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Asymptotic Symmetries and Patient Observers in de Sitter*

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Abstract. The questions regarding the quantum origin of the universe, the black hole information problem, and the cosmological constant problem, can all be argued to involve infrared issues of quantum gravity. In order to make progress towards a solution of these problems, it may therefore make sense to first explore their possible infrared connections. In this proceeding, I will summarise some recent progress we have made in exploring infrared issues in de Sitter analogous to issues with black hole evaporation. In particular, I will here focus on recent work on asymptotic symmetries and patient observers. Patient observers are essentially observers with (gravitational) memory.

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1 Introduction

It has already been argued for a while that a connection between the cosmological constant problem and the quantum origin of the universe might exist. The tiny observed cosmological constant has been argued to be a consequence of environmental selection in the landscape [8, 10, 11], with the need of eternal inflation to populate all false vacua and physically realize all possible universes [9, 12]. Other solutions to the cosmological constant problem relating to infrared issues of inflation has also been speculated upon in the past [13–18].

Similarly, arguments has also been made for a connection between black holes and eternal inflation. In fact, Schwarzschild and de Sitter spacetimes have many similarities. On a heuristic level, they both have horizons, and there are indications of a perturbative breakdown due to the presence of infrared modes in both spacetimes on parametrically similar timescales [4, 6, 7] (see also [19–23]). Below we discuss how the results of [4, 6, 7] can be obtained in the language of

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asymptotic symmetries [1] and how that leads to some new insights regarding patient observers [2, 3].

2 Asymptotic Symmetries and Charges

Accelerating FLRW spacetimes, like de Sitter space, have distinct causal properties from Minkowski space due to the presence of horizons, and so the form of the relevant asymptotic symmetries are very different from the BvdBMS symmetries in flat spacetime [24, 25]. Nevertheless, an analogous procedure for an accelerating FLRW spacetime can be adopted, and the full group of transformations that leaves null infinity invariant can be deduced. The key distinction is that now null infinity coincides with the conformal surface at which both null and timelike lines terminate. The conformal tranformations are slice tranformations leaving the boundary data invariant, and in [26] the residual diffeomorphisms which preserves the asymptotic form of the spactime, but acts nontrivially on the boundary data, were found (in de Sitter) to be just the group of spatial diffeomorphisms. These play the same role as the BvdBMS symmetries in flat spacetime, in the sense that a soft graviton, as it redshifts to infinity, will act on the system as one of these transformations. As reviewed below, this can then be encapsulated by finding the precise diffeomorphism that corresponds to each soft graviton, and then tracking the change in the quantum state by using the Noether charge associated with that transformation.

To find the conserved charges generating the asymptotic symmetry transformations, we consider the perturbed spatially flat FLRW spacetime,

$$ds^2 = -dt^2 + a^2(t)e^{2\zeta}\left[e^{\gamma}\right]_{ij} dx^i dx^j \ , \tag{1}$$

where the perturbations are described by the scalar curvature perturbation ζ and the graviton γ_{ij}. We can further assume γ_{ij} to be transverse and traceless, and we have left out the lapse and shift, as they decay as powers of the scale factor, and so will be unimportant on the asymptotic boundary [27]. The asymptotic symmetries act on this metric as

$$\delta\left(e^{2\zeta}\left[e^{\gamma}\right]_{ij}\right) = \mathcal{L}_{\xi}\left(e^{2\zeta}\left[e^{\gamma}\right]_{ij}\right) \ , \tag{2}$$

where ξ_i is an arbitrary spatial vector, and $\mathcal{L}_{\xi}$ is the Lie derivative with respect to this vector.

The variation of the ζ and γ fields under this transformation are to leading order [1]

$$\delta\zeta = \xi\cdot\partial\zeta + \frac{1}{3}\partial\cdot\xi \tag{3}$$

and

$$\delta\gamma_{ij} = \partial_i\xi_j + \partial_j\xi_i - \frac{2}{3}\partial\cdot\xi\delta_{ij} + \xi\cdot\partial\gamma_{ij} + \mathcal{O}(\gamma^2). \tag{4}$$

