Pattern of perturbations from a coherent quantum inflationary horizon

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It is proposed that a coherent quantum inflationary horizon produces relic primordial curvature perturbations with a distinctive pattern of nonlocal correlations. A holographic quantum model is used to derive an uncertainty for space-time localization on causal diamonds, and to formulate candidate directional symmetries of emergent curvature perturbations on the horizon: the variance on any great circle is equal to the global variance, the mean value on any axis is independent of its equatorial mean, and the values are antipodally anticorrelated. These symmetries are shown to account for some well-known anomalies in cosmic microwave background anisotropy, and to sharpen predictions for new tests that avoid the usual ambiguity imposed by cosmic variance on large scales.

I. INTRODUCTION

A standard cosmological model\cite{1} is now supported by a considerable body of evidence, especially precise measurements of correlations in cosmic microwave background radiation (CMB)\cite{2,13}. The early evolution is generally described by slow-roll inflation\cite{14–16}, during which the repulsive gravity of an exotic scalar inflaton field drives an accelerating expansion. Inflation shapes the structure of the universe on the largest scales—a large, nearly-flat geometry, with nearly-scale-invariant primordial perturbations in curvature that give rise to cosmic structure. In the standard picture, the cosmic perturbations are created by vacuum fluctuations of a quantum field, the inflaton, coupled by linearized gravity to a classical background geometry.

It is also possible that the cosmic perturbations during inflation arise from new fundamental degrees of freedom of Planck scale quantum geometry. In such holographic or “spooky” scenarios\cite{17,18}, the perturbations are associated with the noisy emergence of space-time locality. Unlike the standard inflaton scenario based on quantized field modes, spooky inflation posits that the inflationary horizon of every observer, defined as its past light cone at the end of inflation, is a coherent nonlocal quantum object, similar to the horizon in quantum models of black holes\cite{19,20}. Spatial projections and correlations of quantum states collapse onto null surfaces instead of plane waves, so that primordial perturbations are laid down with nonlocal quantum correlations on the spherical horizon in all directions.

The predicted primordial power spectrum of curvature in the spooky picture is indistinguishable from standard inflation\cite{17,18}: spooky inflation also produces perturbations with a nearly scale invariant, slightly tilted power spectrum, so it preserves the precise match of standard cosmology to a host of measurements that depend only on the power spectrum, including CMB anisotropy spectra and measurements of cosmic large structure over a large range of scales. Unlike standard inflation, the dimensionless perturbation power \( \langle A^2 \rangle \equiv H t_P \) now depends only on \( H \), the expansion rate during inflation, and the Planck time \( t_P \equiv \sqrt{\hbar G/c^5} \). An important new observable effect\cite{18} is that nonlocal entanglement produces perturbations with exotic correlations in phase and direction not present in the standard picture.

This paper formulates a framework for new, precise direct tests of the spooky hypothesis, based on generic directional symmetries of spooky nonlocal geometrical directional correlations, and their measurable signatures in large-angle CMB temperature anisotropy. A simple quantum system, based on a standard spin algebra, is introduced as a model for nonlocal degrees of freedom that govern geometrical nonlocality on light cones in flat space-time. When the space-time structure of its quantum uncertainty is matched with correlations of perturbations on the classical inflationary horizon, the model serves to identify precisely-defined candidate symmetries in the directional pattern of perturbations on the largest scales.

These specific new nonlocal symmetries, not present in the standard perturbations originating from field vacuum states, can lead to distinctive measurable directional correlations of perturbations on the cosmological last scattering surface, including correlations of fine-grained structure at large angular separations. The holographic absence of one independent rotational degree of freedom, together with rotational symmetry in the emergent system, can lead to exact symmetries in correlations of curvature at large angles, particularly for the mean and variance of curvature on great circles and their polar directions, and possibly at other particular angular separations. Perturbations can also display antipodal anticorrelation, similar to some models of quantum black hole horizons, partially broken in the cosmological case by the time asymmetry of the inflationary background.

These symmetries are shown here to lead to specific signatures in the pattern of anisotropy in the CMB. Spooky directional correlations could provide a unified physical explanation of some long-known, seemingly unrelated statistical anomalies in the CMB\cite{4,8,13}. They also lead to predictions for new, more sharply formulated statistical comparisons with standard inflationary models, including new tests with existing datasets.
II. MODEL OF COHERENT QUANTUM CAUSAL DIAMOND FLUCTUATIONS

Hearkening back to Wheeler’s vision that quantum gravity should “take as the basis of our description of nature the elementary concepts of world line and light cone,” \[22\] the coherent quantum elements of geometry are assumed here to have the space-time structure not of quantum fields, but of coherent null surfaces: light cones, causal diamonds and horizons. Classical space-time is an emergent statistical behavior \[23\]–\[26\], and curvature on large scales represents a collective activity of many null elements.

In this view, not only geometry, but locality itself, is an emergent property of a quantum system. To address the physical consequences of geometrical nonlocality, it is necessary to address the connection of quantum mechanics and space-time at a basic level.

This section develops a covariant model of emergent locality based on coherent quantum light cone states. The model invokes new nonlocal geometrical degrees of freedom, not included in standard quantized linear gravity, that could have measurable physical effects on macroscopic scales, both in the laboratory and in cosmology.

A. Precedents for coherent quantum light cones

The “toy model” of coherent light cone states developed here does not address quantum dynamics at the Planck scale, nor is it a substitute for other, arguably more fundamental theories of quantum gravity \[27\]–\[29\]. It is introduced to reveal a particular measurable physical effect of coherent light cone states: the possible new symmetries of correlations arising from geometrical nonlocality in a regime not easily accessible in other approaches, much smaller than the classical radius of curvature but much larger than \(ct\).

Nonlocally coherent geometrical states are motivated by long standing information paradoxes of thermodynamic geometry with field states \[30\]–\[32\], and are required for a universal holographic scaling of geometrical information \[33\]–\[35\]. In some quantum models of black holes \[19\]–\[21\], information paradoxes are resolved with coherent horizon states, along with a semiclassical treatment of back-reaction. Nonlocal coherence of geometry is widely studied for finite intervals in bulk anti-de Sitter space \[36\]–\[39\].

As shown below, our model degrees of freedom have a holographic information content, as required to account for classical gravitation as an emergent behavior \[23\]–\[26\]. In this framework, their excitations are associated with the emergence of classical curvature on large scales, such as the cosmological perturbations considered here. The symmetries proposed here could provide evidence for more comprehensive models of emergent space-time and cosmology based on coherent causal diamonds \[17\].

These degrees of freedom have no dynamics or propagating wave solutions, and do not transport energy or information. However, they create quantum relationships that place measurable physical constraints on nonlocal, spacelike correlations. They describe relational information about positions in time and space that resemble standard non-dynamical spooky quantum relationships among particle states \[40\]. Their physical effects become apparent in measurements that project positions onto time and space coordinates in the rest frame of an observer. This can happen either in certain laboratory experiments, or as discussed at more length here, during inflation.

