Planckian Physics in the Sky

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Abstract

In this talk I review a few ideas on how to connect the physics of quantum gravity with cosmology. In particular, I review a possible modulation of the power spectrum of primordial density fluctuations generated through transplanckian (maybe stringy) effects. I also discuss the prospects of finding effects due to holography during inflation.

1 Introduction

In the coming years we can expect substantial advances in fundamental physics thanks to two new circumstances. The first is that string theory has reached a maturity which makes it possible to construct detailed cosmological models. The second is that the precision of cosmological observations is rapidly increasing. Important areas include the time variation of the expansion rate of the universe and its relation to a possible cosmological constant, as well as the fluctuation spectrum of the CMBR. It is widely believed that cosmology could be the key to the verification of string theory, and that string theory could be what we need to solve several of the present puzzles in cosmology.

Many attempts to describe the early universe are based on inflation, for an excellent review see [1]. The defining feature of inflation is a stage of accelerated expansion which is believed to be necessary to solve a few embarrassing problems of the Big Bang model, in particular the flatness and horizon problems.

The flatness problem relates to the observation that the real density of the universe, ρ, has long been known to be very close to the critical density ρ_c. That is, $\Omega = \frac{\rho}{\rho_c}$ is close to one. To understand the importance of this, it is useful to rewrite the Friedmann equation

$$H^2 = \frac{8\pi G}{3}\rho - \frac{k}{a^2}, \tag{1}$$

in the form

$$\Omega - 1 = \frac{k}{a^2 H^2}. \tag{2}$$

It is easy to see that for any ordinary type of matter, $\frac{1}{a^2 H^2}$ will *increase* with time. One finds, for example, that pressureless dust gives $\sim t^{2/3}$ and radiation gives $\sim t$. From this one concludes that, unless the universe is flat ($k = 0$) and therefore has *exactly* $\Omega = 1$, Ω will rapidly evolve away from $\Omega = 1$. If one starts with a value $\Omega < 1$ the value will decrease towards zero, while if $\Omega > 1$ the value of Ω will increase and even diverge if the expansion stops. In order to have a value close to 1 today, one would therefore expect to have a value of Ω even closer to 1 in the early universe. It is easy to see that the necessary fine tuning is one part in 10^{16} a second after the Big Bang and one part in 10^{60} at planckian times.

This fine tuning is certainly extremely unnatural. A possible way out would be that there is some kind of mechanism at work in the early universe that dynamically drives Ω towards 1. This is where inflation comes in. Inflation corresponds to a period where $\frac{1}{a^2 H^2}$ actually *decreases*. This is the case, for an expanding universe, if the scale factor a, that is the distance between two test objects, increases faster than the horizon radius $1/H$. This is the same as to say that the expansion of the universe is accelerating, i.e. $\ddot{a} > 0$.

Inflation also solves the horizon problem. That is, how seemingly unconnected parts of the universe can look so similar. The reason is, as explained above, that the expansion rate is in a very definite sense faster than the speed of light. Objects in causal contact can, through the expansion, be separated to distances larger than the Hubble radius. Eventually, after inflation stops, the Hubble radius will grow faster than the expansion and the objects will return within their respective horizons. An observer not taking inflation into account will wrongly conclude that these objects have never before been in causal contact.

This is how inflation works, but how do we get inflation? From the Friedmann equation

$$\frac{\ddot{a}}{a} = -\frac{4\pi G}{3}(\rho + 3p) \tag{3}$$

it follows that an accelerated universe requires matter with negative pressure. Luckily, this can be provided by a scalar field, the *inflaton*, which possesses a potential energy. The energy density, for a homogenous inflaton field, is given by

$$\rho = \frac{\dot{\phi}^2}{2} + V(\phi) \tag{4}$$

while the pressure is given by

$$p = \frac{\dot{\phi}^2}{2} - V(\phi). \tag{5}$$

A simple example that yields inflation is the case of a cosmological constant where $p = -\rho$, that is, de Sitter space.

