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A dual description of decoherence in de Sitter space

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Abstract: Decoherence associated with super-Hubble modes in de Sitter space may have a dual description, in which it is attributed to interaction of sub-Hubble modes with an "environment" residing just inside the observer's horizon. We present a version of such description, together with some consistency checks, which it is shown to pass.

Keywords: Solitons Monopoles and Instantons, Models of Quantum Gravity, Quantum Dissipative Systems.

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1. Introduction

Properties of de Sitter (dS) spacetime are of considerable interest due to the important role that near-dS spaces play in inflationary cosmology and in the theory of accelerating universe. Of particular interest is what can be called the quantum structure of the dS space, i.e., the structure that is supposed to provide us with a microscopic explanation of the value of the entropy [1]. In the present paper, we want to explore what can be learned about that structure from looking at the behavior of a two-level system in a dS spacetime. Two-level systems are ubiquitous probes of various types of quantum environments, due to their susceptibility to *decoherence* — the decay, as a result of interactions with the environment, of the off-diagonal components of the system's 2×2 density matrix.

A simple way to construct a two-level system in de Sitter space is to consider a scalar field with a degenerate double-well potential, such that the mass scale μ of the potential is sufficiently large in comparison with the Hubble parameter H. In this case, perturbative amplification of modes does not occur, and the field within a given Hubble volume spends most of the time close to one of the minima. Any particles created as a result of transitions between the minima are rapidly blown away by the cosmological expansion.

Transitions between the minima of the potential will occur via tunneling, but they will be completely incoherent (meaning that no periodic oscillations of probability will emerge) [2]. In global coordinates, this decoherence can be attributed to the super-Hubble modes of the field. On the other hand, an individual observer in dS spacetime can access only a single Hubble volume, so it is natural to ask if he can attribute the decoherence to physics at — or just inside — his own cosmological horizon. The requirement that such an alternative, or *dual*, description exists is analogous to the "complementarity principle" proposed in various forms in refs. [3, 4].

In the dual description, a statistical system (the "environment"), which lives on a sphere just inside the observer's horizon, will detect transitions between the states of the two-level system and cause decoherence in the usual way, through system-environment interactions. Only sub-Hubble modes of the field will be involved. Since the interaction of the field with the environment will typically be of the form envisioned in the AdS/CFT correspondence [5], it makes sense to explore if, in our case, a similar correspondence can be used to learn something about the "environment". (This proposal, with the boundary located just inside the horizon, is completely unrelated to the dS/CFT correspondence discussed in ref. [6].)

We will see that in our case an AdS/CFT-like correspondence is subject to various consistency checks, including a rather stringent one, which follows from the observation [2] that decoherence in de Sitter, as measured by the decay of the overlap between different evolution histories, is very strong. We show that this consistency check is nevertheless passed, due to a conspiracy of two factors: the high temperature of the environment (a consequence of the near-horizon blueshift) and the form of the environment's spectral density, which turns out to be Ohmic.

Before we proceed to the main point in section 3, we briefly discuss, in section 2, tunneling between degenerate minima in de Sitter spacetime. Using the thin-wall approximation, we describe tunneling in the hamiltonian formalism and compute the amplitude in the WKB approximation. Comparing the results with those obtained by the instanton calculus on the euclidean four-sphere, we learn about the role of the back-reaction (which the thin-wall formalism allows us to take into account) and the nature of the turning points.

Throughout the paper we consider three spatial dimensions (plus time), but the construction also works in two.

2. Tunneling in the two-level system

Because tunneling between degenerate minima of a heavy field in de Sitter space may not be entirely familiar, we begin with a brief discussion of it here. In this discussion, we consider the extreme case, when the mass scale μ of the potential is not merely somewhat larger than the Hubble parameter, but

$$\mu \gg H \,. \tag{2.1}$$

This allows us to apply the thin-wall approximation. Our subsequent results on decoherence do not depend on the limit (2.1).