B. Dirac light cone function

Locality on light cones is conventionally defined by the covariant four-dimensional generalization of the one-dimensional Dirac \(\delta\)-function (ref. \[41\], §75):

\[
\Delta(x) = 2\delta(x^\mu x_\mu)x_0/|x_0|, \tag{1}
\]

where \(x^\mu = (t, \vec{r})\) represents 4-position. It vanishes at the origin, and is nonvanishing on past and future light cones from the origin. It is odd in timelike directions and even at spacelike separations, with a 4D point-parity antisymmetry that combines time and space,

\[
\Delta(-x) = -\Delta(x). \tag{2}
\]

It has a purely imaginary transform, of the same functional form,

\[
\tilde{\Delta}(k) \equiv \int d^4x \Delta(x)e^{ikx} = 4\pi^2i\Delta(k). \tag{3}
\]

C. Locality of quantum field states

The standard model of locality for field states is based on point localization in a classical space-time background. To quantize fields \[41\]–\[42\], the light cone function is used to write covariant commutation relations for field operators \(\hat{A}(x)\),

\[
[\hat{A}_\mu(x), \hat{A}_\nu(x')] = g_{\mu\nu}\Delta(x - x'), \tag{4}
\]

where \(g_{\mu\nu}\) denotes a (classical) tensor. The explicit geometrical coefficients on the right hand side are all classical objects: geometry is not part of the quantum system. There is an unwritten quantum operator on the right side which is just the identity operator on a field state. Since the only position dependence on the right side comes in classical functions, there is no entanglement of field states with geometry.

For field dynamics, it is necessary to include conjugate momentum or derivative operators. This leads to
propagating states in the form of plane waves, with a commutator for the transform of quantized wave modes

$$[\hat{A}_\mu(\vec{k}), \hat{A}_\nu(\vec{k}')] = ig_{\mu\nu}\delta(\vec{k} - \vec{k}')/4\pi^2 k_0,$$

where $g_{\mu\nu}$ again denotes a (classical) tensor. The coefficient depends on the spin or helicity of the field, and relates internal quantum degrees of freedom to the classical inertial frame. In the case of linearized gravity the propagating quanta are gravitons, a spin 2 tensor field with a very small self coupling set by the Planck scale.

In standard inflation, perturbations arise from linearized gravity coupled to quantum fluctuations of field vacuum states. The “collapse” of these states into a classical metric occurs coherently in $\ell_P$ space, for each spatially infinite mode. In holographic or spooky inflation, these perturbations are subdominant to those from new fundamental geometrical quantum degrees of freedom that underlie holographic, emergent gravity[17, 18].

D. Model of coherent causal diamonds

The new quantum-geometrical degrees of freedom can be modeled with a different quantum system, based on geometrical position operators $\hat{P}_\kappa$. These quantities represent nonlocally quantized event positions, labeled by expectation values of position $x_\kappa$:

$$\langle \hat{P}(x_\kappa) \rangle \equiv \langle \hat{P}_\kappa \rangle = x_\kappa.$$

A contraction with the antisymmetric Levi-Civita 4-tensor $\epsilon_{\kappa\lambda\mu\nu}$ allows us to write a Lorentz covariant generalized rotational commutative algebra with the same light cone structure as the field commutator (Eq. 4):

$$[\hat{P}_\kappa, \hat{P}_\lambda] = i\epsilon_{\kappa\lambda\mu\nu} \hat{P}^\mu \Delta(x') \ell_P.$$

The imaginary coefficient in Eq. 7 allows for superposition and entanglement of the geometrical states. In Eq. 7, geometrical quantum operators on both sides share the same degrees of freedom, so there are new nonlocal quantum relations among spatial directions, and between space and time, that cannot be described by Eq. 4, even with a linearized tensor field to represent quantum gravity.

Eq. 7 is not a fully consistent noncommutative quantum geometry, since the light cone function $\Delta(x')$ is not a quantum operator and a classical metric has been used for raising and lowering the indices. It projects the geometrical state onto a classical metric, and the operator labels onto a classical inertial frame. In the physical interpretation elaborated below, the classical metric corresponds to a measurement, and thereby a choice of observer world line and inertial frame.

The localization scale $\ell_P$ has the same dimensions as $P$. It fixes the information content of the new degrees of freedom in physical units. In our physical interpretation of this system, $\Delta(x')$ is a function of physical space-time event positions, so $\ell_P$ represents the coherence scale of light cone states in the frame of an emergent observer—a finite resolution in proper time. As shown below, for a system that obeys the holographic principle, such as emergent gravity, $\ell_P$ is the Planck length.

1. Nonlocal measurement, projection and uncertainty

Even though the degrees of freedom represented by the $\hat{P}_\kappa$’s have no local or dynamical effects, their fluctuations in time and direction affect correlations in nonlocal measurements.

Consider projection onto a 3D spacelike surface of constant $x^0 \neq 0$. The light cone function $\Delta(x^0)$ is then a $\delta$-function on a 2-sphere of radius $|x_0|$, which coincides with the surface of a causal diamond. For spatial positions on this surface, we can set $\nu = 0$ in Eq. 7 to obtain a standard spin algebra in three dimensions, with relabeled indices $i, j, k$ taking values $1, 2, 3$,

$$[\hat{P}_i, \hat{P}_j] = i\epsilon_{ijk} \hat{P}_k \ell_P.$$

Thus, when a light cone function $\Delta(x^0)$ is used to project the four dimensional space-time operators onto three dimensions — an eigenstate of proper time— it creates a fully consistent quantum algebra that entangles positions in three dimensions.

We now recall some standard results of quantum mechanical spin in three dimensions (e.g., [43, 44]). There are rotations about each spatial axis,

$$[\hat{P}_1, \hat{P}_2] = i\hat{P}_3 \ell_P$$

and its cyclic permutations. A radius operator $\hat{P}^2 = \sum \hat{P}_i^2$ commutes with all direction components:

$$[\hat{P}_i, \hat{P}^2] = 0.$$

The discrete states of the system are assembled using raising and lowering operators. For direction 3,

$$\hat{P}_{3\pm} = \hat{P}_3 \pm i\hat{P}_2$$

with the commuting properties

$$[\hat{P}_3, \hat{P}_{3\pm}] = \pm \ell_P \hat{P}_{3\pm},$$

raising and lowering operators for the other directions again obtained by cyclical permutations of $1, 2, 3$. In a conventional integer-spin representation, $l = 0, 1, 3, \ldots$ denotes the principal quantum number, and $m = -l \ldots + l$ denotes projections onto a chosen axis, say 3. The eigenvalues of $\hat{P}^2$ are $l(l+1)\ell_P^2$, and the eigenvalues of $\hat{P}_3$ are $m\ell_P$.

