting. Unlike other efforts in combining QFT with QRFs, we use the latter to provide novel mathematical and conceptual foundations for the former. We focus on scalar fields in Minkowski spacetime and discuss the emergence of relational local observables and fields from the consideration of Poincaré-covariant frame observables defined over the space of inertial reference frames. We recover a relational notion of Poincaré covariance, with transformations on the system directly linked to the state preparations of the QRF. We introduce and analyse various causality conditions. The theory makes direct contact with established foundational approaches to QFT: the frame smearing functions describing the QRF's localisation uncertainty plays the role of the Wightmanian test functions, and the algebras generated by relational local observables suitably extend the core axioms of Algebraic QFT.

Relativistic Quantum Reference Frames

Start with a classical inertial reference frame, which can be thought of as a clock and rods, sharply localised in spacetime and with a definite Lorentz orientation. Such a tetrad lives in the (orthonormal) frame bundle over Minkowski $F(\mathbb{M})$.



Suppose you want to describe a system in state $\rho = \rho^{(X)} \in \mathscr{D}(\mathcal{H}_\mathcal{S})$ with system Hilbert space $\mathcal{H}_\mathcal{S}$, relative to your inertial reference frame $X \in F(\mathbb{M})$. The state relative to another inertial reference frame $X' = (x, \lambda) \cdot X$ is $\rho^{(X')} = U_\mathcal{S}(x,\lambda)^\dagger \rho U_\mathcal{S}(x,\lambda)$.

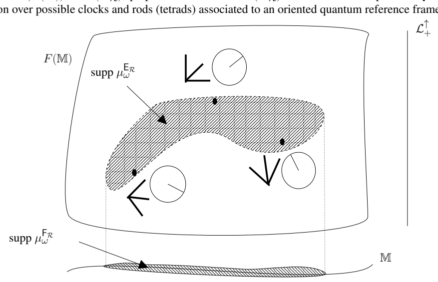
Perhaps you have some uncertainty regarding your inertial reference frame: you assign a probability distribution $\mu \in \mathrm{Prob}(F(\mathbb{M}))$ so that the proper mixed state that you should assign to describe experiments is

$$\rho^{(\mu)} = \int_{F(\mathbb{M})} \rho^{(X)} d\mu(X) = \int_{\mathcal{P}_+^\uparrow} U_\mathcal{S}(x,\lambda)^\dagger \rho U_\mathcal{S}(x,\lambda) d\mu(x,\lambda)$$

or, in a Heisenberg-like picture with fixed states, the observables now look like

$$\phi^{(\mu)} = \int_{\mathcal{P}_+^\uparrow} U_\mathcal{S}(x,\lambda) \rho U_\mathcal{S}(x,\lambda)^\dagger d\mu(x,\lambda)\,.$$

Suppose now that your uncertainty stems from the fact that your measurement apparatus (clock and rods) is fundamentally quantum mechanical: based on $\mathcal{H}_\mathcal{R}$, with Poincaré covariant POVM $\mathsf{E}_\mathcal{R} : \mathrm{Bor}(F(\mathbb{M})) \to \mathcal{E}(\mathcal{H}_\mathcal{R})$, prepared in state $\omega \in \mathscr{D}(\mathcal{H}_\mathcal{R})$. We have a Born probability distribution over possible clocks and rods (tetrads) associated to an oriented quantum reference frame.



Then the **relational local observable** relative to your **quantum reference frame** (QRF) is given by

$$\yen_\omega^\mathcal{R}(\phi) = \phi^{\left(\mu_\omega^{\mathsf{E}_\mathcal{R}}\right)} = \int_{\mathcal{P}_+^\uparrow} U_\mathcal{S}(x,\lambda)\rho U_\mathcal{S}(x,\lambda)^\dagger d\mu_\omega^{\mathsf{E}_\mathcal{R}}(x,\lambda)\,.$$

The relational local observable that one recovers after preparing one's detector (QRF) in a state ω is $\yen_\omega^\mathcal{R}(\phi)$, i.e. to every system operator $\phi \in \mathcal{B}(\mathcal{H}_\mathcal{S})$ and detector modelled by the QRF $\mathcal{R} = (U_\mathcal{R}, \mathsf{E}_\mathcal{R}, \mathcal{H}_\mathcal{R})$ which can be prepared in states ω we consider the **relational quantum field**

$$\hat{\Phi}^\mathcal{R} : \mathscr{D}(\mathcal{H}) \ni \omega \mapsto \yen_\omega^R(\phi) = \int_{\mathcal{P}_+^\uparrow} U_\mathcal{S}(x,\lambda)\phi U_\mathcal{S}(x,\lambda)^\dagger \, d\mu_\omega^{\mathsf{E}_\mathcal{R}}(x,\Lambda) \in \mathcal{B}(\mathcal{H}_\mathcal{S})\,.$$

How does that compare to the usual notion of "quantum field" in physics? The "fuzziness" in the space of inertial reference frames projects to the localisation in spacetime. Consider the disintegration $d\mu_\omega^{\mathsf{E}_\mathcal{R}}(x,\lambda) = d\nu_\omega^{\mathsf{E}_\mathcal{R}}(\lambda \mid x) d\mu_\omega^{\mathsf{F}_\mathcal{R}}(x)$, where $\mathsf{F}_\mathcal{R} : \mathrm{Bor}(\mathbb{M}) \to \mathcal{E}(\mathcal{H}_\mathcal{R})$.

$$\hat{\Phi}^\mathcal{R}(\omega) = \int_\mathbb{M} \hat{\phi}_\omega^\mathcal{R}(x) d\mu_\omega^{\mathsf{F}_\mathcal{R}}(x) = \int_\mathbb{M} \hat{\phi}_\omega^\mathcal{R}(x) f_\omega^{\mathsf{F}_\mathcal{R}}(x) d^d x$$

where $f_\omega^{\mathsf{F}_\mathcal{R}} \equiv \frac{d\mu_\omega^{\mathsf{F}_\mathcal{R}}}{d^d x}$ is the **frame smearing function**, and

$$\hat{\phi}_\omega^\mathcal{R} : \mathbb{M} \ni x \mapsto \int_{\mathcal{L}_+^\uparrow} U_\mathcal{S}(x,\lambda)\phi U_\mathcal{S}(x,\lambda)^\dagger \, d\nu_\omega^{\mathsf{E}_\mathcal{R}}(\lambda \mid x) \in \mathcal{B}(\mathcal{H}_\mathcal{S})$$

is a (pointwise) **relational local quantum field**. It is an integral kernal for $\hat{\Phi}^\mathcal{R}$ which always exists.