The Noether charge can be constructed in the usual way [28]. Although it is unproblematic to use the Noether charge when it appears in commutators with fields, its main technical disadvantage is its volume divergence. However, in the language of Brown and York [29] their charge is a boundary term and local counterterms can be added to yield a finite quantity [30, 31]. We will need these counterterms to regulate the volume divergences. To see the equivalence of the Noether charge and the Brown and York charge we follow [1] and note that the Noether charge is given by

$$Q_N(\xi) = \frac{1}{2}\int d^3x\ \left[\Pi_\zeta \delta\zeta + \Pi_\gamma^{ij}\delta\gamma_{ij}\right] + h.c.\ , \tag{5}$$

where $\Pi_{\zeta,\gamma} = \delta\mathcal{L}/\delta\left(\dot{\zeta}, \dot{\gamma}_{ij}\right)$ are the canonical momenta and $\delta\zeta, \delta\gamma_{ij}$ are the transformations of the canonical variables ζ and γ_{ij} under the diffeomorphism ξ. The Hermitian conjugate must be added to ensure that the charge is real when dealing with quantum operators.

Focussing for simplicity on pure de Sitter and using that $2\tilde{\Pi}^{ij} = M_p^2\sqrt{g}(K^{ij} - Kg^{ij})$, it can be shown that the charge can be written as [1]

$$Q_N = \frac{M_p^2}{2}\int d^3x\sqrt{g}(K^{ij} - Kg^{ij})\nabla_{(i}\xi_{j)}\ , \tag{6}$$

where, K_{ij} is the extrinsic curvature, and the index-free version is traced with the full spatial metric, noting that this charge is defined exactly for the full theory. This form makes it very easy to see the equivalence between the Noether charge as defined in 5 and the Brown-York charge [1]. If the spatial integration is over a 3-dimensional hypersurface Σ, this can be integrated by parts to yield a boundary term (also using the momentum constraint, $\nabla_i\Pi^{ij} = 0$) associated with the 2-dimensional boundary $\partial\Sigma$ [1]

$$Q_{BY} = \int_{\partial\Sigma} d^2x\sqrt{\sigma}n^i\xi^j T_{ij}^{BY}\ , \tag{7}$$

where σ is the induced metric in $\partial\Sigma$, and n^i is a vector tangent to Σ and normal to $\partial\Sigma$. This is the Brown-York charge with the associated Brown-York tensor

$$T_{ij}^{BY} = M_p^2\left(K_{ij} - Kh_{ij}\right)\ . \tag{8}$$

Having an expression for the Noether charge for a generic diffeomorphism, we can find the unique charge corresponding to a precise soft mode by matching the effect the charge has on short wavelength modes encoded in the commutator $[Q, \mathcal{O}] = -i\delta\mathcal{O}$. From (3), (5) we have

$$[Q, \zeta(x)] = -i\xi\cdot\partial\zeta(x)\,. \tag{9}$$

Thinking of Q as the operator corresponding to a infinitesimal linearised symmetry transformation of the form $x \to x - \xi^L$, with ξ^L just being given by the linear term

$$\xi_i^L = \frac{1}{2}\gamma_{ij}^L x^j \ , \tag{10}$$

then, after successive applications of the commutator and letting the differential operators act on them self (but assuming still that the long mode, γ^L, is constant), one finds that a finite large symmetry transformations becomes the expected coordinate transformation [1]

$$e^{-iQ}\zeta(x)e^{iQ} = e^{-i[Q,\cdot]}\zeta(x) = e^{-\xi^L\cdot\partial}\zeta(x) = \zeta(xe^{\gamma^L/2}) \ . \tag{11}$$

3 Change in the Local Vacuum

We have just seen that the charge generates a large gauge transformation which is physically equivalent to the creation of soft modes (gravitons or scalars). This transformation captures the local evolution of the vacuum where modes become superhorizon as time evolves.

In expressions involving explicit expressions of the charge (not just commutators of the charge), one has to deal with the fact that the Noether charge is formally infinite (only in commutators of the charge, as above, the infinities cancels out and the problem is absent) being proportional to the spatial volume. This is a generic feature of charges constructed from spontaneously broken symmetries.