For each $l$ there are $2l+1$ directional projection eigenstates, so the number of degrees of freedom $N$ scales holographically, as the surface area in Planck units:

$$N = \sum_{\ell = 0}^{l} (2l + 1) \approx (|P| / \ell_P)^2,$$
where the approximation applies in the large $l$ limit. Thus, a Planck scale normalization $(\ell_p \approx ct_\nu)$ agrees with holographic emergent gravity\cite{24,26}. We interpret $N$ as the geometrical information in the whole (3+1D) volume enclosed by a (2+1D) causal diamond defined by a 1D proper time interval.

Using standard algebraic methods\cite{18,43}, it can be shown that in an eigenstate of $\hat{P}_3$,

$$\langle \hat{P}_1 \rangle = \langle \hat{P}_2 \rangle = 0 \quad \text{(14)}$$

and

$$\langle \hat{P}_1^2 \rangle = \langle \hat{P}_2^2 \rangle = (\hat{P}^2 - \hat{P}_3^2)/2 = \ell_p^2/((l+1)-m^2)/2. \quad \text{(15)}$$

As a result, we can write a generalized uncertainty principle for quantum fluctuations in any three orthogonal directions:

$$\langle \delta P_i^2 \rangle = \langle \delta \hat{P}_i^2 \rangle + \langle \delta \hat{P}_2^2 \rangle + \langle \delta \hat{P}_3^2 \rangle \geq l \ell_p^2 > |\mathcal{P}|\ell_p, \quad \text{(16)}$$

where

$$\delta \hat{P}_i^2 \equiv \langle \hat{P}_i^2 \rangle - \langle \hat{P}_i \rangle^2. \quad \text{(17)}$$

For large $l$, the uncertainty (Eq. \ref{16}) is much larger than $\ell_p^2$. An eigenstate of time—such as a causal diamond in flat space-time—has a 2D bounding surface radius with radial quantum fluctuations of this magnitude, entangled among all three directions. A similar magnitude of macroscopic geometrical uncertainty is estimated in other models of holographic space-time\cite{17}.

2. Coherent macroscopic correlations

The quantum system in Eq. \ref{9} represents transverse degrees of freedom in time and space unlike any classical dynamical degrees of freedom: departures from classical locality coherently shared between time and space everywhere on a causal diamond. Their macroscopic fluctuations display new properties: they are much larger than those produced by Planck scale quantum field fluctuations, and they are directionally coherent.

As shown above (Eq. \ref{16}), a lower limit on fluctuation variance can be expressed as an uncertainty principle,

$$\langle \delta P_i^2 \rangle \geq \mathcal{P}\ell_p, \quad \text{(18)}$$

which increases with the magnitude of $\mathcal{P}$, the separation between points. The amplitude of the displacement increases with scale, but the fractional effect on emergent relative positions gets smaller:

$$\langle \delta P_i^2 / \mathcal{P}^2 \rangle \geq \ell_p / \mathcal{P} \quad \text{(19)}$$

so space-time relationships on large scales approximate a classical metric with small coherent fluctuations. The model of inflationary correlations developed here is based on the hypothesis that correlations in the metric on all scales preserve some symmetries of these coherent causal diamonds in flat space-time.

III. PERTURBATIONS ON THE INFLATIONARY HORIZON

We now extrapolate from a flat space-time background to model symmetries of perturbations that originate on coherent quantum horizons during slow-roll inflation.

A. Causal structure of inflationary space-time

An unperturbed inflationary universe has a Friedmann-Lemaitre-Robertson-Walker metric, with space-time interval

$$ds^2 = a^2(t)[c^2 d\eta^2 - d\Sigma^2], \quad \text{(20)}$$

where $t$ denotes proper cosmic time for any comoving observer, $d\eta \equiv dt/a(t)$ denotes a conformal time interval, and $a(t)$ denotes the cosmic scale factor, determined by the equations of motion. The spatial 3-metric in comoving coordinates is

$$d\Sigma^2 = dr^2 + r^2 d\Omega^2, \quad \text{(21)}$$

where the angular interval in standard polar notation is $d\Omega^2 = d\theta^2 + \sin^2 \theta d\phi^2$. Future and past light cones from an event are defined by a null path,

$$d\Sigma = \pm cd\eta. \quad \text{(22)}$$

Causal diagrams for an inflationary metric are shown in Figures \ref{1} and \ref{2}). The end of inflation $t_f$ is defined by when the expansion changes from accelerating $\ddot{a} > 0$ to decelerating $\ddot{a} < 0$. A causal diamond for an observer $O$ with boundary at $t_1$ corresponds to an interval with equal conformal time before and after $t_f$.

The inflationary horizon $H$ is an inbound null surface that arrives at an observer $O$ at the end of inflation. The exact choice of null surface (and $t_1$) does not matter; the important thing is that $H$ forms the future boundary of a series of causal diamonds of nearly constant area $4\pi(c/H)^2$ during the slow-roll phase. We will posit that exotic geometrical correlations on the light cone $H$ imprint correlations on the emergent curvature of comoving world lines as they pass through $H$.

B. Generation of relic classical perturbations

The covariant model of geometrical nonlocality developed above is based on position operators $\mathcal{P}$ referred to a flat metric. The premise of spooky inflation is that the nonlocal uncertainty of position required by nonlocality on causal diamond surfaces appears in the form of scalar curvature perturbations on the horizon of the emergent inflationary metric.

Adopting standard conventions for linear perturbations\cite{13,16}, denote the invariant curvature perturbation\cite{15} in comoving 3-space by $A(\vec{r})$, and
FIG. 1. Causal structure of an inflationary universe, with two spatial dimensions suppressed. Vertical axis represents the world line of an observer $O$, horizontal axis represents the end of inflation $t_I$, and left and right halves represent antipodal spatial directions. The horizon $H$ forms the outer boundary of causal diamonds that end on $O$ before $t_I$. Dashed lines represent spatial hypersurfaces of constant cosmic time, and dotted lines are world lines of constant comoving position. In standard inflation, a plane wave mode freezes out acausally along a dashed line when its wavelength matches the horizon, so quantum states of geometry collapse into eigenstates of wave modes on 3D spacelike hypersurfaces of constant $x_0$; in spooky inflation, geometrical states collapse on 2D boundaries of causal diamonds defined by the horizon.

its spectral transform in comoving wavenumber space $\vec{k}$ on surfaces of constant cosmological time by

$$\hat{A}(\vec{k}) = \int d\vec{r} A(\vec{r}) e^{i\vec{k} \cdot \vec{r}} = |\hat{A}(\vec{k})| e^{i\theta(\vec{k})}. \quad (23)$$

The following analysis concentrates on directional correlations, so perturbations will be described in the polar coordinates adopted for the metric, $A(r, \theta, \phi)$. Usually, quantum coherence is assigned and matched to the gravitational effect of plane-wave modes of amplitude $\hat{A}(\vec{k})$ on infinite spacelike hypersurfaces (Fig. 1).