Consider a thin-wall bubble with surface tension σ nucleating in de Sitter space. The problem can be studied using the well-developed junction formalism [7]; we use conventions of ref. [8]. The metric inside the bubble is taken in the form

$$ds^2 = -(1 - H^2 r^2) dt^2 + \frac{dr^2}{1 - H^2 r^2} + r^2 d\Omega^2 \,, \tag{2.2}$$

and the metric outside in the form

$$ds^2 = -f(r') dt'^2 + \frac{dr'^2}{f(r')} + r'^2 d\Omega^2 \,, \tag{2.3}$$

with

$$f(r') = 1 - H^2 r'^2 - \frac{2kE}{r'} \,. \tag{2.4}$$

Here Ω is the solid angle, k is Newton's constant, and E is the energy of the bubble. Because we consider a strictly degenerate potential, the Hubble parameter H is the same inside and outside the bubble.

Continuity of the metric requires that at the location of the bubble wall $r = r'$. Away from the wall, the radial coordinates r and r' differ: in the direction r increases, r' decreases, and vice versa. The junction conditions lead to an equation of motion for the wall [7]. In our case, that equation can be adapted from the general form given in ref. [8] and reads

$$E = M(\dot{r}^2 + 1 - H^2 r^2)^{1/2} - k\frac{M^2}{2r} \,, \tag{2.5}$$

where $M = 4\pi\sigma r^2$ is the wall's "bare" mass, and dot denotes a derivative with respect to the wall's proper time.

Values of the turning points are obtained from $\dot{r} = 0$. Vacuum tunneling, the case of main interest to us, corresponds to $E = 0$. In this case, the turning points are at $r = 0$ (no bubble) and $r = r_c$, where

$$H^2 r_c^2 = \left(1 + 4\pi^2 k^2 \frac{\sigma^2}{H^2}\right)^{-1} \equiv \frac{1}{b^2} \,. \tag{2.6}$$

The second term in the bracket represents the back-reaction of the wall on the metric. If we were to neglect this term, we would obtain $r_c = r_H \equiv H^{-1}$, that is the wall of a critical bubble would emerge precisely at the horizon. With the back-reaction included, $r_c < r_H$.

We can now compute the amplitude of tunneling in the WKB approximation. This is most easily done using the canonical formalism, where the right-hand side of (2.5) is identified with the hamiltonian $\mathcal{H}$ of the bubble. As it stands, $\mathcal{H}$ is written inconveniently in terms of the proper time derivative, but it can be rewritten in terms of a canonical momentum.[1] We define the cosmic time $t_{\cos}$ inside the bubble, so that the metric (2.2) becomes

$$ds^2 = -dt_{\cos}^2 + \exp(2Ht_{\cos})(d\rho^2 + \rho^2 d\Omega^2) \,, \tag{2.7}$$

where ρ is the comoving radius, and express $\dot{r}$ through the wall velocity $v = dr/dt_{\cos}$:

$$\dot{r}^2 = \frac{v^2}{1 - (v - Hr)^2} \,. \tag{2.8}$$

Now it can be verified that

$$\mathcal{H} = \sqrt{p^2 + M^2} + Hrp - k\frac{M^2}{2r} \,, \tag{2.9}$$

where

$$p = M\frac{v - Hr}{[1 - (v - Hr)^2]^{1/2}} \tag{2.10}$$

[1]The procedure is quite familiar for the case of a dust shell in otherwise empty space; for a recent discussion, see ref. [9]. The presence of a cosmological constant in our case introduces only minor technical complications.

is the canonical momentum. The latter turns out to be the standard expression for a relativistic momentum, except that it appears here in terms of the relative velocity of the wall and the expanding background, $v_{\rm rel} = v - Hr$.