Much of contemporary cosmology has dealt with the construction of phenomenologically viable inflationary models with various potentials and number of inflatons. A remaining question is how to fit inflation into the rest of physics. Recently string theory has come up with an interesting candidate for inflation using branes. The idea is to use two stacks of branes separated by a certain distance, corresponding to the inflaton, in a higher dimensional space. This is a rapidly developing subject – for an early review see [2], for some more recent discussions, see [3].

How do we test inflation? The key is structure formation. One reason to invoke inflation was to make the universe smooth and flat. However, in the real universe there is certainly a large amount of structure. This structure can be traced back to subtle fluctuations in the matter distribution during the time when the CMBR was released. A naive application of inflation does, however, exclude such non-uniformity. So, from where does it come? Actually, inflation itself provides the answer provided we take quantum mechanics into account.

Inflation magnifies microscopic quantum fluctuations into cosmic size, and thereby provides seeds for structure formation. The details of physics at the highest energy scales is therefore reflected in the distribution of galaxies and other structures on large scales. The fluctuations begin their life on the smallest scales and grow larger (in wavelength) as the universe expands. Eventually they become larger than the horizon and freeze. That is, different parts of the wave can no longer communicate with each other since light can not keep up with the expansion of the universe. This is a consequence of the fact that the scale factor grows faster than the horizon, which, as we have seen, is a defining property of an accelerating and inflating universe. At a later time, when inflation stops, the scale factor will start to grow slower than the horizon and the fluctuations will eventually come back in again. The imprint of the quantum fluctuations can be studied in the CMBR revealing important clues about physics at extremely high energies in the early universe.

Inflation is a wonderful opportunity to connect the physics of the large with physics of the small. In the rest of my talk I will discuss two possibilities that I find especially intriguing: transplanckian signatures and holography. Transplanckian signatures refers to the intriguing idea of using inflation as a giant microscope to study physics near the Planck scale. Perhaps effects of physics beyond the Planck scale might be visible on cosmological scales in the spectrum of the cosmic microwave background (CMBR) fluctuations? I will also discuss holography as the means of connecting physics at the largest and smallest scales.

2 Transplanckian signatures

As mentioned above, quantum fluctuations play an important role in the theory of inflation. But how is the structure of these microscopic fluctuations determined? In a time dependent background – where there are no global timelike Killing vectors – the definition of a vacuum is highly non-trivial. In the ideal situation one has an initial, asymptotically Minkowsky like region, where one can define an in-vacuum. This vacuum will then time evolve through an intermediate time dependent era and then end up in a final Minkowsky like region. Typically the initial vacuum will not evolve into the final vacuum but rather correspond to an excited state with radiation. One example is a star that collapses into a black hole where it is well known that one obtains Hawking radiation. Technically, the excited state is related to the vacuum through a Bogolubov transformation.

It is easy to see that a similar phenomenon can be expected during inflation. But here the situation is, in a sense, even tricker since the universe (in Robertson-Walker coordinates) *always* is expanding. How do we choose the initial state in an unambiguous way? Luckily the very same property – the acceleration – which makes inflation so successful in other respects, can also help us here. If we follow a given fluctuation backwards in time, its wavelength will become arbitrarily smaller than the horizon radius. This means that deviations from Minkowsky space will become less and less important and the vacuum essentially unique – *the Bunch-Davies vacuum*. The fact that a unique vacuum is picked out is an important property of inflation and is an example of how inflation does away with the need to choose initial conditions.

But, and this is the main point, this argument relies on our ability to follow a mode to infinitely small scales which comes into conflict with the presence of a fundamental scale – Planckian or stringy – where physics could be completely different. Clearly this is an assumption that needs to be carefully scrutinized. Could there be effects of new physics that will change the predictions of inflation? In particular one could worry about changes in the predictions of the CMBR fluctuations. Several groups have investigated various ways of modifying high energy physics in order to look for such modifications, see e.g. [4-13].