In the spooky model, the relic curvature perturbations are matched to fluctuations of the quantum system projected onto spherical causal diamond boundaries of the inflationary horizon $H$ (Fig. 2). Inside $H$, the system is in a coherent state. The state of each causal diamond nested in $H$ determines the potential of $O$ in relation to events on its surface. Inflationary fluctuations freeze in as differences of potential $A$ from the observer when a comoving separation $r$ passes through the horizon. A classical solution with conserved local curvature in comoving coordinates is assumed outside $H$ for any observer.

From Eq. (19), the fractional fluctuation power of dimensionless perturbations on a horizon of radius $H = c/H$ is given by

$$\langle A^2 \rangle = \langle \delta P^2 \rangle / P^2 = H t_p, \quad (24)$$

where observed perturbations $A^2 \sim 10^{-9}$. Perturbations that scale like Eq. (24) with slowly varying $H$ typically produce a nearly-scale-invariant power spectrum indistinguishable from standard cosmology, and agree with spectral measurements of CMB anisotropy and cosmic structure for a suitable choice of inflation model. Unlike the standard model based on coherent quantum field wave modes, the coherent horizon has correlations among different spatial directions, and over a broad band of $k$.

C. Candidate symmetries

We seek to test the hypothesis that the pattern of relic curvature $A(r, \theta, \phi)$ is a coherent projection of the quantum light-cone state when the comoving sphere $r$ leaves $H$. Our approach is to identify new candidate directional symmetries, motivated by models of black hole horizons and emergent nonlocality on light cones, that cannot be produced in the standard scenario.
1. Constant variance on great circles

One set of candidate directional symmetries is based on the idea that emergent scalar curvature perturbations in different directions have nonlocal relationships determined by coherent states of a quantum system with rotational symmetry. In Eqs. [7] and [8], the commutator on the left hand side that is responsible for the exotic quantum fluctuation has projected space and time components orthogonal to the components on the right hand side, related by the antisymmetric tensors $\epsilon_{\lambda \mu}$ and $\epsilon_{ijk}$, and leads to an uncertainty (Eq. [10]): an eigenstate with $\delta P_3 = 0$ has an irreducible uncertainty in the sum of orthogonal components, $\langle \delta P_1^2 + \delta P_2^2 \rangle$.

A similar symmetry applied to curvature during inflation would relate perturbation power in orthogonal directions. A possible consequence is that a directional average in the plane normal to any direction, say 3, obeys the same relation:

$$\langle \mathcal{A}^2 \rangle_{\perp 3} = \langle \delta P_1^2 + \delta P_2^2 \rangle / P^2 \geq \langle \mathcal{A}^2 \rangle.$$  \hspace{1cm} (25)

Suppose that cosmology preserves the symmetry of statistical isotropy, that is, statistical quantities are independent of direction. In order for Eq. (25) to hold for any direction in a single statistically isotropic distribution, the bound must saturate in all directions: if any direction were to exceed the overall average, another direction would need to have less than the lower bound. This leads to a symmetry of the azimuthal average:

$$\langle \mathcal{A}^2 \rangle_{\phi} \equiv \int d\phi' [A(\phi', \theta' = \pi/2)]^2 = \langle \mathcal{A}^2 \rangle,$$  \hspace{1cm} (26)

for any orientation of polar coordinates $(\theta', \phi')$. That is, the variance of curvature perturbations on any great circle is equal to the variance for the whole sky.

The symmetry represented by Eq. (26) does not hold for standard random-phase noise. It arises from the rotational symmetry of the emergent system, and the holographic absence of one independent rotational degree of freedom. As shown below, it has distinctive signatures that can differentiate the distribution from the random gaussian noise predicted in the standard model.

As usual, there is a zero mean variation averaged over all directions $\theta$,

$$\langle A \rangle_{\theta} = \int d\phi d\theta \sin(\theta) A(\phi, \theta) = 0,$$  \hspace{1cm} (27)

so $A$ can be decomposed into a linear sum of spherical harmonic components $Y^m_\ell$. To satisfy the symmetry of great circle variance (Eq. 26) they must have a coherent relationship that appears as a conspiracy of alignments and amplitudes. It is useful to illustrate with an example that uses a few low order spherical harmonics: octopole ($\ell = 3$), quadrupole ($\ell = 2$), and dipole ($\ell = 1$).

Suppose there is a dipole aligned along the $z$ axis, with $A_1 \propto \cos \theta$. It represents the intrinsic dipole of curvature in the polar direction, as viewed in the local cosmic rest frame. By itself, it obeys Eq. [26] for all great circles that pass through the pole, with normal directions in the $x, y$ plane.

For the dipole alone, the variance on the equator $\theta = \pi/2$ vanishes so it does not satisfy Eq. (26). The other multipole moments must organize around this direction to have perturbations satisfying Eq. (26) in all directions. The fit improves if we add to the dipole a precisely aligned sectoral octopole ($\ell = m = 3$), with angular dependence $A_3 \propto e^{3\phi} \sin^3 \theta$, and a precisely aligned sectoral quadrupole ($\ell = m = 2$), with angular dependence $A_2 \propto e^{2\phi} \sin^2 \theta$. The dipole variance vanishes along the equator, while the octopole and quadrupole variances are maximized for an equatorial great circle. To create a pattern consistent with Eq. (26) in all directions requires correlations and alignments of higher multipoles, although the symmetry could apply over only a limited range of $\ell$, depending on the angular coherence scale of correlations. Some manifestations of multipole alignment in the CMB are outlined below.

2. Antipodal anticorrelation and parity violation

Another candidate feature of spooky inflationary perturbations is antipodal correlation. In our model of quantum nonlocality, antipodal point-parity antisymmetry appears directly in the Dirac light cone structure (Eq. [2]) and carries over into positional operators:

$$\tilde{\mathcal{P}}(x) = -\mathcal{P}(-x).$$  \hspace{1cm} (28)

The nonlocal information represented by antipodal antisymmetry helps to resolve information paradoxes in quantum models of eternal black holes[19–21], where antipodes on the horizon are actually identified, and time-reversed conjugate particle states are entangled at opposite spatial poles, even on macroscopic scales. In our model of emergent inflation, the antipodal outgoing states of geometry on the horizon become curvature perturbations of the emergent metric, and can have similar nonlocal spooky correlations.