For a given energy eigenvalue $\mathcal{H} = E$, the momentum as a function of r reads

$$p = \frac{-Hr\left(E + k\frac{M^2}{2r}\right) \pm \left[\left(E + k\frac{M^2}{2r}\right)^2 - M^2(1 - H^2r^2)\right]^{1/2}}{1 - H^2r^2} . \tag{2.11}$$

Classically forbidden regions correspond to negative values of the expression in the square bracket. The tunneling amplitude in the WKB approximation is proportional to $\exp(-S_{\rm WKB})$, where

$$S_{\rm WKB}(E) = \int_{r_1(E)}^{r_2(E)} dr (1 - H^2r^2)^{-1} \left[M^2(1 - H^2r^2) - \left(E + k\frac{M^2}{2r}\right)^2\right]^{1/2} , \tag{2.12}$$

and r_1 and r_2 are the turning points, for which the expression in the square brackets vanishes. Recall that M depends on r: $M = 4\pi\sigma r^2$.

For $E = 0$, when $r_1 = 0$ and $r_2 = r_c$, the integral in (2.12) is readily calculated:

$$S_{\rm WKB}(0) = \frac{\pi^2\sigma}{H^3 b}\left(b - \sqrt{b^2 - 1}\right)^2 , \tag{2.13}$$

where b is defined by (2.6). Another corollary to eq. (2.12) is

$$\left.\frac{\partial S_{\rm WKB}}{\partial E}\right|_{E=0} = -\frac{1}{H}\arctan\frac{1}{\sqrt{b^2 - 1}} . \tag{2.14}$$

So, if we were to look at tunneling at nonzero temperature, when the rate is an integral over E of $\exp[-2S_{\rm WKB}(E) - \beta E]$ (with some prefactor), we would find that as long as the temperature remains smaller than $1/\beta_{\rm cr}$,

$$\beta_{\rm cr} = \frac{2}{H}\arctan\frac{1}{\sqrt{b^2 - 1}} , \tag{2.15}$$

vacuum tunneling ($E = 0$) is more important than the nearby nonvacuum contributions ($E > 0$).

We can now compare these results to those from the euclidean instanton calculus. A quantitative comparison can only be made for the limiting case of vanishing back-reaction, since this is the only case for which the instanton action has been computed in ref. [2]. In this limit, $b \to 1$, and the WKB answer (2.13) becomes

$$S_{\rm WKB}(0) = \frac{\pi^2\sigma}{H^3} . \tag{2.16}$$

In the euclidean approach, on the other hand, the instanton is a four-sphere of radius r_H with the domain wall positioned along an equator. This is a complete periodic instanton, in the terminology of ref. [2], and to obtain the amplitude of tunneling we need the action of a half of it. The required half is obtained by cutting the four-sphere with a plane perpendicular to the equator. The boundary produced by this cutting has the topology of a three-sphere and can be viewed as a result of gluing together the spatial metrics of the patches (2.2) and (2.3) along the line $r = r' = r_c$. These two patches therefore correspond to the two turning points of the periodic instanton (at euclidean times $\tau = it = 0$ and $\tau' = it' = \pi$).

The requisite euclidean action is now simply the wall tension times half of the volume of the equator (the limit (2.1) is implied):

$$S_{\mathrm{inst}} = \frac{\pi^2 \sigma}{H^3} . \tag{2.17}$$

This coincides with the WKB action (2.16) obtained in the hamiltonian approach.

In conformal coordinates ρ and $\eta = -H^{-1} \exp(-Ht_{\cos})$, with ρ and $t_{\cos}$ defined by (2.7), the surface of constant $t = t_{\mathrm{nucl}}$ corresponds to [2]

$$\eta^2 - \rho^2 = H^{-2} \exp(-2Ht_{\mathrm{nucl}}) \equiv \eta_{\mathrm{nucl}}^2 . \tag{2.18}$$

Thus, η_{nucl} can be viewed as the moment of conformal time when the *center* of the bubble ($\rho = 0$) nucleates. On the other hand, the bubble wall at $r = r_c$ corresponds to $\rho = -Hr_c\eta$ and therefore nucleates at $\eta = \eta_{\mathrm{nucl}}/\sqrt{1 - H^2 r_c^2}$, i.e., *earlier* than the center (recall that $\eta < 0$). For weak back-reaction, one can show that by the time the center nucleates, the bubble is already of a super-Hubble size.