Covariance

$$U_\mathcal{S}(a,\Lambda)\hat{\Phi}^\mathcal{R}(\omega)U_\mathcal{S}(a,\Lambda)^\dagger = \hat{\Phi}^\mathcal{R}(U_\mathcal{R}(a,\Lambda)\omega U_\mathcal{R}(a,\Lambda)^\dagger) \quad \Big(\sim (a,\Lambda)\cdot\hat{\Phi}_W(f) = \hat{\Phi}_W(\{a,\Lambda\}\cdot f)\Big)$$

$$U_\mathcal{S}(a,\Lambda)\hat{\phi}_\omega^\mathcal{R}(x)U_\mathcal{S}(a,\Lambda)^\dagger = \hat{\phi}^\mathcal{R}_{U_\mathcal{R}(a,\Lambda)\omega U_\mathcal{R}(a,\Lambda)^\dagger}(\Lambda x + a) \quad \Big(\sim (a,\Lambda)\cdot\hat{\phi}(x) \text{``}=\text{''} \hat{\phi}(\Lambda x + a)\Big)$$

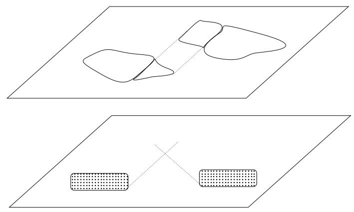
Relational causality

A weak epistemic condition (which is analogous to local commutativity in Wightman QFT) for relational causality is **relational Einstein causality**:

$$\text{supp } \mu_{\omega_1}^{\mathsf{F}_{\mathcal{R}_1}} \perp\!\!\!\perp \text{supp } \mu_{\omega_2}^{\mathsf{F}_{\mathcal{R}_2}} \Longrightarrow \left[\hat{\Phi}^{\mathcal{R}_1}(\omega_1), \hat{\Phi}^{\mathcal{R}_2}(\omega_2)\right] = \left[\hat{\Phi}^{\mathcal{R}_1}(\omega_1)^\dagger, \hat{\Phi}^{\mathcal{R}_2}(\omega_2)\right] = 0\,.$$

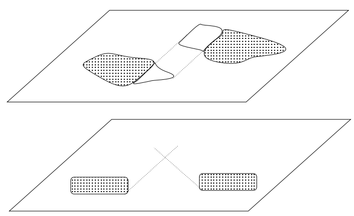
A (stronger) ontic condition (now at the level of kernels) is **weak relational microcausality**:

$$\text{supp } \mu_{\omega_1}^{\mathsf{F}_{\mathcal{R}_1}} \perp\!\!\!\perp \text{supp } \mu_{\omega_2}^{\mathsf{F}_{\mathcal{R}_2}} \Longrightarrow \left[\hat{\phi}_{\omega_1}^{\mathcal{R}_1}(x_1), \hat{\phi}_{\omega_2}^{\mathcal{R}_2}(x_2)\right] = \left[\hat{\phi}_{\omega_1}^{\mathcal{R}_1}(x_1)^\dagger, \hat{\phi}_{\omega_2}^{\mathcal{R}_2}(x_2)\right] = 0 \quad \forall x_1, x_2 \in \mathbb{M}\,.$$



An even stronger ontic condition is **strong relational microcausality**:

$$x_1 \perp\!\!\!\perp x_2 \Longrightarrow \left[\hat{\phi}_{\omega_1}^{\mathcal{R}_1}(x_1), \hat{\phi}_{\omega_2}^{\mathcal{R}_2}(x_2)\right] = \left[\hat{\phi}_{\omega_1}^{\mathcal{R}_1}(x_1)^\dagger, \hat{\phi}_{\omega_2}^{\mathcal{R}_2}(x_2)\right] = 0 \quad \forall \omega_1, \omega_2\,.$$



Finite-precision ($(x-y)^2 > -\sigma$) causality conditions with respect to operationally meaningful preparations of the frame ($\mathfrak{S}_\mathcal{R} \subset \mathscr{D}(\mathcal{H}_\mathcal{R})$), leading to (strong/weak) $(\mathfrak{S}_\mathcal{R}, \sigma)$-(micro)causality can also be written down and a causal covariant relational quantum field can be constructed explicitly [1].

Relational QFT and Algebraic QFT

Let $\mathcal{R}$ be a relativistic QRF, $\mathfrak{S}_\mathcal{R} \subseteq \mathscr{D}(\mathcal{H}_\mathcal{R})$ be convex, $\sigma \geq 0$ and $\mathcal{O}_\mathcal{S} \subseteq \mathcal{B}(\mathcal{H}_\mathcal{S})$ be $(\mathfrak{S}_\mathcal{R}, \sigma)$-causal. Given a subset $U \subseteq \mathbb{M}$, we call

$$\mathcal{A}_{\mathcal{O}_\mathcal{S}}^{(\mathfrak{S}_\mathcal{R},\sigma)}(U) := \{\hat{\Phi}^\mathcal{R}(\omega) \mid \phi \in \mathcal{O}_\mathcal{S}, \omega \in \mathfrak{S}_\mathcal{R} \text{ s.t. supp } \mu_\omega^{\mathsf{F}_\mathcal{R}} \subset U\}''$$

a *relational local algebra*. Then

1. (Isotony) For all $U \subseteq V \subseteq \mathbb{M}$, $\mathcal{A}_{\mathcal{O}_\mathcal{S}}^{(\mathfrak{S}_\mathcal{R},\sigma)}(U) \subseteq \mathcal{A}_{\mathcal{O}_\mathcal{S}}^{(\mathfrak{S}_\mathcal{R},\sigma)}(V)$.
2. (Covariance) For all $(a,\Lambda) \in \mathcal{P}_+^\uparrow$ and $U \subseteq \mathbb{M}$, $(a,\Lambda)\cdot\mathcal{A}_{\mathcal{O}_\mathcal{S}}^{(\mathfrak{S}_\mathcal{R},\sigma)}(U) = \mathcal{A}_{\mathcal{O}_\mathcal{S}}^{((a,\Lambda)\cdot\mathfrak{S}_\mathcal{R},\sigma)}((a,\Lambda)\cdot U)$.
3. (Causality) For all $U \perp\!\!\!\perp_\sigma V \subset \mathbb{M}$, $\left[\mathcal{A}_{\mathcal{O}_\mathcal{S}}^{(\mathfrak{S}_\mathcal{R},\sigma)}(U), \mathcal{A}_{\mathcal{O}_\mathcal{S}}^{(\mathfrak{S}_\mathcal{R},\sigma)}(V)\right] = 0$.