The spooky model creates a possible mechanism to produce a fundamental asymmetry between even and odd parity relationships between opposite points in the sky. In standard inflation, the underlying physics is symmetric, due to statistical homogeneity and isotropy of the background. Any measured asymmetry in antipodal perturbation power is entirely due to “cosmic variance” from the zero expected value of the ensemble. In the range of possible realizations, the bulk of the point-parity antisymmetry (or symmetry) is usually contributed by a small number of harmonic modes on large angular scales, especially the intrinsic dipole and quadrupole.

In the spooky model, there can be a universal physics-based parity violation, and global symmetries governing antipodal points. Indeed there may be little or no cosmic variance in the ratio of odd to even perturbation power;
its value may be set by the underlying physics. Gener-
ically, it is plausible to have a significant imbalance be-
tween odd and even power, given the parity violation of
known matter fields, and the fact that the slow-roll in-
flationary background cosmology violates temporal sym-
metry. In general, the spectrum of the asymmetry ex-
tends to higher $\ell$ than the standard picture: violation is
not confined to low-$\ell$ modes. We will explore signatures
of antipodal anticorrelation, or a bias of odd over even
power, characterized by a parity violating parameter as
a function of angular scale.

To start with, consider the extreme case: frozen metric
perturbations with exact antisymmetry similar to eternal
black holes, so that the curvature in direction $\theta$ satisfies
\begin{equation}
\mathcal{A}(\theta) = -\mathcal{A}(-\theta),
\end{equation}
an exactly odd point parity of perturbations. Microscopic
spookiness is manifested macroscopically: points in op-
site directions “know about each other” like nearby ones
do. The quantum weirdness does not separate scales: it
applies to fine-grain angular detail (that is, high resolu-
tion $\ell$), even at large angular separation.

In our model, antipodal anticorrelations on the infla-
tionary horizon lead directly to antipodal perturbation
power. In this extreme example, the all-sky distribution
has an antipodal variance equal and opposite to the
single-point variance:
\begin{equation}
\langle \mathcal{A}(\tilde{\theta})\mathcal{A}(-\tilde{\theta}) \rangle = -\langle \mathcal{A}^2 \rangle.
\end{equation}

We should allow for the possibility of a less extreme
imbalance of odd and even perturbations. In cosmol-
yogy, as in a realistic time-asymmetric black hole, an ex-
act antipodal antisymmetric black hole can be broken by a back-
ground system that is not in a time-symmetric equilib-
rium or ground state. The exact antipodal antisymmetry
of flat-space-time may not apply for inflationary per-
turbations, since the classical inflationary background
breaks the time-displacement and boost symmetries of
classical relativity, and the vacuum matter fields that
couple to gravity also in general violate parity. The mag-
nitude of time-direction symmetry breaking is related to
how much the inflationary solution departs from that of
the maximally symmetric de Sitter inflationary solution.
One direct measure of time asymmetry comes from the
slight tilt of the power spectrum from exact scale invari-
ance, measured\cite{10,11} to be $1 - n_S = 0.035 \pm 0.004$, which
arises from the small fractional decrease in the expan-
sion rate during each $e$-folding of inflation\cite{18}. We adopt
a cosmological symmetry breaking parameter $\mathcal{E} < 1$, so
that Eq. (30) becomes
\begin{equation}
\langle \mathcal{A}(\tilde{\theta})\mathcal{A}(-\tilde{\theta}) \rangle = -\langle \mathcal{A}^2 \rangle (1 - \mathcal{E}),
\end{equation}
where the ratio of even to odd perturbation power is $\mathcal{E}$. In
general, $\mathcal{E}$ depends on angular wave number $\ell$.

3. Global azimuthal symmetries

The same orthogonal commutator structure (Eq. 8),
that may give rise to constant great circle curvature vari-
ance (Eq. 26), could also produce an exact global equa-
torial symmetry:
\begin{equation}
\langle \mathcal{A}(\tilde{\theta})\mathcal{A}_{\perp\tilde{\theta}} \rangle = 0.
\end{equation}
where $\langle \rangle_{\perp\tilde{\theta}}$ denotes the azimuthal mean on the great cir-
cle normal to direction $\tilde{\theta}$, and $\langle \rangle_{\tilde{\theta}}$ denotes an average
over all directions $\tilde{\theta}$. This equatorial symmetry follows
from exactly odd point parity ($\mathcal{E} = 0$), or if all great
circles have vanishing mean. However, it does not re-
quire either of these to hold: it specifies an orthogonal
rather than an antipodal relationship, and can hold ex-
actly even if antisymmetry does not ($\mathcal{E} \neq 0$). It could
arise from independently-generated perturbations in or-
thogonal directions from any observer.

Since by definition the point average vanishes,
\begin{equation}
\langle \mathcal{A}(\tilde{\theta}) \rangle_{\tilde{\theta}} = 0,
\end{equation}
Eq. (32) follows if $\langle \mathcal{A} \rangle_{\perp\tilde{\theta}}$ on each great circle is uncor-
related with the sum of its polar values, $\mathcal{A}(\tilde{\theta}) + \mathcal{A}(-\tilde{\theta})$.
This simple symmetry could be a natural consequence in
a cosmology with an observer-dependent emergent back-
ground, in which even-parity perturbations of one ob-
server are absorbed in the mean curvature of another\cite{18}. In
standard cosmology, a coherent plane wave perturba-
tion mode on a pre-existing background generally affects
all parts of the sky, and a global conspiracy of mode
phases is required for even-parity modes to cancel in this
way.

This equatorial symmetry is a special case of more gen-
eral (that is, more restrictive) symmetries that could ap-
ply for circles on other planar slices normal to a polar
direction. A more general formulation is that a global
directional average of azimuthal averages at a polar an-
gle $\Theta$ has vanishing correlation:
\begin{equation}
\langle \mathcal{A}(\tilde{\theta})\mathcal{A}_{\Theta,\tilde{\theta}} \rangle = 0,
\end{equation}
where $\langle \rangle_{\Theta,\tilde{\theta}}$ denotes an azimuthal mean on a circle at
a constant polar angle $\Theta$ about direction $\tilde{\theta}$. Depending
on the coherence of the emergent light cones, the inter-
sections of causal diamond surfaces could lead to special
angles of frozen correlations at $\Theta = 30$ and 60 degrees,
which could eventually appear as characteristic scales in
the temperature correlation function (see Fig. 3).

D. Symmetries of the correlation function

A rigorous derivation of the holographic origin of these
symmetries requires development of physical models of
emergent inflation based on coherent causal diamonds
applies even at high angular resolution in a spooky model with the exact symmetry given by Eq. \([32]\); it can occur by chance in some realizations of a conventional cosmology, but only very rarely to a very high precision. More general symmetries of azimuthal averages (Eq. \([34]\), Fig. \([3]\)) could in principle lead to null predictions for \(C_A(\Theta)\) at other particular angles.