3. A dual description of decoherence

After a bubble nucleates, it begins to expand. The standard way to follow the subsequent evolution of the bubble is to switch from the static coordinates (2.2) to some global, e.g., conformal, coordinates. In such global coordinates, decoherence between the states built near the minima of the potential — we now refer to these states as $|L\rangle$ and $|R\rangle$ — can be attributed to super-Hubble modes of the field [2]. On the other hand, in the dual description, decoherence is attributed to interaction of sub-Hubble modes with an "environment" located inside the observer's horizon, so it will be sufficient for our purposes to work with a single Hubble patch described by the metric (2.2). In this section, it will be convenient for us to analytically continue to the euclidean signature, so that $t = -i\tau$ with real τ, and measure all distances in units of $r_H = H^{-1}$. The metric becomes

$$ds_E^2 = (1 - r^2) d\tau^2 + \frac{dr^2}{1 - r^2} + r^2 d\Omega^2 . \tag{3.1}$$

Note that τ is periodic with period $\beta = 2\pi$.

If the environment is placed on a sphere of radius $r_0 < 1$, the metric on that sphere will be

$$ds_0^2 = (1 - r_0^2) d\tau^2 + r_0^2 d\Omega^2 . \tag{3.2}$$

We see that the environment can be considered as being at inverse temperature

$$T_0^{-1} = 2\pi \sqrt{1 - r_0^2} . \tag{3.3}$$

In what follows we take r_0 to be close to 1, so that T_0 is very high.

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To cause decoherence between states of a scalar field ϕ the environment has to interact with the field somehow. We assume an interaction of the form

$$S_{\rm int} = \int d\tau d\Omega \Phi(\Omega,\tau)\mathcal{O}(\Omega,\tau)\,, \tag{3.4}$$

where

$$\Phi(\Omega,\tau) = \phi(r_0,\Omega,\tau) \tag{3.5}$$

is the boundary value of the field, and $\mathcal{O}$ is some operator constructed from variables of the environment. In this section, ϕ denotes small fluctuations of the scalar field near one of its minima or, in fact, any other field with similar properties, such as a component of the metric tensor (a gravitational wave).

Note that the interaction (3.4) is of the same form as bulk-boundary interactions envisioned in the AdS/CFT correspondence [5]. In that case, known dynamics of fields in the bulk can be used to explore the properties of operators in the boundary theory. Here we apply a similar idea to obtain the two-point correlation function of the operator $\mathcal{O}$.

The linearized equation of motion for the field ϕ (assuming minimal coupling to gravity) reads

$$\frac{1}{r^2}\partial_r\left[r^2(1-r^2)\partial_r\phi\right] - \frac{\hat{L}^2\phi}{r^2} + \frac{\partial_\tau^2\phi}{1-r^2} - \mathcal{M}^2\phi = 0\,. \tag{3.6}$$

Here $\hat{L}$ is the angular momentum, and $\mathcal{M}$ is the mass of the field. Similarly to the AdS/CFT case, we require that the corresponding bilinear action matches the effective action for Φ obtained in the boundary theory:

$$S[\phi] \sim \int d\tau d\tau' d\Omega d\Omega' \Phi(\Omega,\tau)\Phi(\Omega',\tau')\langle\mathcal{O}(\Omega,\tau)\mathcal{O}(\Omega',\tau')\rangle\,. \tag{3.7}$$

In global coordinates, solutions for a scalar field in the dS background are of course well-known [10]. However, these solutions represent traveling waves, while to make use of the correspondence (3.7) we will need the standing waves. Rather than expand one set of solutions in another, it will be easier for us to find the requisite properties of the standing waves directly from eq. (3.6).