However, if we restrict to any finite region with volume V, all expressions will be finite. The price we pay for this is losing the property that the charge is conserved: since we deal with a finite region, charge can leak in or out of the region under consideration corresponding to the emission of a soft mode with time evolution as expected. In practice, we utilise the equivalent formulation of the charge in terms of the Brown-York charge, and renormalise the infrared volume divergences by adding adequate counterterms on the boundary by considering [1]

$$S = M_p^2 \left[\int d^4x\sqrt{-g^{(4)}}\frac{R}{2} + \int_{\mathcal{I}^+} d^3x\sqrt{-g^{(3)}}\left(-K[g]+ \right.\right.$$
$$\left.\left. \underbrace{-2c_1 H + \frac{c_2}{2}R[g] + c_3 R[g]^2 + c_4 R[g]_{ij}R[g]^{ij}\Big)\Big]}_{\text{Counter-terms}} \right. \ . \tag{12}$$

and determining the boundary term coefficients, c_i, such that the final result is independent of the finite volume, V, we choose, and our result becomes independent of the infrared regulator. Interestingly, the UV divergences appearing in the computations are also renormalised by the same choice of counter terms, such that our IR renormalisation condition simultaneously renormalise the UV divergences.

Consider the initial local vacuum state denoted by $|0\rangle$. In the presence of a long wavelength field the state undergoes the symmetry transformation

$$|0'\rangle = e^{iQ}|0\rangle \ , \tag{13}$$

where Q here and in the following denotes the renormalized charge.

In order to compare how much the two states differ we compute their overlap

$$\langle 0'|0\rangle = \langle 0|e^{iQ}|0\rangle = \langle 0|\left(1 + iQ - \frac{Q^2}{2} + \mathcal{O}\left(Q^3\right)\right)|0\rangle. \tag{14}$$

When this matrix element becomes much smaller than one it means that our standard description of the vacuum is no longer accurate. The overlap becomes

$$|\langle 0|0'\rangle| = \left|1 - \frac{1}{2}\langle Q_3^2\rangle_{\text{NC}} + \cdots\right| \ , \tag{15}$$

where $\cdots$ are terms suppressed by powers of $(H/M_p)^2$ and the subscript NC denotes non-contact terms, and the cubic term of the charge is

$$Q_3(t) = \frac{a^3 M_p^2}{16}\int_V d^3x\, \dot{\gamma}_{ij}\gamma^L_{ab}x^b\partial^a\gamma_{ij} + h.c. \tag{16}$$

Following our discussion we have considered the charge associated with a volume V in a space-like surface at the future infinity of de-Sitter space. After renormalisation, explicit evaluations yields [1]

$$\langle Q_3^2(t)\rangle_{\text{NC}} = \frac{8\pi^2}{45}\langle\gamma^L\gamma^L\rangle. \tag{17}$$

Recalling that γ^L is a particular soft mode of size larger than the volume V, we would now would like to average the overlap over many spheres of size V inside a bigger region, which we could think of as the total inflated patch, where the variations in the long mode become relevant. This double averaging procedure, $\langle\langle\cdots\rangle\rangle$, is similar to the procedure of [4], and is equivalent to applying the charge continuously over a given range of soft modes, as follows from the ergodicity of the process. The averaging of $\langle\gamma^2(x)\rangle$ is the variance of the long mode in the big volume and is given by [4]

$$\langle\gamma^2(x,t)\rangle = -\frac{H^2}{2\pi^2 M_p^2}\log\left(\frac{\Lambda_{\text{IR}}}{a(t)H}\right), \tag{18}$$

where $\Lambda_{\text{IR}} = H$ is the largest comoving scale, the initial scale factor having been set to 1. Therefore,

$$\langle\langle Q_3^2\rangle\rangle \simeq \frac{4}{45}\frac{H^2}{M_p^2}\log\left(a(t)\right), \tag{19}$$

which means, that the average overlap between the two vacuum goes to zero on the time scale [1]

$$\frac{H^2}{M_p^2}\log\left(a(t)\right)\simeq 1\,, \tag{20}$$

which, using $\log(a(t)) = Ht$, is exactly the equivalent of the Page time for the de Sitter: $t = M_p^2/H^3$, consistent with the earlier findings of [4].

4 Recovering Consistency Relations

In order to get a feel for the physical meaning of the change in the vacuum generated by the charge, we use the charge to recover several known consistency conditions for inflationary correlators [1]. This also serves as a consistency check.

We start by the derivation of the graviton 3-point function in the squeezed limit

$$\langle\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}\rangle\,, \tag{21}$$

where s_i are the polarizations and $q_1 \ll q_2, q_3$.