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Sebastian Schuster

Canonical Quantum Gravity without Canonical General Relativity

Canonical Quantum Gravity without Canonical General Relativity

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Looking for me?

Want to ask questions or discuss, but I am not at my poster?
Please look for the person in the picture! ☺
(Though with longer & unrulier beard…)
(From the back, look for a long, low ponytail!)

Context, Goals, and Methods

(Beyond) Canonical Quantum Gravity

- **Canonical quantum gravity:**
  1. Assume global hyperbolicity,
  2. Canonically quantize Hamiltonian formulation of general relativity,
  3. Arrive at Wheeler–DeWitt (WDW) equation,
  $$\hat{H}|\Psi\rangle\rangle = 0,$$
  for state of the Universe Ψ.
- It has the **problem of time**, as the WDW equation is a time-independent Schrödinger equation with 0 energy.
- **Successor theories can be richer** than the assumptions going into their derivation.
- Our approach to **study time travel:**
  1. Assume WDW equation,
  2. Use periodic (relational) clocks and periodic systems as constituents of full Hamiltonian constraint $\hat{H}$,
  3. Model time travel with these toy models.
- Importantly, the global hyperbolicity in the original derivation of the WDW equation is explicitly *not* assumed.

Goals

- **Distinguish** the following cases of systems with relational dynamics:
  1. System with a periodic clock;
  2. System with a periodic clock and a memory/calendar;
  3. System undergoing time travel (without Novikov self-consistency condition);
  4. System undergoing time travel (with Novikov self-consistency condition).
- *Can* we even distinguish cases 2 and 3?
- What role do **interactions** play?
- First, study toy models.
- Aim for more general results about time travel when faced with the problem of time of canonical quantum gravity.
- Demonstrate usability of the WDW equation in absence of the classical assumptions that went into its derivation.

Method: The Page–Wootters Formalism

No parameter time, but (sub)systems can evolve w.r.t. a **relational/emergent time** [PW83].

- Essentially:
  1. Separate full Hilbert space: $\mathcal{H} = \mathcal{H}_C \otimes \mathcal{H}_S$
  2. Introduce a clock Hamiltonian $\hat{H}_C$ of a subsystem
  3. Fix the clock state ψ_C at some chosen, initial time
  4. Define time through evolution of this state with $\hat{H}_C$
  5. Measure time evolution of an operator $\hat{A}$, stationary w.r.t. $\hat{H}_C$, as
  $$E(A|\tau) = \mathrm{tr}\Big(\hat{A}\hat{P}_\tau\hat{\rho}\Big) / \mathrm{tr}\Big(\hat{P}_\tau\hat{\rho}\Big),$$
  where
  $$\hat{P}_\tau = |\psi_C(\tau)\rangle\langle\psi_C(\tau)| \otimes \mathbb{1}_S, \qquad \text{and} \qquad \hat{\rho} \in \mathcal{L}(\mathcal{H})$$
- Early counterarguments have recently been tackled, providing a unified picture, with clocks as a gauge to be chosen. See [Pau80; UW89; Kuc11; GLM15; MV17; HSL21; Cha+26].

A First Toy Model and First Extensions [ASV24]

A Physicist's Answer to Everything

- Harmonic oscillators!
- This time, two: Uncoupled, for variables a ('clock') and χ ('system'):
$$\hat{H}\Psi(a,\chi) = (\underbrace{\frac{\partial^2}{\partial a^2} - \omega_a^2 a^2}_{\hat{H}_C} \underbrace{- \frac{\partial^2}{\partial \chi^2} + \omega_\chi^2 \chi^2}_{-\hat{H}_S})\Psi = 0$$
- Normalizability of Ψ gives commensurability condition:
$$\frac{\omega_a}{\omega_\chi} = \frac{2n_\chi + 1}{2n_a + 1}, \qquad n_a, n_\chi \in \mathbb{N}_0$$
- Now:
  - Apply the **PW formalism** for the cyclic time operator given through the **POVM methods** (see next box).
  - Commensurability implements a notion of **Novikov self-consistency** [Nov92]—how restrictive is this?

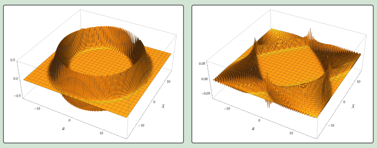

Figure: Example states Ψ fulfilling the WDW of the first toy model.

Time Operators & POVMs

- The **harmonic oscillator 'clock'** [BGL94; BGL95]:
  - Clock Hamiltonian
  $$\hat{H}_C = \hat{n}_C + \frac{1}{2}\mathbb{1}_C.$$
  - Define non-unitary $\hat{W}$ through
  $$\hat{a} = \hat{W}\widehat{|a|}, \qquad \text{with} \qquad \widehat{|a|} := \hat{n}^{1/2}$$
  having improper eigenstates $|\theta\rangle$
  $$\hat{W}|\theta\rangle = e^{i\theta}|\theta\rangle, \qquad \text{with} \qquad |\theta\rangle = \sum_{n\geq 0} e^{in\theta}|n\rangle.$$
  - The relevant POVM:
  $$B_0(X) := \frac{1}{2\pi}\int_X \mathrm{d}\theta\, |\theta\rangle\langle\theta| = \sum_{n,m\geq 0}\frac{1}{2\pi}\int_X e^{i(n-m)\theta}\,\mathrm{d}\theta\, |n\rangle\langle m|.$$
  - Giving one of many possible time operators as:
  $$\hat{T}_0 = B_0(\theta) = \sum_{n\neq m\geq 0}\frac{1}{i(n-m)}|n\rangle\langle m| + \pi\mathbb{1}.$$
  - The index 0 indicates the freedom to choose where on the unit circle the clock starts.
- As the clock is periodic, this avoids already many of the usual counter-arguments to time-operators.

Results

- Time evolution happens w.r.t. a chosen initial clock state $|\psi(0_*)\rangle_C \in \mathcal{H}_C$ which the clock Hamiltonian evolves to
$$|\psi(\theta_*)\rangle_C := e^{-i\hat{H}_C\theta_*}|\psi(0_*)\rangle_C,$$
with associated projectors
$$\hat{P}_{\theta_*} := (|\psi(\theta_*)\rangle_C\, {}_C\langle\psi(\theta_*)|) \otimes \mathbb{1}_S.$$
- Observables should not change the state of the Universe; focus on observables compatible with the system Hamiltonian $\hat{H}_S$.
- For a general state fulfilling the constraint,
$$|\Psi\rangle\rangle = \sum_{n,n'} A_{n,n'}|n\rangle_C \otimes |n'\rangle_S,$$
the evolution of the system according to clock time is given by conditional probabilities
$$E(\hat{H}_R|\theta_*) = \frac{\mathrm{tr}\Big(\hat{H}_S\hat{P}_{\theta_*}\hat{\rho}_\Psi\Big)}{\mathrm{tr}\Big(\hat{P}_{\theta_*}\hat{\rho}_\Psi\Big)}.$$
- After some calculation:
$$E(\hat{H}_S|\theta_*) = E(\hat{H}_S|0_*).$$

Self-consistency [Nov92] meets boring physics conjecture [Vis96]; consistent with results in [Cha+26].