Expanding ϕ in the eigenfunctions of $\hat{L}^2$ and in $e^{in\tau}$, with $n = 0, \pm 1, \ldots$, and making the change of variables $x = r/\sqrt{1-r^2}$, we bring (3.6) to the form

$$\frac{\sqrt{1+x^2}}{x^2}\frac{\partial}{\partial x}\left(\frac{x^2}{\sqrt{1+x^2}}\partial_x\phi\right) - \frac{l(l+1)\phi}{x^2(1+x^2)} - \frac{n^2\phi}{1+x^2} - \frac{\mathcal{M}^2\phi}{(1+x^2)^2} = 0\,, \tag{3.8}$$

$l = 0, 1, \ldots$. Note that if r runs from 0 to some r_0 close to 1, then x runs from 0 to large

$$R = \frac{r_0}{\sqrt{1-r_0^2}}\,. \tag{3.9}$$

Thus, the boundary is now at $x = R \gg 1$.

At $x \ll 1$, eq. (3.8) reduces to the equation for a field with mass squared $\mathcal{M}_n^2 = \mathcal{M}^2 + n^2$ in a flat tree-dimensional space. The solution regular at $x = 0$ grows at larger x. For large $\mathcal{M}_n$ this growth is exponential: $\phi = I_{l+1/2}(\mathcal{M}_n x)/\sqrt{x}$, where I_ν is the modified Bessel function. So, in the heuristic picture of the boundary radiating various fields into the bulk, only fields with sufficiently low masses will be able to reach the observer without being exponentially suppressed. At $x \to \infty$, eq. (3.8) becomes

$$\frac{1}{x}\partial_x\left(x\partial_x\phi\right) - \frac{n^2}{x^2}\phi = 0\,. \tag{3.10}$$

Except possibly for some special values of $\mathcal{M}$ and l, the solution that was regular at $x = 0$ will continue to grow at $x \gg 1$. According to (3.10), the growth is $\phi \sim x^{|n|}$ for $n \neq 0$, and $\phi \sim \ln x$ for $n = 0$. This growth is a reflection of the coordinate singularity of the original eq. (3.6) near the horizon.

Expanding the boundary value $\Phi(\Omega, \tau)$ in $Y_{lm}(\Omega)$ and $e^{in\tau}$, we can write each component of the field ϕ at $x \gg 1$ as

$$\phi_{nlm}(x) \approx \left(\frac{x}{R}\right)^{|n|}\Phi_{nlm} \tag{3.11}$$

for $n \neq 0$, or $\phi_{0lm}(x) \approx \Phi_{0lm}\ln x/\ln R$, for $n = 0$. Substituting these expressions in the bilinear bulk action, we find that the main contribution indeed comes from the large x region, with the result that

$$S[\phi] \sim \sum_{nlm}|n||\Phi_{nlm}|^2\,. \tag{3.12}$$

In particular, the contribution from $n = 0$ is suppressed as $1/\ln R$ and therefore vanishes in the large R limit. Note that for this calculation the field was assumed to be canonically normalized, which can always be achieved by a suitable redefinition.

The correspondence (3.7) now tells us that

$$\langle\mathcal{O}_{nlm}\mathcal{O}^{\dagger}_{nlm}\rangle \sim |n|\,, \tag{3.13}$$

or in the coordinate space, for $|\tau - \tau'| \ll 2\pi$,

$$\langle\mathcal{O}(\Omega,\tau)\mathcal{O}(\Omega',\tau')\rangle \sim \delta(\Omega-\Omega')\frac{1}{(\tau-\tau')^2}\,. \tag{3.14}$$

Eq. (3.13) can be analytically continued to the lorentzian signature:

$$G^R_{lm}(\omega) \sim i\omega\,. \tag{3.15}$$

Note that this analytical continuation gives a *retarded* correlator (as we indicate by superscript R). So, the imaginary part of (3.15) is, up to a numerical factor, the environment's spectral density in the corresponding channel: $\rho_{lm}(\omega) \sim \omega$. We see that the spectral density is Ohmic.