In the standard picture, this correlator is zero at tree-level and one needs to go to first order in the in-in formalism and insert one interaction Hamiltonian. In the charge picture it is the vacuum itself which is changed by the symmetry transformation. Then, in the new vacuum

$$\langle 0'|\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}|0'\rangle = \langle 0|e^{-iQ}\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}e^{iQ}|0\rangle \equiv \langle\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}\rangle_Q\,. \tag{22}$$

We proceed by expanding 22 in powers of Q

$$\langle 0|e^{-iQ}\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}e^{iQ}|0\rangle = \langle\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}\rangle - i\langle[Q,\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}]\rangle + \mathcal{O}(Q^2). \tag{23}$$

Evaluating the commutator explicitly using our expressions for the charge, one obtains [1]

$$\langle\gamma_{q_1}^{s_1}\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}\rangle_Q = \frac{3}{2}\,\epsilon_{ab}^{s_1}(q_1)\hat{q}_2^a\hat{q}_2^b\,\langle\gamma^L\gamma^L\rangle'\langle\gamma_{q_2}^{s_2}\gamma_{q_3}^{s_3}\rangle'\delta_{s_2,s_3}\delta^{(3)}(q_2+q_3). \tag{24}$$

which agrees with the result found in [33].

Similarly one can also demonstrate that the change in state is able to reproduce loop corrections to the correlators calculated in the literature. We choose to illustrate this for the two point scalar correlator, as computed in [4]. We can compute this by evaluating the expression

$$\langle 0'|\zeta\zeta|0'\rangle = \langle 0|e^{-iQ}\zeta\zeta e^{iQ}|0\rangle \equiv \langle\zeta\zeta\rangle_Q\,. \tag{25}$$

In order to evaluate this we need the scalar part of the charge

$$Q = \int d^3x\,\Pi_\zeta\left(D_L\zeta + \frac{1}{3}\nabla\cdot\xi\right)\,, \tag{26}$$

where $D_L = (\gamma_L/2 + \gamma_L^2/8)_{ij} x^i \partial_j + \ldots$. After explicit calculations, this yields [1]

$$\langle\zeta\zeta\rangle_Q = \left(1 + \langle\gamma_L^2\rangle \left(\frac{2}{5} k \cdot \partial_k + \frac{4}{15}(k \cdot \partial_k)^2\right)\right) \langle\zeta\zeta\rangle + \ldots \tag{27}$$

In agreement with [4]. If the running of the two point function is negligible this can be further reduced to the simple expression

$$\langle\zeta\zeta\rangle_Q = \left(1 + \frac{(4-n_s)(1-n_s)}{15}\langle\gamma_L^2\rangle\right) \langle\zeta\zeta\rangle + \ldots \tag{28}$$

as also discussed in [4], together with generalizations.

5 Bogoliubov Transformation

Here will discuss how the transformation $|0\rangle \to |0'\rangle$ also can be viewed as a Bogoliubov transformation, with the Bogoliubov coefficients defined as

$$\beta_k = \frac{u'_k v_k - v'_k u_k}{W} \tag{29}$$

$$\alpha_k = \frac{u'_k v^*_k - u_k v'^*_k}{W} \tag{30}$$

where $v(k) = u(\tilde{k})$ is the mode function in the shifted coordinates under a large gauge transformation with $\tilde{k}^2 = e^{2\zeta_L} e^{\gamma_L}_{ij} k_i k_j$, and W is the Wronskian. If we now identify u_k with γ_k the Wronskian would be given by $2i/(aM_p)^2$. It can be verified that the charge acts on the vacuum as a Bogoliubov transformation, which has the form [2]

$$|0'\rangle = \prod_k \frac{e^{\frac{-\beta^*}{2\alpha} a^\dagger_k a^\dagger_{-k}}}{|\alpha_k|^{1/2}} |0\rangle. \tag{31}$$

Using the expressions for the Bogoliubov coefficients defined in equations (29) and (30) we obtain, in the dS limit [2],

$$\alpha_k = \frac{1}{2i}\left[\left(\left(\frac{k}{\tilde{k}}\right)^{3/2} - \left(\frac{\tilde{k}}{k}\right)^{1/2}\right)\frac{1}{k\eta} - i\left(\left(\frac{k}{\tilde{k}}\right)^{1/2} + \left(\frac{\tilde{k}}{k}\right)^{1/2}\right)\right] e^{i(\tilde{k}-k)\eta}\,, \tag{32}$$

$$\beta_k = \frac{1}{2i}\left[\left(\left(\frac{\tilde{k}}{k}\right)^{1/2} - \left(\frac{k}{\tilde{k}}\right)^{3/2}\right)\frac{1}{k\eta} - i\left(\left(\frac{\tilde{k}}{k}\right)^{1/2} - \left(\frac{k}{\tilde{k}}\right)^{1/2}\right)\right] e^{i(\tilde{k}+k)\eta}\,. \tag{33}$$