Let us now, following [8], provide a typical example of the kind of corrections one might expect due to changes in the low energy quantum state of the inflaton field due to transplanckian effects. First, we need to find out when to impose the initial conditions for a mode with a given (constant) comoving momentum k. To do this it is convenient to use conformal coordinates rather than the standard Robertson-Walker coordinates where the inflating metric is given by

$$ds^2 = dt^2 - a(t)^2 dx^2, \tag{6}$$

with the scale factor given by $a(t) = e^{Ht}$. The conformal coordinates are obtained by defining $\eta = -\frac{1}{aH}$ and the metric takes the form

$$ds^2 = \frac{1}{H^2\eta^2}\left(d\eta^2 - d\mathbf{x}^2\right). \tag{7}$$

We now note that the physical momentum p and the comoving momentum k are related through

$$k = ap = -\frac{p}{\eta H}, \tag{8}$$

and impose the initial conditions when $p = \Lambda$, where Λ is the energy scale important for the new physics. This scale could be the Planck scale or the string scale. We find that the conformal time when the initial condition is imposed to be

$$\eta_0 = -\frac{\Lambda}{Hk}. \tag{9}$$

As we see, different modes will be created at different times, with a smaller linear size of the mode (larger k) implying a later time.

To proceed we need the equation of motion for the scalar field (ignoring the potential)

$$\ddot{\phi} + 3H\dot{\phi} - \nabla^2\phi = 0, \tag{10}$$

which in terms of the conformal time η, and the rescaled field $\mu = a\phi$, becomes

$$\mu_k'' + \left(k^2 - \frac{a''}{a}\right)\mu_k = 0. \tag{11}$$

When quantizing this system (a nice treatment can be found in [14]) one needs to introduce oscillators $a_k(\eta)$ and $a^\dagger_{-k}(\eta)$ such that

$$\begin{aligned} \mu_k(\eta) &= \frac{1}{\sqrt{2k}}\left(a_k(\eta) + a^\dagger_{-k}(\eta)\right) \qquad (12) \\ \pi_k(\eta) &= \mu_k'(\eta) + \frac{1}{\eta}\mu_k(\eta) = -i\sqrt{\frac{k}{2}}\left(a_k(\eta) - a^\dagger_{-k}(\eta)\right), \end{aligned}$$

obey standard commutation relations. The crux of the matter is that these oscillators are time dependent, and can be expressed in terms of oscillators at a specific moment only using the Bogolubov transformations

$$\begin{aligned} a_k(\eta) &= u_k(\eta)\, a_k(\eta_0) + v_k(\eta)\, a^\dagger_{-k}(\eta_0) \qquad (13) \\ a^\dagger_{-k}(\eta) &= u_k^*(\eta)\, a^\dagger_{-k}(\eta_0) + v_k^*(\eta)\, a_k(\eta_0), \end{aligned}$$

where

$$|u_k(\eta)|^2 - |v_k(\eta)|^2 = 1. \tag{14}$$

From this it is clear that the choice of vacuum is a highly non-trivial issue. Without knowledge of the high energy physics we can only list various possibilities and investigate whether there is a typical size or signature of the new effects. A useful example is to choose the vacuum determined by

$$a_k(\eta_0)\,|0, \eta_0\rangle = 0. \tag{15}$$

This vacuum should be viewed as a typical representative of other vacua besides the Bunch-Davies. It can be characterized as a vacuum corresponding to a minimum uncertainty in the product of the field and its conjugate momentum, [14], the vacuum with lowest energy (lower than the Bunch-Davies) [5], or as the instantaneous Minkowsky vacuum[1]. Therefore, it is as special as the Bunch-Davies vacuum and there is no a priori reason for

[1] As observed in [15] the exact caracterization of the vacuum depends on the canonical variables used.

planckian physics to prefer one over the other. Note that for $\eta_0 \rightarrow -\infty$ the Bunch-Davies vacuum is recovered.