Contributions of odd and even spherical harmonics to the correlation function are given by the standard formula (e.g., \([4, 46, 47]\))

\[
C_A(\Theta) = \frac{1}{4\pi} \sum_{\ell} (2\ell + 1) C_\ell P_\ell(\cos \Theta),
\]

where \(P_\ell\) are the Legendre polynomials; \(P_0(0) = 0\) for odd \(\ell\), \(P_0(0) \neq 0\) for even \(\ell\). With only odd harmonics in Eq. \([37]\), the odd parity of odd Legendre polynomials leads to an exact reflection antisymmetry of \(C_A\) around \(\pi/2\):

\[
C_A(\pi/2 + \delta\Theta) = -C_A(\pi/2 - \delta\Theta).
\]

This property is violated by even modes, but only in second order in \(\delta\Theta\). From Eq. \([37]\), the derivative of every even-parity contribution in \(C_A\) vanishes at \(\Theta = \pi/2\), so departures from Eq. \([38]\) are predicted to be of order \(|\mathcal{E}|(|\mathcal{E} + 1)|\delta\Theta^2 C_A(0)|/2\), where \(|\mathcal{E}|(|\mathcal{E} + 1)|\) is the fraction of total fluctuation power in even modes. One simple consequence is that the mean of \(C_A\) in a band of width \(\pm\delta\Theta\) around \(\pi/2\) vanishes to a precision of order \(|\mathcal{E}|(|\mathcal{E} + 1)|\delta\Theta^2 C_A(0)|/2\).

The values of \(C_A\) at zero lag, \(C_A(0) \equiv C_0\), and at 180 degrees, \(C_\pi\), are directly related to the asymmetry in perturbation power parity. With parity violation, Eq. \([33]\) leads to an antipodal anticorrelation:

\[
C_A(\pi) = -(1 - \mathcal{E}) C_A(0).
\]

Thus, a measurable ratio of total perturbation power to antipodal power,

\[
R_{0\pi} = \frac{C_0}{C_\pi} = \frac{1}{1 - \mathcal{E}},
\]

directly related to the ratio of even to odd perturbation power,

\[
\mathcal{E} = \frac{R_{0\pi} + 1}{R_{0\pi} - 1} = \frac{C_0 + C_\pi}{C_0 - C_\pi}.
\]

If even and odd perturbations have the same total power, \(C_\pi = 0\), the expected value in the standard scenario. In standard cosmology, antipodal perturbation power mostly arises from large angle, large scale modes; the large number of high-\(\ell\) modes average to contribute close to the expected zero net parity in a typical realization. (In most ensemble predictions of \(C(\Theta)\), the true expected average is obscured by the subtraction of the unobserved dipole, which also contributes much of the cosmic variance in predictions at large \(\Theta\).) In a spooky cosmology, antipodal correlation power can also manifest at higher map resolution \(\ell\), that is, smaller scale structure also tends to be anticorrelated. This property is useful for comparisons with standard cosmology.
IV. PATTERNS IN CMB ANISOTROPY

On large angular scales, CMB temperature anisotropy is dominated by perturbations within a relatively thin sphere at the epoch of last scattering, with some secondary effects from the intervening volume[48–50]. At low $\ell$ the primary temperature anisotropy is a direct map of the primordial potential at last scattering on each scale, in the sense that there is an approximately direction-independent transfer function between the scalar curvature perturbation and observed temperature — the Sachs-Wolfe approximation[48]. On these scales, primordial directional correlations of $A$ on a sphere then apply approximately to $\delta T/T$.

A. Interpretation of well known CMB anomalies

In the Sachs-Wolfe approximation, the spooky model provides a unified physical interpretation that accounts in a general way for several long-studied features of measured CMB anisotropy at low $\ell$ that are statistically anomalous in the standard model[4, 8, 13, 46, 51–53]. We first summarize a proposed interpretation of some well known empirical anomalies, then suggest more sharply defined theory-motivated tests that can differentiate spooky inflation from the standard picture.

1. Axes defined by the quadrupole and octopole are closely aligned. The WMAP all-sky maps[2–4] revealed a remarkably close agreement in direction for quadrupole ($\ell = 2$) and octopole ($\ell = 3$) harmonics. The aligned direction is defined by the axis that maximizes the sum of the squares of $a_{\ell,0}$ and $a_{\ell,-\ell}$ spherical harmonic coefficients, that is, maximizes polar asymmetry. A variety of studies have confirmed the close alignment to be highly unlikely in the standard model[4]. In our model, as discussed above, the principal axes of harmonics could be aligned to satisfy the constraints imposed by a holographic information deficit with rotational symmetry: either constant variance on great circles (Eq. 26), or vanishing global azimuthal equatorial averages (Eq. 32).

Since the principal axis is associated with a physical primordial dipolar curvature mode on the horizon scale, whose orientation correlates with the principal axes of other large scale modes, the model accounts for why “secondary” Integrated Sachs-Wolfe (ISW[48–50]) contributions from gravitational effects in the intervening volume do not spoil the precise alignment.

2. The shape of the octopole is unusually “planar”[40]. In our model, the primordial large scale directional pattern arises from a collapse of a quantum state with a principal axis, that of the intrinsic dipole, along with correlated higher multipoles. Directional components of multipoles are not random with respect to the dipolar axis, but reflect the coherence of the quantum state, like an atomic orbital eigenmode. In the illustrative example above, as part of this pattern, the octopole takes a sectoral form dominated by $|m| = \ell = 3$ specifically to compensate for the vanishing equatorial variance of the dipole, to match the constraint of constant great-circle variance (Eq. 26). A test of this interpretation is outlined below.

3. The two-point temperature correlation function is small at large angular separation[8, 13, 46]. An unexpected lack of large angle correlation power has been apparent since the first measurements with COBE[47]. In the WMAP analysis of $C(\Theta)$, based on 7 years of data[4], the authors comment on the (true) fact that there is no significant conflict with the standard random-phase scenario, and no significant deficit of large-scale power: $C(\Theta)$ lies within the 95% confidence range of the best-fit ΛCDM model for all $\Theta$, as determined by Monte Carlo simulations. This supports the conclusion that there is no statistically significant lack of large-scale power on the full sky.” At the same time, compared with a standard ensemble of random-phase CMB realizations, $C(\Theta)$ is both anomalously close to zero around 90 degrees, and anomalously negative near 180 degrees.

By contrast, the WMAP measurement appears to agree remarkably well with the simple and exact global equatorial symmetry (Eq. 30) of spooky inflation: the measured $C_\ell(\Theta)$ has a value at 90 degrees very close to zero, and a nearly antisymmetric dependence for a significant range around that (Eq. 38).