It is easy check that the normalization condition $|\alpha_k|^2 - |\beta_k|^2 = 1$ is satisfied, and that on super-horizon scales, where the mode function u_k scales as $k^{-3/2}$, one has indeed

$$|u_{\tilde{k}}|^2 = \left(\frac{k}{\tilde{k}}\right)^3 |u_k|^2 \,, \tag{34}$$

which proves that the form of α_k and β_k are self consistent.

As a further non-trivial check of the Bogoliubov transformation, we can again consider a 3-point function in the squeezed limit, $k_L \ll k_1 \sim k_2$,

$$\langle 0|\gamma_{k_L}\zeta_{\tilde{k}_1}\zeta_{\tilde{k}_2}|0\rangle \tag{35}$$

where ζ_k has the usual mode expansion

$$\zeta_k = u_k a_k + u_k^* a_{-k}^\dagger \tag{36}$$

and the tilde quantities, which live in the background of the long mode, are related to the vacuum solutions by

$$u_{\tilde{k}} = \alpha_k u_k + \beta_k u_k^* \,. \tag{37}$$

Taking the super-horizon limit, $-k\eta \to 0$, of the expressions for α_k and β_k we obtain

$$\begin{aligned} &\langle 0|\gamma_{k_L}\zeta_{\tilde{k}_1}\zeta_{\tilde{k}_2}|0\rangle \\ &= (2\pi)^3\delta^{(3)}(k_1+k_2)\langle\gamma_{k_L}|\alpha_k|^2 + |\beta_k|^2 + 2\mathrm{Re}(\beta_k\alpha_k^*))\rangle|u_{k_1}|^2 \\ &= (2\pi)^3\delta^{(3)}(k_1+k_2)\langle\gamma_{k_L}\left(\frac{k}{\tilde{k}}\right)^3\rangle|u_{k_1}|^2 \,. \end{aligned} \tag{38}$$

Now using that to leading order $\tilde{k}^2 = k^2 - \gamma_{ij}^L k_i k_j$ and $\gamma_{ij}^L = \gamma_{k_L}\epsilon_{ij}$, as well as remembering that we treat the super horizon mode γ_{k_L} classically, i.e. $\langle\gamma_{k_L}^2\rangle$ should be thought of a classical averaging over the domain of interest and the use of the ergodicity theorem of quantum mechanics is implicit, we arrive at

$$\langle 0|\gamma_{k_L}\zeta_{\tilde{k}_1}\zeta_{\tilde{k}_2}|0\rangle = (2\pi)^3\delta^{(3)}(k_1+k_2)\frac{3}{2}\epsilon_{ij}\hat{k}_i\hat{k}_j|\gamma_{k_L}|^2|u_{k_1}|^2 \tag{39}$$

which is easily seen to reproduce the result of [33].

6 Patient Observers

The next logical question to ask is if this change in the local vacuum is observable to any physical observer? To answer the question, we remind ourselves that all ground states $|0\rangle$, $|0'\rangle$ are equivalent since they are connected by symmetry transformations, and yet they look physically inequivalent, having different expectation values for the charge. The resolution to this apparent paradox is related to measurement theory: Any apparatus that can measure the charge must be able to distinguish between the different states and therefore break the symmetry that connects those states. If we only allow for detectors that preserve the symmetry of the system, then the different ground states will be indistinguishable

This leads to our definition of a patient observer [2]: A patient observer is an observer that has been in existence prior to the time when a given long mode of interest was created, and is equipped with some measurement device, necessarily not invariant under the large gauge transformation, to record its state for the entire duration. In other words a patient observer can be said to carry gravitational memory.