We are now in a position to calculate the expected fluctuation power spectrum. Following [8] one finds

$$\begin{aligned} P(k) &= \left(\frac{H}{\dot{\phi}}\right)^2 \left\langle |\phi_k(\eta)|^2 \right\rangle = \left(\frac{H}{\dot{\phi}}\right)^2 \frac{1}{a^2} \left\langle |\mu_k(\eta)|^2 \right\rangle \\ &= \left(\frac{H}{\dot{\phi}}\right)^2 \left(\frac{H}{2\pi}\right)^2 \left(1 - \frac{H}{\Lambda} \sin\left(\frac{2\Lambda}{H}\right)\right). \end{aligned} \tag{16}$$

This result should be viewed as a typical example of what to be expected from transplanckian physics if we allow for effects which at low energies reduce to changes compared to the Bunch-Davies case. We note that the size of the correction is linear in H/Λ, and that a Hubble constant that varies during inflation gives rise to a modulation of the spectrum. As argued in [8], the modulation is expected to be a quite generic effect that is present regardless of the details of the transplanckian physics. (See also [10] for a discussion about this). After being created at the fundamental scale the modes oscillate a number of times before they freeze. The number of oscillations depend on the size of the inflationary horizon and therefore changes when H changes. A varying Hubble constant is crucial for a detectable signal, since a Hubble constant that does not vary during inflation would just imply a small change in the overall amplitude of the fluctuation spectrum which would not constitute a useful signal. Luckily, since the Hubble constant is expected to vary, the situation is much more interesting.

Let me now turn to possible observable consequences. I will discuss what happens using slow roll parameters, see, e.g., [1], where in particular,

$$\varepsilon = \frac{M_{pl}^2}{2} \left(\frac{V'}{V}\right)^2, \tag{17}$$

with $M_{pl} = 1/\sqrt{8\pi G} \sim 2 \cdot 10^{18} GeV$ as the (reduced) Planck mass. It is not difficult to show (using that H is to be evaluated when a given mode crosses the horizon, $k = aH$) that

$$\frac{dH}{dk} = -\frac{\varepsilon H}{k}, \tag{18}$$

which gives

$$H \sim k^{-\varepsilon}. \tag{19}$$

The k dependence of H will translate into a modulation of $P(k)$, with a periodicity given by

$$\frac{\Delta k}{k} \sim \frac{\pi H}{\varepsilon \Lambda}. \tag{20}$$

To be more specific, let me consider a realistic example. In the Hořava-Witten model [16], unification occurs roughly at the same scale as a fifth dimension becomes visible and also comparable to the string scale and the higher dimensional Planck scale. For a discussion and references see, e.g., [17] or [7]. As a rough estimate we therefore put $\Lambda = 2 \cdot 10^{16}$ GeV. This is a rather reasonable possibility within the framework of the heterotic string and corresponds to $\gamma = 0.01$. The Hubble constant during inflation can not be much larger than $H = 7 \cdot 10^{13}$ GeV, corresponding to $\varepsilon = 0.01$. Using these values we find

$$\frac{H}{\Lambda} \sim 0.004$$
$$\frac{\Delta k}{k} = \Delta \ln k \sim 1.$$

This means one oscillation per logarithmic interval in k, which fits comfortable within the parts of the spectrum covered by high-precision CMBR observation experiments.

As I have already emphasized, it is important to note that the transplanckian effects, regardless of their precise nature, have a rather generic signature in form of their modulation of the spectrum. If it had just been an overall shift or tilt of the amplitude it would not have been possible to measure the effect even if it had been considerably larger than the percentage level. The shift would just have gone into a changed value or slope of, e.g., H. With a definite signature we can use several measurement points throughout the spectrum, as discussed in more detail in [13]. There it was argued that the upcoming Planck satellite might be able to detect transplanckian effects at the 10^{-3} level, which would put the Hořava-Witten model within range, or at least tantalizingly close. In this way one can also beat cosmic variance that otherwise would have limited the sensitivity to about 10^{-2} at best. (To make a more definitive statement, a more careful analysis of covariance of the transplanckian signature with a large number of other cosmological parameters will be needed, something we leave for future work.)

There has been extensive discussions of these results in the literature and their relevance for detectable transplanckian signatures. As pointed out in [9] the initial condition approach to the transplanckian problem allows for a discussion of many of the transplanckian effects in terms *the α-vacua.*

These vacua have been known since a long time,[18], and corresponds to a family of vacua in de Sitter space that respects all the symmetries of the space time.

