Horizon Physics and String Theory

Based on work of <u>many</u> contributors and collaborators.

(New results mentioned in this talk based on work w/M.Dodelson, Kang; Flauger, Mirabayi, Senatore; Munchmeyer/Peiris...*Planck*)

Plan:
I. Horizons in relativity and observations
II. Subtleties with featureless horizon physics: (how) do low energy field theory and general relativity break down?
Applications to black hole physics and inflationary cosmology

General Relativity:

``Curvature = Stress-energy"

$$S = \int d^4x \sqrt{g}\, \frac{R}{G_N} + S_{matter} \rightarrow R_{\mu\nu} - \frac{1}{2} g_{\mu\nu} R = 8\pi G_N T_{\mu\nu}$$

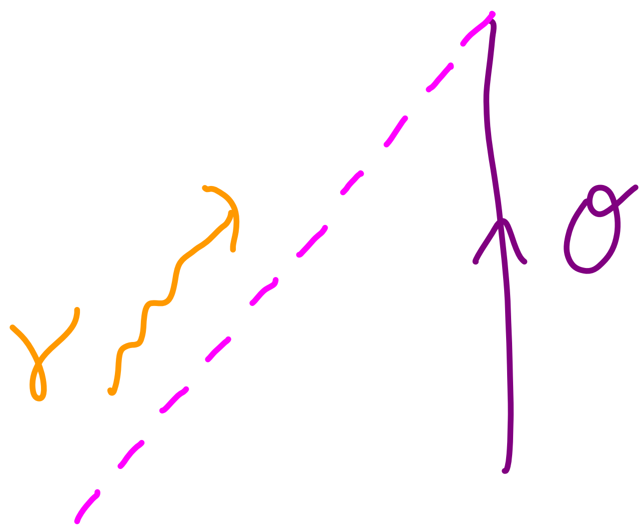
This leads to curved spacetimes with horizons, e.g.:

*black hole formation sourced by sufficiently dense matter

*accelerated expansion of the universe sourced by approximately homogeneous stress-energy

Horizons:

The Einstein equations have solutions in which observers lose causal contact with some regions of spacetime.



The energetics of probes strongly affected: get large redshifts and relative boosts.

For example, de Sitter (exponentially expanding) cosmology:

$$ds^2 = -dt^2 + a(t)^2 ds^2_{space}$$

+ perturbations

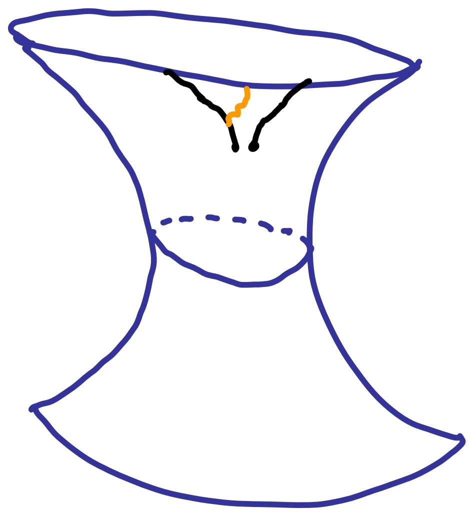
$$H(t)^2 \equiv \left(\frac{\dot{a}}{a}\right)^2 = \frac{\rho}{3 M_p^2} - \frac{k}{a^2}$$

$$\left(M_p^2 = \frac{\hbar c \equiv 1}{8\pi G_N} \right)$$

$$\rho \simeq \text{const} \Rightarrow a \to a_0 e^{Ht}$$

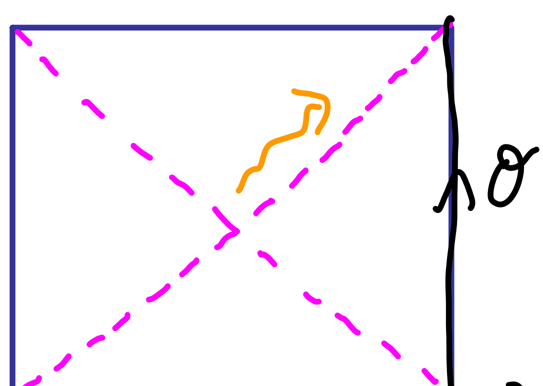
late times

$$ds^2 = -d\tau^2 + L^2\cosh^2\frac{\tau}{L}\,d\Omega^2$$



$$= \frac{1}{\cos^2 T}\left(-dT^2 + d\Omega^2\right)$$

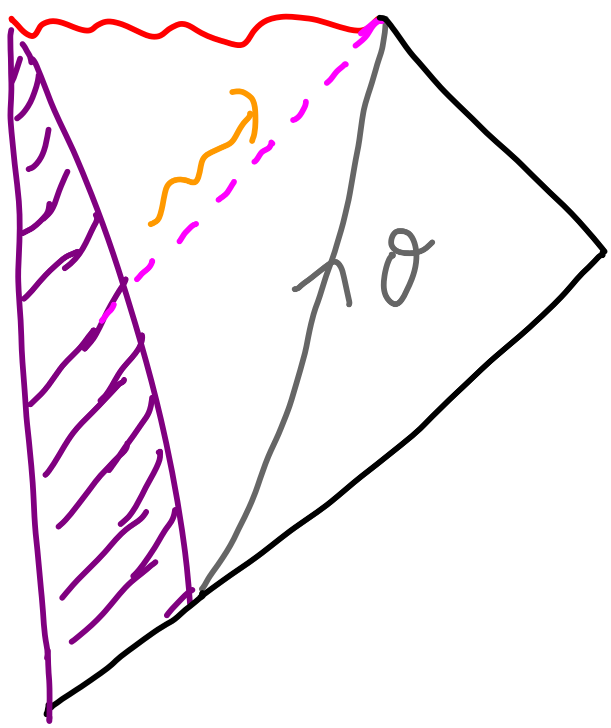
Light ray : $ds^2 = 0$ $\left(\overset{1}{\overset{|||}{c}}\,dt = dx \text{ in flat spacetime}\right)$



$$= -\left(1-\frac{r^2}{r_0^2}\right)dt^2 + \frac{dr^2}{1-\frac{r^2}{r_0^2}} + r^2 d\Omega^2$$