Formally, the local in space behavior of (3.14) is a consequence of the angular pieces dropping out from the asymptotic equation (3.10). So, strictly speaking, the delta-function in (3.14) is coarse-grained: if l in eq. (3.8) becomes of order R (thus involving angular scales of order R^{-1}), eq. (3.10) will no longer apply. Accordingly, (3.14) applies only for $|\Omega - \Omega'| \gg R^{-1}$. Still, since R is large, eq. (3.14) implies that the correlation length is short.

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Because the result (3.15) was obtained in the large R limit, we need to establish how far in the low ω region it actually holds. The minimal ω corresponds to the minimal nonzero energy transfer that the environment can accept. On a sphere of radius $r_0 \approx r_H$, the minimal momentum transfer is of order $r_H^{-1} \sim H$. Applied perpendicular to a typical momentum of order $T_0 \sim R$, this momentum transfer results in the minimal energy transfer of order

$$\omega_{\min} \sim \frac{H^2}{T_0} . \tag{3.16}$$

4. Consistency checks

An immediate consistency check of the above construction is the question if a correlator like (3.15), Ohmic in time (frequency) domain and local in space, can be obtained for a reasonable environment. We notice at once that this is precisely the behavior characteristic of hydrodynamic modes of a thermal system. So, if the scalar couples to some combination of the energy and pressure (and the graviton to the traceless part of the stress tensor), the form (3.15) will emerge, with the coefficient in front of ω proportional to one of the viscosities of the environment.[2]

A somewhat unusual property of eq. (3.15) is its overall magnitude: as follows from the calculation via (3.7), the correlator does not contain any large factors. In particular, powers of R present in (3.11) cancel out those from the integration over x. This does not impose any bounds on the viscosities, unless one knows precisely the coefficient of proportionality between the operator $\mathcal{O}$ and a component of the stress tensor. However, there is no such freedom in the calculation of decoherence, and so we arrive at another consistency check.

We now have two calculations of decoherence for the two-level system considered in the preceding section. One is the calculation in the global coordinates [2], where decoherence is attributed to super-Hubble modes, and the other is the dual calculation, where it is attributed to system-environment interaction. The results should agree. This requirement is nontrivial because the calculation in global coordinates shows that the effect is very large, while the magnitude of (3.15), as we have just seen, is quite modest.

To quantify what we mean by a large effect, consider the overlap between the states obtained from one of the basis states of the two-level system, say $|L\rangle$, via two different evolution histories. These are constrained histories, or rather, paths in the functional integral: we prescribe times at which tunneling has occurred but let all the other degrees of freedom evolve freely. For example, suppose the first history is when nothing happens, and the second is when the system tunnels from $|L\rangle$ to $|R\rangle$, spends time Δt there, and then tunnels back. Using the corresponding evolution operators, we can write the overlap as

$$\langle L | U_2^\dagger(t_1, t_2; \Delta t) U_1(t_1, t_2) | L\rangle \sim \exp[-Q(\Delta t)] , \tag{4.1}$$

where $t_{1,2}$ are some initial and final moments of cosmic time, and Q will be called the decoherence exponent. The states $U_1|L\rangle$ and $U_2|L\rangle$ are almost indistinguishable within

[2] Hydrodynamics of a finite-temperature boundary was studied recently [11] using a real-time version of the standard *anti*-de Sitter/CFT correspondence. It is curious that here we find a connection with hydrodynamics for the very different (and much less understood) de Sitter case.

a single Hubble volume but, once Δt exceeds some minimal value, are very different at super-Hubble scales: the second state has two expanding bubble walls, while the first has none. As a result, the overlap (4.1) (including all modes, both sub- and super-Hubble) rapidly decreases to zero as a function of t_2.