In [7][19] it has been pointed out that there could be tricky problems with field theories based on non trivial vacua of this sort. None of these problems are, however, necessarily relevant to the issue of transplanckian physics in cosmology for a very simple reason, [12]. The whole point with the transplanckian physics, as explained in the introduction, is to find out whether effects beyond quantum field theory can be relevant for the detailed structure of the fluctuation spectrum of the CMBR. In the real world we do expect quantum field theory to break down at high enough energy to be replaced by something else, presumably string theory. The modest proposal behind [8] is simply that we should allow for an uncertainty in our knowledge of physics near planckian scales.

Another problem is that the α-vacua are non-thermal. One can therefore worry that they will not survive but instead quickly relax to the thermal Bunch-Davies vacuum. This is one of the questions that will be answered when we now to turn to the subject of entropy and holography.

3 Holography

3.1 Holographic bounds

After the rather phenomenological approach to the problem of finding observable hints of quantum gravity that I so far have followed, I will now consider a more specific approach. The idea is to see if holography can give some new insights. Holography has its origin in black hole physics and the discovery in the 70's by Bekenstein that black holes carry an entropy proportional to the area of the horizon, [20]. Bekenstein further argued that there are general bounds on the amount of entropy that can be contained in matter. The entropy bound that will serve as a starting point for our discussion states that in asymptotically flat space, [21],

$$S \leq S_B = 2\pi E R, \tag{21}$$

where E is the energy contained in a volume with radius R. This is the *Bekenstein bound*. There are several arguments in support of the bound when gravity is weak [22], and it is widely believed to hold true for all reasonable physical systems. Furthermore, in the case of a black hole where

$R = 2El_p^2$, we have an entropy given by

$$S_{BH} = \frac{A}{4l_p^2} = \frac{\pi R^2}{l_p^2}, \tag{22}$$

which exactly saturates the Bekenstein bound. We will consequently put $\hbar = c = 1$, but explicitly write the Planck length, $l_p = \sqrt{\frac{G\hbar}{c^3}}$, to keep track of effects due to gravity.

Beginning with [23], there have been many attempts to apply similar entropy bounds to cosmology and in particular to inflation, [24]. The idea has been to choose an appropriate volume and argue that the entropy contained within the volume must be limited by the area. An obvious problem in a cosmological setting is, however, that for a constant energy density a bound of this type always will be violated if the radius R of the volume is chosen to be big enough. However, as was explained in [25], it is not reasonable to discuss radii which are larger than the Hubble radius in the expanding universe. See also [26]. This, then, suggests that the maximum entropy in a volume of radius $R > r$, where r is the Hubble radius, is obtained by filling the volume with as many Hubble volumes as one can fit – all with a maximum entropy of $\frac{\pi r^2}{l_p^2}$. This gives rise to the *Hubble bound*, which states that

$$S < S_H \sim \frac{R^3}{r^3}\frac{r^2}{l_p^2} = \frac{R^3}{rl_p^2}. \tag{23}$$

The Hubble bound is a bound on the entropy that can be contained in a volume much larger than the Hubble radius. It is, therefore, a bound that gives measurable consequences only if inflation stops allowing scales larger than the inflationary Hubble radius to become visible. One should also note that the notion of a cosmological horizon, and its corresponding area, does not play an important role from this point of view since all entropy is present in matter, possibly in the form of black holes.

If we, on the other hand, want to discuss things from the point of view of what a local observer, that does not have time to wait for inflation to end, can measure, we must be more careful. In this case one has a cosmological horizon with an area that it is natural to give an entropic interpretation [27]. Since the area of the horizon grows when matter is passing out towards the horizon, from the point of view of the local observer, it is natural to expect the horizon to encode information about matter that, in its own reference frame, has passed to the *outside* of the cosmological horizon of the local observer. From the point of view of the observer, the matter will never be seen to leave but rather become more and more redshifted. The outside of the cosmological horizon should, therefore, be compared with the inside of a black hole. It follows that the horizon only indirectly provides bounds on entropy within the horizon as is nicely exemplified through the *D-bound* introduced in [28]. The cosmological horizon area in a de Sitter space with some extra matter is smaller than the horizon area in empty space. If the matter passes out through the horizon, the increase in area can be used to limit the entropy content in matter. This is the content of the D-bound which turns out to coincide with the Bekenstein bound. The D-bound, therefore has not, necessarily, that much to do with de Sitter space or cosmology. It is more a way to use de Sitter space to derive a constraint on matter itself.