Relationally, nothing ever happens.

Possible Extensions

- **Generalize:** Subdivide the Hilbert space up to 4-partite:
$$\mathcal{H} = \mathcal{H}_{\mathrm{env}} \otimes \mathcal{H}_C \otimes \mathcal{H}_S \otimes \mathcal{H}_M$$
where
  - $\mathcal{H}_{\mathrm{env}}$: **Environment**,
  - $\mathcal{H}_C$: **Clock**,
  - $\mathcal{H}_S$: **System**,
  - $\mathcal{H}_M$: **Memory** or calendar.
- Add interactions. *Work in progress:*
  - Adapt a **Jayne–Cummings model**,
  $$\left(\frac{\omega_C}{2}\hat{\sigma}_z - \omega_S\hat{a}^\dagger\hat{a} + g\big[\hat{\sigma}_+\hat{a} + \hat{\sigma}_-\hat{a}^\dagger\big]\right)|\Psi\rangle\rangle = 0,$$
  to the Wheeler–DeWitt equation.
  - Vanishing interaction g again yields $E(\hat{H}_S|\theta_*) = E(\hat{H}_S|0_*)$.
- Could our approach be extended to **relational rulers** to study, *e.g.*, **topology change instead of time travel** as a move beyond the globally hyperbolic origin of the WDW equation?

The Long Game: Observational Entropy?

Can we find **entropic arguments for/against time travel**?

- Entropy in closed systems (like the WDW equation) is difficult to define.
- Potential option: **Observational entropy** [ŠDA19] provides a notion of entropy for closed systems where
$$\hat{H} = \hat{H}^{(1)} + \cdots + \hat{H}^{(m)} + \epsilon\hat{H}_{\mathrm{int}}$$
on a Hilbert space
$$\mathcal{H} = \mathcal{H}^{(1)} \otimes \cdots \otimes \mathcal{H}^{(m)}:$$

1. Pick conserved quantities $\hat{A}_1, \ldots, \hat{A}_n$, spectral decompose them, form spectral projection operators $\hat{P}_i$;
2. Define
$$V_{i_1,\ldots,i_n} := \mathrm{tr}\Big(\hat{P}_{i_n}\ldots\hat{P}_{i_2}\hat{P}_{i_1}\hat{P}_{i_2}\ldots\hat{P}_{i_n}\Big),$$
$$p_{i_1,\ldots,i_n} := \mathrm{tr}\Big(\hat{P}_{i_n}\ldots\hat{P}_{i_1}\hat{\rho}\hat{P}_{i_1}\ldots\hat{P}_{i_n}\Big);$$
3. Get **coarse-grained entropy** even for closed systems as
$$S_{O(\mathcal{C}_1,\ldots,\mathcal{C}_n)}(\hat{\rho}) = -\sum_{i_1,\ldots,i_n} p_{i_1,\ldots,i_n}\ln\left(\frac{p_{i_1,\ldots,i_n}}{V_{i_1,\ldots,i_n}}\right),$$
where $O(\mathcal{C}_1,\ldots,\mathcal{C}_n)$ is an ordered set of collections C_i of above projection operators.

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Tadeusz Adach

κ-deformed Discrete Symmetries

κ-deformed discrete symmetries

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Introduction

Quantum gravity is anticipated to profoundly alter the structure of spacetime. It can be hypothesized that there exists a physical regime wherein quantum gravity effects lead to a replacement of local quantum field theory with a **quantum-deformed field theory**, with the deformation scale identified with the Planck energy scale. A potential realization of this hypothesis involves replacing the Poincaré algebra of Minkowski spacetime symmetries with its deformed Hopf algebra counterpart, the **κ-Poincaré algebra**, which has become a prominent tool in **quantum gravity phenomenology**. We investigate the impact of κ-deformation on discrete symmetries of free fields.

1. κ-Minkowski

κ-Minkowski spacetime is characterized by the relations

$$[x^0, x^j] = \frac{i}{\kappa} x^j, \qquad [x^j, x^k] = 0$$

Its symmetries are described by the **κ-Poincaré quantum group**. Its corresponding **Hopf algebra** can be realized in different bases, which express the deformation in different ways. In our approach, the classical basis is used, in which the algebra sector of κ-Poincaré is undeformed, but translations and boosts have non-primitive coproducts (or, in plain language, deformed Leibniz rules), e.g.

$$P_j \triangleright (\phi \star \psi) = (P_j \triangleright \phi) \star \left(\frac{1}{\kappa} \sqrt{\kappa^2 + P_0^2 - \mathbf{P}^2} \triangleright \psi \right) + \phi \star (P_j \triangleright \psi)$$

with respect to the (non-commutative) $\star$-product.

2. κ-momentum space

The momentum space of κ-Minkowski is a **submanifold of dS_4** with κ playing the role of curvature radius:

$$p_0^2 - \mathbf{p}^2 - p_4^2 = -\kappa^2$$

with the additional constraints

$$p_+ = p_0 + p_4 > 0, \qquad p_4 > 0.$$

While the mass-shell relation is undeformed for both particles and antiparticles, the Hopf-algebraic properties of translation generators introduce an asymmetry in the **positive** and **negative** energy modes:

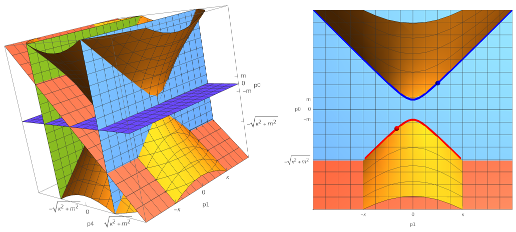

Fig. 1: (1+1)D depiction depiction of κ-momentum space

The negative energy modes are obtained from the positive through the **antipodal map**:

$$S(p_0) = -p_0 + \frac{\mathbf{p}^2}{p_+}, \qquad S(p_j) = -\frac{\kappa}{p_+} p_j$$

and their momenta are bounded by κ.