4. The quadrupole and other even harmonics are smaller than expected. An excess of odd over even fluctuation power[8, 13, 46], measured in harmonic decomposition, shows anomalous antipodal anticorrelation on angular scales smaller than the dipole, which also appears as a significant negative correlation in $C_\ell(\Theta)$ near 180 degrees. Both are interpreted here as direct signatures of antipodal anticorrelation on the horizon, with $E < 1$ over a wide range of $\ell$ (Eq. 31).

B. Tests and comparisons with current data

Because there is no cosmic variance in an exact symmetry, precision tests are possible on large angular scales without the usual cosmic-variance penalty on significance: predictions for symmetries are not influenced by the random variables of realizations as they are in the standard scenario. This feature allows for powerful parameter-free comparisons with standard theory.

1. Equatorial antisymmetry

Again, consider a low pass filter $\ell < 20$, and use the Sachs-Wolfe approximation. The exact prediction for curvature correlation (Eq. 36) translates into a vanishing temperature anisotropy correlation $C_{TT}$ at 90 degrees,

$$C_{TT}(\pi/2) = 0,$$  \hspace{1cm} (42)

with an antisymmetric shape given by Eq. 38,

$$C_{TT}(\pi/2 + \delta \Theta) = -C_{TT}(\pi/2 - \delta \Theta).$$  \hspace{1cm} (43)
One example of a robust null prediction is that the mean of $C_{T\ell}(\Theta)$ vanishes in a band $\pm \delta\Theta$ around 90 degrees, to precision $|E/(E+1)| \delta^2 C_{T\ell}(0)/2$.

This prediction does not depend on the intrinsic dipole, which is unmeasured in CMB maps. It is also insensitive to significant extra sources of fluctuations not included in the Sachs-Wolfe approximation, from velocities of the primordial plasma; these become important only at small angular separations, $\Theta < 1$. The main obstacles to a robust null test appear to be on the measurement side, particularly from well known difficulties in foreground subtraction. They are more important here than in large scale tests of the standard picture, whose predictions have a large scatter from cosmic variance.

2. Antipodal anticorrelation

The antipodal anticorrelation of curvature (Eq. 39) leads to

$$C_{T\ell}(\pi) = -(1 - \mathcal{E}_{T\ell}) C_{T\ell}(0),$$

where $\mathcal{E}_{T\ell}$ denotes a parity violation parameter for temperature at resolution $\ell$. The latter is guaranteed to be larger (that is, closer to unity and less antisymmetric) than the parameter $E$ for curvature, because of the unmeasured intrinsic dipole, whose restoration adds to $C_{T\ell}(0)$ and subtracts from $C_{T\ell}(\pi)$ by the same amount. Nevertheless, if $\mathcal{E} < 1$ then antipodal anticorrelation should still produce $\mathcal{E}_{T\ell} < 1$ in the higher ($\ell > 1$) harmonics, observable even after removal of the intrinsic dipole, so the measured $C_{T\ell}(\pi)$ is expected to be negative, and not very close to zero. This contrasts with the standard expectation in which antipodal power at high $\ell$ averages to zero, the bulk of the antipodal power in any realization is dominated by the cosmic variance of low order harmonics, and the measured value is expected to be positive because of the unmeasured dipole.

This prediction on its own does not allow for a highly precise null test, but can add to significance in a comparative test of theories. From Eq. (37), restoration of the unobserved dipole will add a contribution to $C_{T\ell}(\Theta)$ proportional to

$$C_{\text{dipole}}(\Theta) \propto \cos(\Theta),$$

which will reduce the value of $\mathcal{E}_T$ (that is, decrease the fraction of even-parity power in Eq. (44), while preserving Eq. (42). It can only add to the substantial antipodal anticorrelation already measured by WMAP and Planck in finer-grained structure in the higher order multipole spectrum.

A certain amount of fine-grained antipodal anticorrelation appears to be real. One measurement is the parity asymmetry in the power spectrum, which can be defined as a function of filtering scale $\ell$. A spectral parity asymmetry has been measured to high significance for harmonics up to $\ell > 30$. In the Planck measurement (Fig. 25 of ref. [13], where our $\mathcal{E}_{T\ell}$ is shown as $R_{TT}$), a minimum value $\mathcal{E}_{T\ell} \approx 0.3$ occurs at $\ell$ of order a few, and the power asymmetry $1 - \mathcal{E}_{T\ell}$ appears to fall off approximately exponentially at higher wavenumbers, with an $e$-folding or coherence scale of roughly $\ell \approx 30$. The probability of obtaining such a large asymmetry in the standard picture reaches a minimum of about 0.2 percent at around $\ell \approx 23$. This estimate of perturbation power asymmetry is roughly in accord with the overall power asymmetry that produces the anticorrelation in $C_{T\ell}(\Theta)$ at large angles, approaching 180 degrees, in the WMAP measurement.

In the spooky picture, the value and spectrum of the asymmetry is not due to cosmic variance, but is a basic physical property that depends on how parity violation affects the process of emergent inflation. The spectral coherence scale could in principle be related to the spectral tilt connected with time asymmetry of the background, and/or the anomalous deficit in overall power spectrum in the range $20 < \ell < 30$, but we do not have a theory of these relationships.

3. Great circle variance

We do not measure directly the intrinsic dipole, but a constant great circle variance allows us to uniquely reconstruct both its axis and its amplitude from measured maps. This symmetry predicts that after restoration of the correct dipole to the map of $\delta T$, the pattern should agree with constant variance (Eq. [20]) on any great circle:

$$V_{T\perp}(\bar{\theta}) \equiv \int d\phi_\perp \delta T(\phi_\perp, \theta_\perp = \pi/2)^2 = \langle \delta T^2 \rangle_{\text{sky}},$$

where the coordinates in the integral have polar axis $\bar{\theta}$, so that $V_{T\perp}(\bar{\theta})$ denotes the temperature variance for the great circle with the axis $\bar{\theta}$. The standard picture predicts variations in $V_{T\perp}(\bar{\theta})$ for different directions, due to cosmic variance.