Let me now explain the nature and relations between the various entropy bounds a little bit better. In particular on what scales the entropy is stored. If we assume that all entropy is stored on short scales smaller than the horizon scale r, we can consider each of the horizon bubbles separately and use the Bekenstein bound (or D-bound) on each and everyone of these volumes. We conclude from this that the entropy, under the condition that it is present only on small scales, is limited by

$$S < S_{LB} = 2\pi E r,$$

which we will refer to as the *local Bekenstein bound*. It is interesting to compare this result with the entropy of a gas in thermal equilibrium. One then finds $S_g \lesssim Er$ for high temperatures where $T \gtrsim 1/r$, and $S_g \gtrsim Er$ for low temperatures where $T \lesssim 1/r$. This is quite natural and a consequence of the fact that most of the entropy in the gas is stored in wavelengths of the order of $1/T$. This means that the entropy for low temperatures is stored mostly in modes larger than the Hubble scale and can therefore violate the local Bekenstein bound S_{LB}.

The size of the horizon therefore limits the amount of information on scales larger than the Hubble scale, or, more precisely, the large scale information that once was accessible to the observer on small scales. If the horizon is smaller than its maximal value this is a sign that there is matter on small scales and the difference limits the entropy (or information) stored in the matter. This is the role of the D-bound. We conclude, then, that a system with an entropy in excess of S_{LB} (but necessarily below S_H) must include entropy on scales larger than the horizon scale.

While the entropy bounds above are rather easy to understand, the way entropy can flow and change involve some more subtle issues. In the case of a diluting gas the expansion of the universe implies a flow of entropy out through the horizon, but as the gas eventually is completely diluted the flow of entropy taps off. Whether or not the horizon radius is changing, one will never be able to violate the Hubble bound or get an entropy flow through an apparent horizon violating the bound set by the area. A potentially more disturbing situation is obtained if we consider an empty universe (apart from a possibly changing cosmological constant), which can be traced arbitrarily far back in time, with entropy generated through the quantum fluctuations that are of importance for the CMBR. As discussed in several works, [14][29], there is an entropy production that can be associated with these fluctuations and one can worry that this will imply an entropy flow out through the horizon that eventually will exceed the bound set by the horizon. This is the essence of the argument put forward in [30].

To understand this better one must have a more detailed understanding of the cause of the entropy. Entropy is always due to some kind of coarse graining where information is neglected. In the case of the inflationary quantum fluctuations we typically imagine, as I have explained, that the field starts out in some pure state – defined by some possibly transplackian physics – with a subsequent unitary evolution that keeps the state pure for all times. This is true whether we take the point of view of a local observer or use the global FRW-coordinates. To find an entropy we obviously must introduce a notion of coarse graining. Various ways of coarse graining have been proposed, but they all imply an entropy that grows as the state get more and more squeezed, [14][29]. It can be shown that most of this entropy is produced at large scales (when the modes are larger than the horizon), and well below the Hubble bound.

This is all in terms of the FRW-coordinates, but let us now take the point of view of the local observer. In this case the freedom to coarse grain is more limited. In order to generate entropy we must divide the system into two subsystems and trace out over one of the subsystems in order to generate entropy in the other. As an example consider a system with N degrees of freedom divided into two subsystems with N_1 and N_2 degrees of freedom, respectively, with $N = N_1 + N_2$ and $N_2 > N_1$. If the total system is in a pure state it is easy to show that the entropy in the larger subsystem is limited by the number of degrees of freedom in the smaller one, i.e. $S_2 < \ln N_1$.[2] Applied to our case, this means that the entropy flow towards the horizon must be balanced by other matter with a corresponding ability to carry entropy within the horizon. Since the amount of such matter is limited by the D-bound, the corresponding entropy flow is also limited. As a consequence, there can not be an accumulated flow of entropy out towards the horizon that is larger than the area of the horizon. For a

[2] A simple proof can be found in [32] in the context of the black hole information paradox.

similar conclusion see [31]. This does not mean that inflation can not go on for ever, nor that there can not be a steady production of entropy on large scales, but it does imply that the local observer will not be able to do an arbitrary amount of coarse graining.