Crucially, the Hopf algebra structure also alters the notion of **plane wave conjugation**:

$$e^{ipx} \star \left(e^{ipx}\right)^\dagger = e^{ipx} \star e^{iS(p)x} = 1$$

which affects how the positive and negative energy modes of fields are constructed:

$$\phi(x) = \int \frac{d^3p}{\sqrt{2\omega_\mathbf{p} p_4/\kappa}} \left(a_\mathbf{p} e^{-ipx} + b_\mathbf{p}^\dagger e^{-iS(p)x} \right)$$

3. Charge conjugation

Because of the deformed wave conjugation, the translation operator acting on the field $(\partial_\mu \phi)$ produces **different momenta for the $a_\mathbf{p}$ and $b_\mathbf{p}$ modes**. By extension, the Noether charges are $\mathcal{C}$-asymmetric (in the natural sense of $\mathcal{C} : a_\mathbf{p} \mapsto b_\mathbf{p}$). For instance, the translation charge takes the form

$$P_\mu = \int \frac{d^3p}{p_4/\kappa} \left[\frac{p_+^3}{\kappa^3} p_\mu a_\mathbf{p}^\dagger a_\mathbf{p} - S(p_\mu) b_\mathbf{p}^\dagger b_\mathbf{p} \right]$$

"Particles" and "antiparticles" are thus **physically distinguishable** even in a free theory, living in two different momentum spaces, and **$\mathcal{C}$ symmetry is broken**.

One could of course try to enforce $\mathcal{C}$-invariance by symmetrizing the Lagrangian $(\phi^\dagger \star \phi + \phi \star \phi^\dagger)$, but inconsistent ordering turns out to be incompatible with κ-Poincaré invariance. Charge conjugation asymmetry appears to be necessary for a κ-Lorentz-invariant theory.

4. Space and time inversions

While $\mathcal{C}$-breaking seems to be inherent in the theory, whether parity and time reversal are broken is largely a matter of chosen definition, as there is a lot of **conceptual ambiguity** there.

The better question is whether the combined $\mathcal{CPT}$ symmetry can survive. As we have recently shown, at least one well-motivated **prescription exists where $\mathcal{CPT}$ remains an exact symmetry** (at least for free scalars and spinors). It is based on an undeformed parity transformation

$$\mathcal{P} : \phi(t, \mathbf{x}) \mapsto \phi(t, -\mathbf{x})$$

and modified time reversal, such that the property of antiunitarity is extended to deformed wave conjugation:

$$\mathcal{T} : \phi(t, \mathbf{x}) \mapsto \phi^{\dagger *}(-t, \mathbf{x})$$

where schematically

$$\phi(x) = \int dp\, \phi(p) e^{-ipx}, \qquad \phi^{\dagger *}(x) = \int dp\, \phi(p) e^{iS(p)x}$$

In effect, $\mathcal{T}^{-1} \left(\int \phi \star \psi \right) \mathcal{T} \propto \int \psi \star \phi$, and so $\int \phi \star \phi^\dagger \xrightarrow[\mathcal{CT}]{} \int \phi \star \phi^\dagger$

5. Extended charge algebra

The deformed symmetry behavior of translation (P_μ), rotation (M_j) and boost (N_j) Noether charges can be summed up in the following table (with the undeformed case on the left for comparison):

Charge	$\mathcal{C}$	$\mathcal{P}$	$\mathcal{T}$	$\mathcal{CPT}$
P_0	$+$	$+$	$+$	$+$
P_j	$+$	$-$	$-$	$+$
M_j	$+$	$+$	$-$	$-$
N_j	$+$	$-$	$+$	$-$

Charge	$\mathcal{C}$	$\mathcal{P}$	$\mathcal{T}$	$\mathcal{CPT}$
P_0^κ	$+\mathfrak{S}$	$+$	$+\mathfrak{S}$	$+$
P_j^κ	$+\mathfrak{S}$	$-$	$-\mathfrak{S}$	$+$
M_j^κ	$+\mathfrak{S}$	$+$	$-\mathfrak{S}$	$-$
N_j^κ	$+\mathfrak{S}$	$-$	$+\mathfrak{S}$	$-$

The involutive $\mathfrak{S}$ operation corresponds to the swapping of p and $S(p)$ momentum spaces while preserving sign.

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Tamal Guha

Minimal Assistance from Environment Unlocks the Encoding Strength

Environment Assistance Unlocks Optimal Encoding Strength

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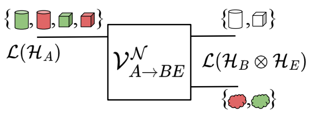
Highlights!

- **Generalizing the classical communication utility** in terms of encoding strength.
- **Environment assistance classical communication** is not optimal for all quantum channels.
- The traditional notion is *contradicted* when viewed from the lens of **encoding strength unlocking**.

Environment Assisted Classical Communication

- Quantum channel with $d_i > d_o$: Input symbols partially leaked into the environment
- Channel $\mathcal{N} : A \to B \implies$ Isometry $\mathcal{V}_{\mathcal{N}} : A \to BE$
- Limited assistance from environment can in principle retrieved these lost symbols



- Traditionally this *Environment Assisted Classical Capacity* (EACC) is measured in terms of the Mutual Information (MI) between the input random variable X and the jointly decoded output random variable Y
- Example - For any quantum channel EACC ≥ 1-bit:
  For any $\mathcal{V}_{\mathcal{N}}$ and for any pair of $|\psi_1\rangle \perp |\psi_2\rangle$, $\mathcal{V}_{\mathcal{N}}(|\psi_1\rangle) \perp \mathcal{V}_{\mathcal{N}}(|\psi_2\rangle)$
  Any pair of orthogonal states are LOCC distinguishable [Walgate *et.al.*, PRL (2000)]
- There are instances where EACC is always sub-optimal ($< \log d_i$-bits)
  [Watrous, PRL (2005); Duan *et.al.*, IEEE Trans. (2009)]
- But why should we measure EACC only in terms of MI? In single-shot scheme MI is not the unique measure of communication utility

Generalized Communication Framework

- Input Random Variable X: Output Random Variable $Y \implies \{p(y_j|x_i)\}_{i,j}$
- Channel Matrix: $[M(n,m)]_{i,j} = p(y_j|x_i)$, $\forall i \in [n]$ & $j \in [m]$.
- $\mathcal{P}^{n\to m}(\mathcal{N}(\mathcal{Q}_{d_A})) := \{$All $n \times m$ channel matrices generated by the channel $\mathcal{N} : \mathcal{L}(\mathcal{H}_{d_A}) \to \mathcal{L}(\mathcal{H}_{d_B})\}$