We have presented two suggestions for patient observers in de Sitter, as thought experiments on how one would go about measuring the effects of long modes [2, 3]. While naively there seem to be no problem in constructing such patient observers, we do however encounter more subtle problems. Each time we are stymied by the requirement that we build our machine out of physically realizable matter. This does not constitute a proof that a measurement device of sufficiently clever design cannot be envisioned, but it does give some indication that there may be a version of "cosmic censorship" at play, preventing a patient from being able to discriminate the $|0\rangle$, $|0'\rangle$ vacuums in practice. This situation is reminiscent of [32], where it is argued that no machine capable of observing a single graviton may ever be constructed.

The first example is an array of satellites connected radially to some concentric point through some wires, just like a carousel. The rigidity of the wires prevents the radial expansion with the expansion of the universe but allows for angular displacement of the satellites. Alternatively we could think of the same array of satellites inside some rigid circular tube, like a hula hoop, where they would be allowed to move frictionless along the tube if some shear acts on them.

As a graviton passes through this setup, it induces a shear stress, and so displaces the satellites from their original angular positions. The relative change in the position is given by

$$\left(\frac{\Delta x_{\text{IR}}}{x}\right)^2 \sim \langle\gamma_L^2\rangle \sim \frac{H^3}{M_p^2}t. \tag{40}$$

Though this effect grows with time, to address its observability we need to compare it with the fundamental uncertainty in position coming from the Heisenberg uncertainty principle, $\Delta p \Delta x_q \gtrsim 1$. This implies

$$\Delta x_q^2 \gtrsim \frac{t}{m} \sim \frac{M_p^2}{mH^3}\left(\frac{\Delta x_{IR}}{x}\right)^2 \sim \left(\frac{L_H}{r_s}\right)\left(\frac{L_H}{x}\right)^2 \Delta x_{\text{IR}}^2 \tag{41}$$

where r_s is the Schwarzschild radius associated with the detector of mass m, and L_H is the horizon size. Therefore, in order for the effect from long modes to be greater than the quantum uncertainty, either the detectors must be separated by a distance greater than the horizon, or they must be so massive that their Schwarzschild radii are larger than the horizon. Thus, quantum effects prevent a local observer from measuring long modes in this way [2, 3].

Another realization of a patient observer consists of comoving satellites exchanging electromagnetic signals, in hopes of measuring a relative time delay induced by the long modes through the redshift of light. Each satellite has its own clock and both are synchronized when the device is built. During the time the satellites are in causal contact they will measure the redshift and deduce from there either a change in time or in the Hubble constant through Hubble's law $z = Ht$. The time shift induced by the soft gravitons can be read directly from the metric to be

$$\Delta t_{\text{IR}} \sim \frac{\langle \gamma_L^2 \rangle}{H} \sim \frac{H^2 t}{M_p^2}\,. \tag{42}$$

However, in order for the clocks to have a precision Δt, they must have an energy uncertainty $\Delta E \gtrsim 1/\Delta t$. Then, as pointed out in [34], physical clocks must interact with each other gravitationally, causing time dilation effects that will shift the time between clocks by an uncertain amount

$$\Delta t_{\text{dilation}} \sim \frac{\Delta E}{M_p^2 x} t\,. \tag{43}$$

Comparing this to the effect of the long mode, we find

$$\Delta t_{\text{dilation}} \sim \left(\frac{\Delta E}{H}\right)\left(\frac{L_H}{x}\right)\Delta t_{IR} \tag{44}$$

Therefore, when the physical requirements that the detector be able to fit inside the Hubble horizon and the minimal energy probed by the satellite cannot be less than the Hubble rate are imposed, we find that these again exactly forbid the effect to be measured [2, 3].

We have performed various extensions of these simple systems, as well as several other novel strategies for measurement, and each time we have run into the same problem. It therefore seems likely that the conjecture, that the effect is not observable to a single physical observer during the de Sitter like expansion, holds true.

7 Conclusions

We have reformulated the soft theorems for cosmological correlation functions [33, 35, 36] and loops [4] in terms of asymptotic symmetries and charges [1], confirming the findings that the quantum state changes by an order one factor after a Page time has elapsed even in de Sitter [4,6,7]. The soft modes becomes physical in questions where it cannot be completely gauged away. Previously we have discussed such situations where inflation ends and modes re-inter the horizon [6], and also global measures of the change in the geometry [7]. We recently revisited the question regarding the physical observer during the de Sitter like phase in terms of a patient observer with gravitational memory [2,3]. However, in this case, we found subtle quantum mechanical fluctuations preventing the observation of soft modes, which seems to support a new cosmic censorship for patient observers [2,3].

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