To summarize: *from a local point of view the production of entropy in quantum fluctuations is limited by the ability to coarse grain; from a global point of view entropy is created on scales larger than the Hubble scale.*

3.2 Complementarity

I have now argued that holography, in the sense of putting limits on the entropy, does not constrain cosmology in any new way. It might still be a useful principle, but it does not contain anything beyond the Bekenstein bound and the generalized second law which, in turn, seem to be automatically obeyed by the ordinary laws of physics.

If we want to find truly new effects, we must go one step further and turn to the principle of *complementarity*. I will therefore investigate the possibilities of an information paradox and compare with the corresponding situation in the case of black holes. In black hole physics the general view that has emerged is that a kind of complementarity principle is at work implying that two observers, one travelling into a black hole and the other remaining on the outside, have very different views of what is going on. According to the observer staying behind, the black hole explorer will experience temperatures approaching the Planck scale close to the horizon, and as a consequence, the black hole explorer will be completely evaporated and all information transferred into Hawking radiation. According to the explorer herself, however, nothing peculiar happens as she crosses the horizon.

As explained in [33] the apparent paradox is resolved when one realizes that the two observers can never meet again to compare notes. Any attempts of the observers to communicate again, after the outside observer has extracted the information from the Hawking radiation, will necessarily make use of planckian energies and presumably fail.

An interesting question to pose is whether a similar mechanism could be at work also in de Sitter space. In order to investigate such a possibility in the case of de Sitter space, we will consider a scenario where at some point in time de Sitter phase is turned off and replaced by a non-accelerated $\Lambda = 0$ phase with ordinary matter. That is, an inflationary toy model. Our discussion is therefore of great relevance to ongoing discussions of whether or not holography and other effects of quantum gravity are of importance for inflation. As I will argue later, and this is one of my main conclusions, the time-scales involved in our thought experiment suggest that direct effects of holography and complementarity are likely to be heavily suppressed in a physically realistic model of inflation.

A possible information paradox, comes about if one assumes that an object receding towards the de Sitter horizon of an inertial de Sitter observer, will return its information content to the observer in the form of de Sitter radiation. If the cosmological constant turns off, the object itself will eventually return to the observer's causal patch, and one has the threat of an information paradox.

To come to terms with the paradox, let us focus on what an observer actually would see as an object recedes towards the horizon, [34]. Since the rate of the photons (emerging from the horizon) received by our observer is of order $1/R$, the time it would take for her to see the object being thermalized will be extremely long. To find out how long, we will now try to find out what actually happens to the object (according to the observer). To do that we think of the horizon as an area consisting of R^2/l_{Pl}^2 Planck cells, and remember that the photon has a wavelength of order Planck scale when emitted and can indeed resolve specific Planck cells.

Now, let us first assume the object in question to be something really simple, corresponding to an information content much smaller than the R^2 number of degrees of freedom of the horizon. This would mean that only a few of the Planck cells are involved in encoding the object. In the extreme case of an object with entropy of $\mathcal{O}(1)$, one would need to wait until of the order R^2 photons have been emitted to be reasonably sure to see a photon coming from the burning of the object. In the other extreme one can think of an object consisting of the order R^2 degrees of freedom. In this case it is clear that one has to wait until of the order R^2 photons have been emitted, in order for all parts of the object to have been burnt. Regardless of the size of the object, one has, therefore, to wait a time,

$$\tau \sim \frac{R^2}{l_{Pl}^2} R = \frac{R^3}{l_{Pl}^2} \sim \frac{1}{T^3 l_{Pl}^2}, \tag{24}$$

in order to actually see the destruction. We therefore suggest that the maximal lifetime of non-thermal excitations in de Sitter space is given by $\tau \sim R^3/l_{pl}^2$. It is important to emphasize that this relatively long time scale follows only if other non-gravitational interactions are frozen. If this is not the case, and those interactions occur faster than the expansion of the universe, the characteristic thermalization time will scale like the naive $1/T$. In that case we do not need any reference to holography or complementarity. Thermalization occurs regardless of perspective.