** $\mathcal{P}^{n\to m}(\mathcal{Q}_{d_A})$ means the channel $\mathcal{N}$ is identity channel

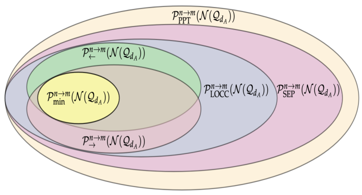
Instrumental Preliminaries

Bound on the Channel Matrix Set

Proposition 1. *For any quantum channel $\mathcal{N} : \mathcal{L}(\mathbb{C}^{d_A}) \mapsto \mathcal{L}(\mathbb{C}^{d_B})$, and for all $n, m \in \mathbb{N}$; $\mathcal{P}^{n\to m}(\mathcal{N}(\mathcal{Q}_{d_A})) \subseteq \mathcal{P}^{n\to m}(\mathcal{Q}_d)$, where $d = \min\{d_A, d_B\}$.*

Assistance from Environment

Depending upon the restrictions on decoding measurement performed on the environment and the output system at the receiver, there could be multiple degrees of environmental assistance:



**Minimal Assistance:* The decoding measurement can be implemented with no communication between the receiver and the environment.
$P \in \mathcal{P}^{n\to m}_{\mathrm{min}}(\mathcal{N}(\mathcal{Q}_{d_A}))$ whenever, $[P]_{ij} = p(y_j|x_i) = \mathrm{Tr}[(\Lambda^B_l \otimes \Lambda^E_k)\mathcal{V}_{\mathcal{N}}(\rho_i)]$, where $j := f(k,l)$ and $\{\Lambda^B_l\}_l$ and $\{\Lambda^E_k\}_k$ are the respective decoding POVMs performed on the Bob's and environment quantum system.
** *Trivially, if an information processing task is perfectly accomplished under minimal assistance, then it will remain true for any other setting.*

Unlocking of Encoding Strength

Definition 1. *Consider a channel $\mathcal{N} : \mathcal{L}(\mathbb{C}^{d_A}) \mapsto \mathcal{L}(\mathbb{C}^{d_B})$ with $d_A > d_B$, and $d' \leq d_B$ is the minimum dimension of a quantum system, such that, $\mathcal{P}^{n\to m}(\mathcal{N}(\mathcal{Q}_{d_A})) \subseteq \mathcal{P}^{n\to m}(\mathcal{Q}_{d'})$ for all $n, m \in \mathbb{N}$. Environment assistance is said to unlock the encoding strength for $\mathcal{N}$, whenever there exists a $P \in \mathcal{P}^{n'\to m'}_{\mathrm{env}}(\mathcal{N}(\mathcal{Q}_{d_A}))$ but $P \notin \mathcal{P}^{n'\to m'}(\mathcal{Q}_{d'})$.*

Bound on the Env. Asst.-Channel Matrix Set

Lemma 1. *For any quantum channel $\mathcal{N} : \mathcal{L}(\mathcal{H}_A) \mapsto \mathcal{L}(\mathcal{H}_B)$ with $d_A \geq d_B$, and for all $n, m \in \mathbb{N}$, $\mathcal{P}^{n\to m}_{\mathrm{env}}(\mathcal{N}(\mathcal{Q}_{d_A})) \subseteq \mathcal{P}^{n\to m}(\mathcal{Q}_{d_A})$, even with the complete access to the environment.*

Optimal Unlocking of Encoding Strength

Definition 2. *Any form of environment assistance is said to optimally unlock the encoding strength of $\mathcal{N} : \mathcal{L}(\mathbb{C}^{d_A}) \mapsto \mathcal{L}(\mathbb{C}^{d_B})$ (with $d_A \geq d_B$), if there exist an n-input-m-output $(n, m \in \mathbb{N})$ channel matrix $P^{n\to m}$, such that $P^{n\to m} \in \mathcal{P}^{n\to m}_{\mathrm{env}}(\mathcal{N}(\mathcal{Q}_{d_A}))$, but $P^{n\to m} \notin \mathcal{P}^{n\to m}(\mathcal{Q}_{\tilde{d}})$ for all $\tilde{d} < d_A$.*

Main Results

Channels with suboptimal EACC

Lemma 2 (Duan *et.al.* IEEE Trans. (2009))**.** *For any choice of $i, j \in \{0,1,2\}$ with $i \neq j$, consider an uncountable set of isometries $\mathrm{S}_{i,j}$ defined by,*

$$\mathrm{S}_{i,j} := \{\mathcal{V}_7 \mid \mathcal{V}_7 : \mathbb{C}^7 \mapsto \mathbb{C}^3 \otimes \mathbb{C}^3 \\ \& \ \mathrm{Range}(\mathcal{V}_7) \perp \mathrm{span}\{|\phi_3^+\rangle, |i\rangle \otimes |j\rangle\}\}$$

where $|\phi_3^+\rangle = \frac{1}{\sqrt{3}}\sum_{k=0}^{2} |k\rangle \otimes |k\rangle$.
EACC of all the channels $\mathcal{N}_7^{\mathcal{V}} : \mathcal{L}(\mathbb{C}^7) \mapsto \mathcal{L}(\mathbb{C}^3)$, induced from each isometry $\mathcal{V}_7 \in \mathrm{S}_{i,j}$ is suboptimal for all distinct i, j pair, even if the decoding measurements of both the receiver and environment are separable super-operators (SEP).

Contradiction!

Theorem 1. *All channels $\mathcal{N}_7^{\mathcal{V}}$ achieves optimal encoding strength with minimal assistance of environment.*

*** There exists a one-parameter class of channel matrices $M_7(p)$, such that, $M_7(p) \in \mathcal{P}^{7\to 7}(\mathcal{N}_7^{\mathcal{V}}(\mathcal{Q}_7))$ but $M_7(p) \notin \mathcal{P}^{7\to 7}(\mathcal{Q}_d), \forall d < 7$.