The observer cannot measure directly the overlap (4.1) (since super-Hubble modes are involved), but he can measure the probability $P_L(t)$ to remain in the original state $|L\rangle$ after time t.[3] Because of the rapid decay of the overlap (4.1) (and of similar other ones), this probability is completely incoherent [2]:

$$P_L(t) = \frac{1}{2}(1 + e^{-\Gamma t}), \tag{4.2}$$

where Γ is the tunneling rate. Of course, in the dual description, with the environment residing on a finite-volume sphere (and therefore, in the presence of an ultraviolet cutoff, having a finite number of degrees of freedom), we cannot expect a complete decay of the overlap. (Similar points have been made in ref. [12] and, in a context quite close to ours, in ref. [4].) So, there will be additive corrections to eq. (4.2). By itself, that is not a problem, since the derivation of (4.2) has used classical gravity and so cannot exclude a residual overlap non-perturbative in $1/M_{\rm Pl}$, e.g., of the form $\exp(-{\rm const.} M_{\rm Pl}^2/H^2)$ envisioned in refs. [12, 4]. Still, this is a large suppression, and it has to be reproduced in the dual approach, where Q is viewed as a result of interactions with the environment.

High temperature T_0 is to rescue: the dual calculation of the decoherence exponent involves the imaginary part of the *time-ordered* correlator, and that is enhanced relative to (3.15) by the Bose factor $\coth(\omega/2T_0)$:

$$Q(\Delta t) \sim \sum_{lm} |\Delta\Phi_{lm}|^2 \int_{\omega_{\rm min}}^{\omega_{\rm max}} \frac{d\omega}{\omega^2} {\rm Im} G^R_{lm}(\omega)(1 - \cos\omega\Delta t) \coth\frac{\omega}{2T_0}, \tag{4.3}$$

where $\Delta\Phi_{lm}$ is the difference in the boundary value of the field between the states $|L\rangle$ and $|R\rangle$ (for a review on decoherence due to system-environment interactions, see ref. [13]). For Ohmic dissipation, one can replace $\coth(\omega/2T_0)$ by its high-temperature limit, so the decoherence exponent is proportional to $T_0\Delta t$ at times $\Delta t \ll 2\pi/\omega_{\rm min}$, where $\omega_{\rm min}$ is the minimal frequency (3.16), and reaches a plateau at larger Δt. From (3.3) and (3.9), we see that $T_0 \sim R$, so if we say that $R \sim M_{\rm Pl}$, then the overlap (4.1) computed in the dual theory acquires an appropriately strong exponential suppression: $Q(\Delta t \to \infty) \propto M_{\rm Pl}^2$.

5. Conclusion

Our main result is that an AdS/CFT-like correspondence between de Sitter bulk and a high-temperature "environment" at a near-horizon boundary survives a pair of most basic consistency checks. First, the correspondence requires Ohmic form of certain two-point functions in the boundary theory, and that is consistent with the presence of hydrody-

[3] The time t of the static coordinate system (2.2) can be written in terms of r and the cosmic time $t_{\rm cos}$ as $t = t_{\rm cos} - \frac{1}{2}\ln(1-r^2)$. Note that inside a bubble, the metric can be written in the form (2.2) about any point, not necessarily the bubble's center. So, we assume that the observer is at $r = 0$, where $t = t_{\rm cos}$.

namic modes at the boundary. Second, the large decoherence in a two-level system in de Sitter space is reproduced, up to expected corrections, by a dual calculation in which it is attributed to interaction with the environment.

In relating bulk gravitational dynamics to hydrodynamics at the boundary, there in no a priori reason to consider only equilibrium, thermal boundary environments. One may as well consider nonthermal states, ranging from nearly thermal (which will hopefully describe near-dS spaces) to highly nonthermal, such as turbulence. If the correspondence proposed here holds up under further scrutiny, it would be extremely interesting to obtain the geometric duals of such nonthermal states.

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