Let us now abruptly turn off the de Sitter phase, and let it be followed by a more standard cosmological evolution. When this happens, the object will return to the observer's causal patch at some time in the future [35]. Since the situation between the object and the observer is symmetric, it is clear that the object will be in as good shape as the observer. If the observer has not been thermalized by then, then neither should the object be. Indeed, considering the symmetric situation we have between the observer and the object (being for example another observer) and the fact that they can meet again some time after the de Sitter phase has turned off, seems to imply that the estimated thermalization time should be the same for local objects as for those who approach the horizon, even from the perspective of one single observer.

Now let us now try to estimate the time it takes to thermalize a local object, bound to the observer. To do this, we reconsider the possibility that local interactions do give rise to a thermalization, but only if we take physics near the Planck scale into account. With an interaction rate given by $\Gamma = \sigma n v$, where the cross section is given by $\sigma \sim l_{pl}^2$, the number density of the radiation $n \sim T^3 \sim 1/R^3$ and the relative velocity $v = c = 1$, one finds the typical time τ it takes for this process to occur to be

$$1 \sim \sigma n v \tau \sim l_{pl}^2 \cdot 1/R^3 \cdot 1 \cdot \tau \ \Rightarrow \ \tau \sim \frac{R^3}{l_{pl}^2}. \tag{25}$$

This coincides, up to orders of one, with the previous result. Therefore, regardless of whether local objects or objects falling towards the horizon are concerned, the thermalization time will be the same. We argued above that this must be the case based on the symmetry between the observer and the object and by noting that, if the de Sitter phase is only temporary, they will eventually meet again. We find it encouraging that the above results are in agreement with this assessment.

The above analysis provides a possible escape route from the information paradox, since, as I have argued, it is very difficult for an observer to exist long enough to actually see any object being fully burnt by Hawking radiation. But this is not all, as observed in [35] there is a further obstacle to experiencing an information paradox. It can be shown that the return time for an object that has been falling towards the horizon a time $\tau \sim \frac{R^3}{l_{pl}^2}$, is of the order of the Poincare recurrence time, $\sim e^{R^2/l_{pl}^2}$ of the de Sitter space. That is, it exceeds the Poincare recurrence time of the detector.

What are the implications for inflation? In inflation the Hubble constant is constrained from observations to be no larger than $H \sim 10^{-4}\ m_{pl}$. With this input the thermalization time for non-thermal excitations (α-vacua included) is found to be of order $\tau \sim R^3 = 1/H^3 \sim 10^{12}\ t_{pl}$. Comparing this with the time needed for the required number of e-foldings, which for 70 e-foldings is $t_{infl} \sim 70/H \sim 7 \cdot 10^5\ t_{pl}$, one concludes that the thermalization time allows for visible effects of non-thermal behavior in the CMBR, with room to spare. This is good news for the transplanckian signatures. On the other hand, with fluctuations leaving the horizon so close to the end of inflation, effects from holography and complementarity are expected to be subtle.

4 Conclusions

In this talk I have discussed a few possibilities of using cosmological observations to test string theory. In particular I have argued that physics near the Planck scale might leave an imprint on the CMBR and give the first glimpses of how Nature works on its smallest scales. I have also discussed holography and the open question of whether issues like entropy bounds and complementarity could be of importance for inflation.

So far, our theoretical understanding of string theory and quantum gravity in a cosmological setting is not good enough to make definite predictions of what one might find. A lot of work certainly remains to be done. But from what we already know, I believe it is reasonable to hope that we one day will see planckian physics in the sky.

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