Classical Transmission Fidelity ($\mathcal{F}_c$): An Operational Interpretation

- For a given $n \times n$ channel matrix $M_n :\ \mathcal{F}_c(M_n) = \sum_{x=0}^{n-1} p(y = x|x) = \mathrm{Tr}(M_n)$
- For any quantum channel $\mathcal{N} : \mathcal{F}_c(\mathcal{N}) := \max_{n\in\mathbb{N}} \ \max_{P\in\mathcal{P}^{n\to n}(\mathcal{N})} \mathrm{Tr}[P]$

***Quantifies the maximal numbers of classical symbols that can be communicated with least possible error

- $\mathcal{F}_c^{env}(\mathcal{N}_7^{\mathcal{V}}) > \mathcal{F}_c(\mathcal{Q}_6)$

Strength of Environment Assistance

Theorem 2. *The minimal assistance from environment is powerful than the strongest 2-2-2 nonlocal correlations between sender and receiver:*

$$\mathcal{F}_c^{env}(\mathcal{N}_7^{\mathcal{V}}) > \mathcal{F}_c^{PR}(\mathcal{N}_7^{\mathcal{V}}) \geq \mathcal{F}_c^{SR}(\mathcal{N}_7^{\mathcal{V}})$$

Generalized Channels with suboptimal EACC

Lemma 3 (Watrous, PRL (2005))**.** *For every $d \geq 3$, consider an uncountable set of isometries S_d*

$$\mathrm{S}_d := \{\mathcal{V}_{d^2-1} \mid \mathcal{V}_{d^2-1} : \mathbb{C}^{d^2-1} \mapsto \mathbb{C}^d \otimes \mathbb{C}^d \\ \& \ \mathrm{Range}(\mathcal{V}_{d^2-1}) \perp |\phi_d^+\rangle\} \tag{1}$$

where $|\phi_d^+\rangle = \frac{1}{\sqrt{d}}\sum_{k=0}^{d-1} |k\rangle \otimes |k\rangle$.
EACC of all channels $\mathcal{N}_{d^2-1}^{\mathcal{V}} : \mathcal{L}(\mathbb{C}^{d^2-1}) \mapsto \mathcal{L}(\mathbb{C}^d)$, induced from each isometry $\mathcal{V}_{d^2-1}$ is suboptimal, even if the decoding measurements of both the receiver and the environment are separable super-operators (SEP).

General Contradiction!

Theorem 3. *For all $d \geq 3$, quantum channels $\mathcal{N}_{d^2-1}^{\mathcal{V}}$, corresponding to all $\mathcal{V}_{d^2-1} \in \mathrm{S}_d$, achieve optimal encoding strength under minimal assistance of the environment.*

Poster Based on

"Minimal Help, Maximal Gain: Environmental Assistance Unlocks Encoding Strength"; arXiv:2509.09340 (2025).

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Tomasz Bazylewicz

Decoherence from the light bending interaction

Decoherence from the light bending interaction

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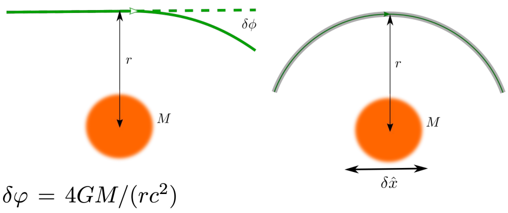
We analyse a decoherence effect, caused by the gravitational interaction between a massive body and the electromagnetic field. Assuming a quantum version of the light bending interaction, we show that it leads to decoherence of the mass if the light is not observed. Using the extreme weakness of the gravitational coupling, we derive explicitly the decoherence lengthscales for general states of the central mass and for both thermal and coherent light. Predictably, the effect is very faint for anything but hugely energetic light, however from the fundamental point of view of co-existence of both gravitation and quantum theories, it is there. Since effectively the studied system is a quantum optomechanical system, we hope our results, properly rescaled, will be also useful in optomechanics.

Gravitational light bending setup



$$\delta\varphi = 4GM/(rc^2)$$

We used tree level interaction between electromagnetic field and a mass [1]. The interaction was put into a scheme where light is in a semicircle waveguide around a massive object at a constant distance. Then interaction is expanded into a first order of quantum fluctuations of a mass.

Hamiltonian of Interaction

$$V_{k\nu} = -g_0(b^\dagger + b) \otimes \hbar\omega_k a^\dagger_{k\nu} a_{k\nu}$$

$$g_0 = \frac{4GM}{r^2c^2}\sqrt{\frac{\hbar}{2M\Omega}}$$

Exactly solvable Unitary Evolution

$$U^I_{int}(t) = U_{SE}(t)\left[\mathbf{1}_S \otimes \Phi(t)\right] \qquad \Phi(t) = \exp\left\{-ig_0^2\left[\Omega t - \sin(\Omega t)\right](H_E/\Omega)^2\right\}$$

$$U_{SE}(t) = \exp\left\{\left(\lambda(t)b^\dagger - \overline{\lambda}(t)b\right) \otimes H_E\right\} \qquad \lambda(t) = \frac{g_0}{\Omega}\left(e^{i\Omega t} - 1\right)$$

We used Magnus series to evaluate this evolution but this model is known in optomechanical literature [2].

Using separated initial state model is solvable

$$\rho_{SE}(0) = \rho_S \otimes \rho_E \longrightarrow \rho_t(x, x') = \int dp \rho_{0t}(x, x'; p) F_t(p; \Delta)$$

Free evolution

Enviorment Impact

$$\rho_{0t}(x, x'; p) = \sqrt{\kappa/\pi^2}\chi_t(-\Delta, p)e^{-i\xi_+ p} \qquad F_t(p; \Delta) = \sum_E p(E)\exp\left\{2iE\left[\gamma_1(t)p + \gamma_2(t)\Delta\right]\right\}$$

Using standart assumptions in open quantum system to analize decoherence we first assume un-correlated state to see how corelations emerge[3]. It turs out that we can solve this problem exactly separating free evolution of a matrix elements from environmental impact.

Results: Coherence Length

Using Talor series

$F_t(p; \Delta)$ $g_0 \ll 1$

High temperature multimode

$$\lambda_{coh} = \sqrt{\frac{\hbar^3\Omega\beta^2}{2g_0^2 M}}$$

Single mode light

$$\lambda_{coh} = \sqrt{\frac{\hbar\Omega}{2g_0^2 M}\frac{e^{\hbar\beta\omega} - 1}{\omega e^{\hbar\beta\omega/2}}}$$

Coherent state

$$\lambda_{coh} = \sqrt{\frac{\hbar\Omega}{2g_0^2\omega^2 M|\alpha|^2}}$$

We obtained Coherence Length shows us on which distance decoherence effect stars to play role then we obtained values of those for some realistic initial states of environment in each of those cases those parameters are small they depend on frequency of oscillator and on light parameters. To see significant effect we would need a temperature $T=10^{20}$ K.