One approach to test this property can be described in terms of great-circle-variance sky maps: for each direction $\bar{\theta}$, calculate the variance for its normal great circle, $V_{T\perp}(\bar{\theta})$. If Eq. (46) holds for the true sky, the map for the measured sky is missing the pattern of variance from the unmeasured intrinsic dipole:

$$V_{T\perp,\text{missing}}(\bar{\theta}) + V_{T\perp,\text{measured}}(\bar{\theta}) = V_{T\perp,\text{true}}.$$  (47)

The map of missing equatorial variance is predicted to be a unique smooth function of position on the sky that depends on only the intrinsic dipole axis and amplitude:

$$V_{T\perp,\text{missing}}(\bar{\theta}) = \text{const} \times \int_0^{2\pi} d\phi_\perp \cos^2(\theta_D(\phi_\perp)),$$  (48)

where $\phi_\perp$ denotes the axial coordinate around the $\bar{\theta}$ direction, and $\theta_D(\phi_\perp)$ denotes the polar angle at each $\phi_\perp$ with
respect to the intrinsic dipole axis. Put another way, the measured (dipole-subtracted) map should approximate a unique smooth function that depends only on an axis and an amplitude,

\[ \mathcal{V}_{TLL,\text{measured}}(\vec{\theta}) = \text{const} \times \int_0^{2\pi} d\phi \sin^2[\theta_D(\phi)] \sin^2[\theta_D(\phi)], \quad (49) \]

when filtered at resolution \( \ell \) comparable to the coherence scale where the symmetry applies. In the standard picture, there is no particular symmetry to bring this about, so the variance maps should generically have more complicated random structure.

This technique might be used to verify that a great-circle-variance symmetry indeed accounts for the apparently-conspiratorial multipole alignments in the real sky: the reconstructed dipole direction should closely align with the direction already singled out by the quadrupole and octopole axes, as well as higher \( \ell \) harmonics, revealing a new “conspiracy”. The symmetry should manifest itself in maps filtered on the coherence scale, estimated previously to be of order \( \ell \approx 30 \).

If the symmetry is real, the best fit to Eqs. (47) and (49) gives an estimated amplitude and axis for the true intrinsic dipole. In principle (if not yet in practice, e.g., ref. 54) the intrinsic dipole can be independently measured, by measuring galaxy flows on very large scales (> 200 Mpc) to account for the dipole contributed by peculiar velocity.

The restoration of the intrinsic dipole also allows an estimate of the true intrinsic \( C(\Theta) \), which could reveal new exact symmetries. For example, it could be that the equatorial antisymmetry generalizes to \( C(\Theta) = 0 \) for a broader range around 90 degrees, or at other specific angles (Fig. 3).

4. Origin of the unperturbed background

In standard inflation, there is a pre-existing uniform background space-time whose origin is considered separately from that of the perturbations. The spooky inflation model also adopts the approximation of a linearly perturbed uniform background, but the pattern of perturbations arises from a quantum system underlying horizons and light cones, or emergent null surfaces around points, rather than a quantum field vacuum.

The quantum origin of the cosmological “background” space-time, or unperturbed metric, is not treated in this approximation. In the emergent picture, locality is observer dependent, and there is not actually a determinate global background metric: on the largest scales, the system remains in an observer-dependent superposition. In a fully consistent model, the properties of the emergent global background should be governed by exact symmetries of the underlying quantum system. It is plausible that the mean curvature of the apparent background is always identically zero, or has a small nonzero dimensionless value determined by the value of the typical curvature perturbations \( \mathcal{A}^2 \), or by the time asymmetry that controls the spectral tilt.

5. Relation to Holometer experiments

Nonlocal coherent fluctuations on causal diamonds, as posited here for inflation, also affect light propagating between massive bodies in flat space on macroscopic scales. It has recently become possible to probe these Planck scale fluctuations directly by laboratory experiments using interferometers with signal measurement bandwidth comparable with or greater than their free spectral range.

If displacements on causal diamonds in the laboratory are governed by the same uncertainty (Eq. 18) as that used for spooky cosmological perturbations (Eq. 24), a differential measurement of light paths of length \( L \) displays fractional fluctuations on timescale \( L/c \) of order

\[ \langle \delta L^2 \rangle / L^2 \approx L / c. \quad (50) \]

They may be observable as cross correlations, with Planck scale power spectral density, between signals in suitably configured interferometers [55–61]. The effect on the signal depends on the directional structure of fluctuations, and the spatial structure of the light paths.

The most sensitive published measurement [62] uses a light-path configuration that constrains certain specifically quadrupolar fluctuations to less than ten percent of the amplitude in Eq. (50). New experimental configurations [61] may respond to generalized rotational fluctuations of light cones, similar to that posited here for inflationary horizons. Detectable cross spectra in time domain signals could then be related directly to correlations of inflationary curvature perturbations. A detection of cosmological spookiness could provide both motivation and design guidance for future experiments.

V. CONCLUSION

The simple geometrical symmetries proposed here are examples of holographic correlations: they apply to the entire sky on all scales, and reduce the independence of perturbations in different directions. It is remarkable that generic, holographically-motivated symmetries naturally produce a pattern that matches several already-measured anomalies of CMB anisotropy, including at least one remarkable feature not generally regarded as significant, the nearly exact vanishing of \( C_T(\Theta) \) at 90 degree angular separation.

The primordial structure of curvature perturbations also survives today in three dimensional large scale structure, but the structural pattern of the galaxy distribution caused by spooky entanglement is more subtle than the CMB pattern [18]. The candidate symmetries of spooky correlations analyzed here are predicted to appear clearly in the CMB on large scales today partly because the last
scattering surface approximates a 2-sphere, like the primordial horizon where the correlations originated. For this reason, CMB anisotropy at present provides the most direct test of the spooky-inflation hypothesis.

The directional symmetries of curvature perturbations derived here (e.g., Eqs. 26, 31, 32) allow the formulation of new predictions accessible to test at large angular scales with existing datasets. Some of them have been summarized here, using a Sachs-Wolfe approximation (e.g., Eqs. 42, 43, 44, 47). The new correlations describe entangled relationships at large separations (θ > 1 in the angular domain), but also in fine grained structural detail (ℓ >> 1 in the wavenumber domain) — a feature distinctly absent in standard inflation.

These symmetries provide a foundation to develop robust tests of whether spooky inflation fits cosmological data better than the standard picture, when phase relationships of the perturbation pattern are taken into consideration. An important new feature of spooky inflation is that its precise symmetries allow predictions for some measurements that are more definite than the standard picture. Some null predictions are subject to negligible cosmic variance, so even large-angle tests θ > 1 can in principle achieve high significance. This property motivates renewed attention to the covariances in fits and background-subtracted maps on large scales: limits on significance from measurement error, including systematics from foreground subtraction, become more important than they have been for tests of cosmic-variance-limited predictions. Improved models of anisotropy, including baryonic matter and radiation transport beyond the Sachs-Wolfe approximation used here, will be needed to quantify errors, and to maximize the useful range of angular wave number ℓ, in order to increase significance and reduce sampling errors. It is possible that changes in covariances among measured quantities from higher order correlations could lead to small changes in fits to standard cosmological parameters related to CMB measurements [63], even though predictions of standard cosmology that depend only on the shape of the power spectrum are unchanged.

Confirmation of the symmetries considered here would lend support to the hypothesis that the primordial perturbations originate from new nonlocal holographic quantum geometrical degrees of freedom. Details of the perturbation pattern could reveal specific signatures of how cosmological initial conditions arise from a quantum system.

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