Conclusions

We studied potential decoherence effects due to the gravitational light bending. We considered a massive harmonic oscillator interacting with quantized light at a fixed distance via classical light bending interaction of general relativity, proposed in [2]. Despite the fact that the impact of this decoherence is small, we show that if the postulated coupling scheme works, gravity mediated interaction will lead to decoherence effects. The effect is somewhat similar to the proposed universal gravitational decoherence [4].

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Varun Kushwaha

How Many Degrees of Freedom Does Field Theory Dynamics Use? Dynamical Compression Before Gravity

QFT With Area Many Degrees of Freedom

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The **holographic principle** states that the maximum entropy **S** that can be accumulated inside a finite region of space **R** (more precisely on the light-sheets of its boundary **δR**) equals the boundary area of that region divided by four times the Planck area[1],

Statistical entropy

$$S_{max} \leq \frac{A_{\delta R}}{4 l_p^2} \longrightarrow \mathcal{N}_{\text{dof},\,R} := \ln(\dim \mathcal{H}_R) \sim A_{\delta R}$$

However, this is inconsistent with standard Quantum Field Theory (QFT) simply because it is built out of volume-many commuting d.o.f.s, i.e., $[Q_r, P_s] = i\delta_{rs}$ with $r, s \sim$ Volume of the region.

That's fine when gravity is switched off. A common lesson learnt from the holographic principle, diffeomorphism invariance of observables[2], and from AdS/CFT[3], is that such algebra is only approximate and perhaps emergent (because QFT works really well!).

It's puzzling how nature manages to describe the low energy physics so well with a theory which apparently requires volume many d.o.f.s despite having a substantially smaller Hilbert space ("room").

Maybe **quantum information tools** can provide some insights?

Idea

Similar to the concept of **overlapping** qubits[4] one can try constructing *volume-many* overlapping field d.o.f.s[5]:

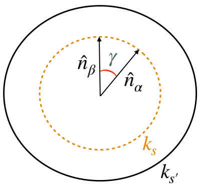
$$A_\alpha^{(s)} \sim \sum_{\ell=0}^{L_s \sim \text{area}} \sum_{m=-\ell}^{\ell} Y_{\ell m}(\hat{n}_\alpha)\, a_{\ell m}^{(s)}$$

Now with the new commutation relations:

$$[A_\alpha^{(s)}, A_\beta^{(s')\dagger}] \sim \delta_{ss'} \underbrace{\sum_{\ell=0}^{L_s} (2\ell+1) P_\ell(\hat{n}_\alpha \cdot \hat{n}_\beta)}_{C_{\alpha\beta}^{(s)}} , \quad [a_{\ell m}^{(s)}, a_{\ell' m'}^{(s')\dagger}] = \delta_{ss'}\delta_{\ell\ell'}\delta_{mm'}$$

Then write the field with a **UV** and **IR** cutoff as:

$$\phi_{\text{app}}(t,\mathbf{x}) = \sum_{s,\alpha} \frac{1}{\sqrt{2\omega_s V}} \left(u_\alpha^{(s)} A_\alpha^{(s)} + u_\alpha^{(s)*} A_\alpha^{(s)\dagger} \right) .$$

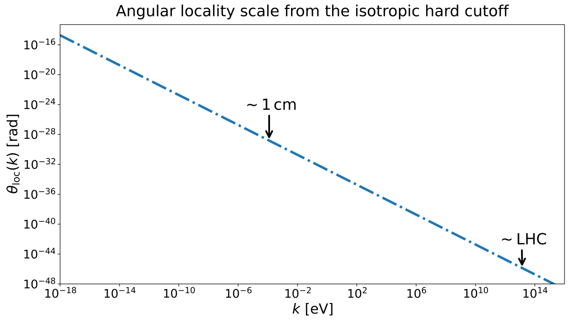


To see the effects one can start by looking at the commutators:

$$[\phi_{\text{app}}(t,\mathbf{x}), \pi_{\text{app}}(t,\mathbf{y})] = i \sum_s \mathcal{K}_s(\mathbf{x},\mathbf{y}) ,$$

$$[\phi_{\text{app}}(t,\mathbf{x}), \phi_{\text{app}}(t',\mathbf{y})] = -i \sum_s \frac{\sin(\omega_s(t-t'))}{\omega_s} \mathcal{K}_s(\mathbf{x},\mathbf{y}) .$$

Angular locality scale from the isotropic hard cutoff



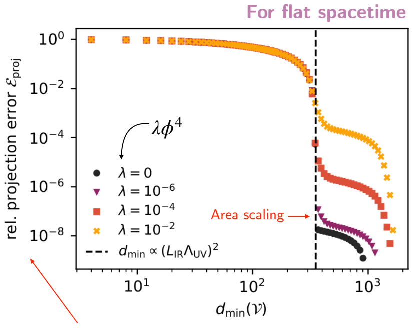
Area scaling from Classical Field Theory

These relations fixes the kinematic side of the theory. An interesting observation is made when we look at the dynamics.

Since, $C_{\alpha\beta}^{(s)}$ is a positive definite matrix, we can split $C = V^\top V$

This split is not unique, as the dimension n of $V_{N\times n}$ is flexible. This n determines the # of underlying d.o.f.s., and if we ask what is the n for which the Hamiltonian evolution of the field is best approximated then we find[6]:

For flat spacetime



$$\min_{V \in \mathbb{R}^{2N\times 2n}} \frac{\|T - VV^\top T\|_F}{\|T\|_F} \quad \text{subject to} \quad V^\top \mathbb{J}_{2N} V = \mathbb{J}_{2n} .$$

this ensures that the construction remains Hamiltonian

See our recent work for curved spacetimes and more details (arXiv:2602.09100)

Thermodynamics?

If the detectors are coupled to apparent d.o.f.s, then from the local density of states[7], obtained from the dyadic Green's function of the photon field, $\rho_{\text{EM}}(\omega,\mathbf{x}) = -\frac{2\omega}{\pi}\mathfrak{I}\,\text{Tr}(G^R_{\text{app}}(\omega;\mathbf{x},\mathbf{x}))$, we get:

this scales as ω_s^2

$$\mathcal{U}(\omega_s, T) = \hbar\omega n_B(\omega_s)\rho_{\text{EM}}(\omega_s) = \frac{2N_s}{V\Delta\omega_s}\frac{\hbar\omega_s}{e^{\beta\hbar\omega_s}-1}$$

This is the standard Planck's distribution!

Outlook

How well can area-many degrees of freedom emulate standard QFT?

So far: surprisingly well --- standard thermodynamic and locality structures looks the same for low energy experiments, furthermore, one can even define one-particle wavepackets with no modification to its dispersion relation.

Next: refine the algebraic/physical motivation and identify where the construction predicts observable deviations.

References