Black Hole horizon

$$ds^2 = -\left(1-\frac{r_s}{r}\right)dt^2 + \frac{dr^2}{1-\frac{r_s}{r}} + r^2 d\Omega^2$$



Classically, exterior geometry largely independent of formation matter, just depending on mass, spin, charge. There is substantial evidence that this is just a coarse-grained description.

Black hole thermodynamics and statistical mechanics

$\Delta A \geq 0$ in classical GR

area of event horizon

$\Delta M = \kappa \, \Delta A$ in classical GR

surface gravity

$T \sim \kappa$ Quantum mechanically (Hawking Radiation)

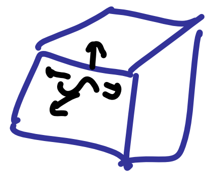
Dynamics of BHs $\Leftrightarrow$ thermodynamics

$\frac{\text{Area}}{G_N} \sim (\underline{\underline{\text{Area}}}\, M_p^2) \Leftrightarrow$ Entropy S

Gravity $\neq$ Local QFT
(d dimensions) (d dimensions)

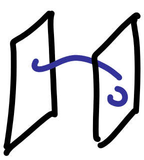
$\uparrow$ $\uparrow$

If you pack too much entropy in local region
$\Rightarrow$ form a BH, $S \sim$ Area



volume's worth of degrees of freedom

String theory: BH Stat. Mech.



Explicit microstate count for tractable BH's ... $\longrightarrow$

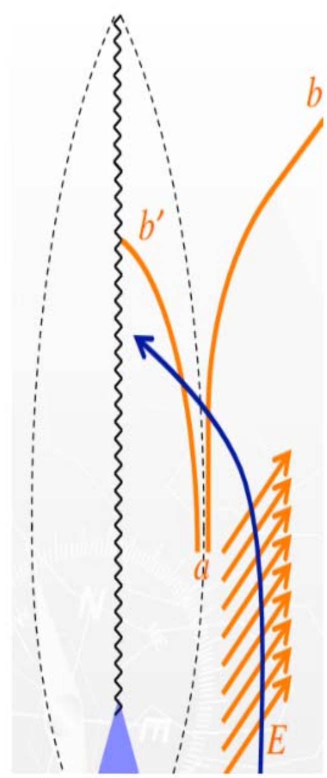
Strominger/Vafa, Maldacena, Klebanov et al ...

Gravity (d dimensions, e.g. AdS) $=$ Local QFT (d-1 dimensions) !

$\Rightarrow$ Unitary BH evolution

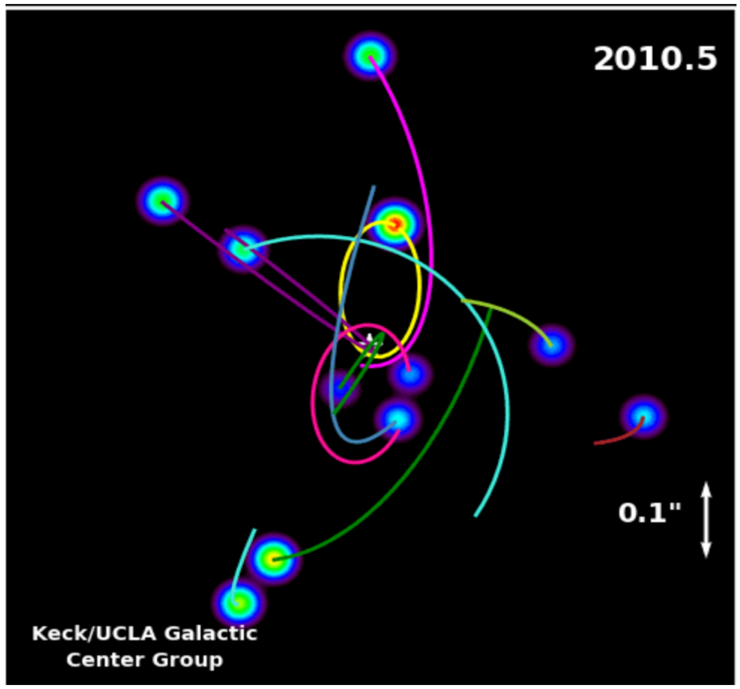
Hawking evaporation calculation gave thermal (information-losing) result => Since AdS/CFT says otherwise, something has to give.

**Hawking radiation thermal (no info); comes from vacuum near horizon: if unitary, then not vacuum near horizon
--monogamy of entanglement
*`Complementarity' was an idea that this problem cannot be seen by a single observer. ...AMPS '12...:
--There are observers who can see too much for `complementarity'
--Thought exp't constraints so far rule out all existing theories, something must give
1)Unitarity: AdS/CFT (Hawking agrees!)
2) EFT outside horizon, smooth infall
How get around above? so far all attempts at smooth horizon violate QM
*Horizon drama for infaller (e.g. `firewall', or variant)? **Need dynamics**

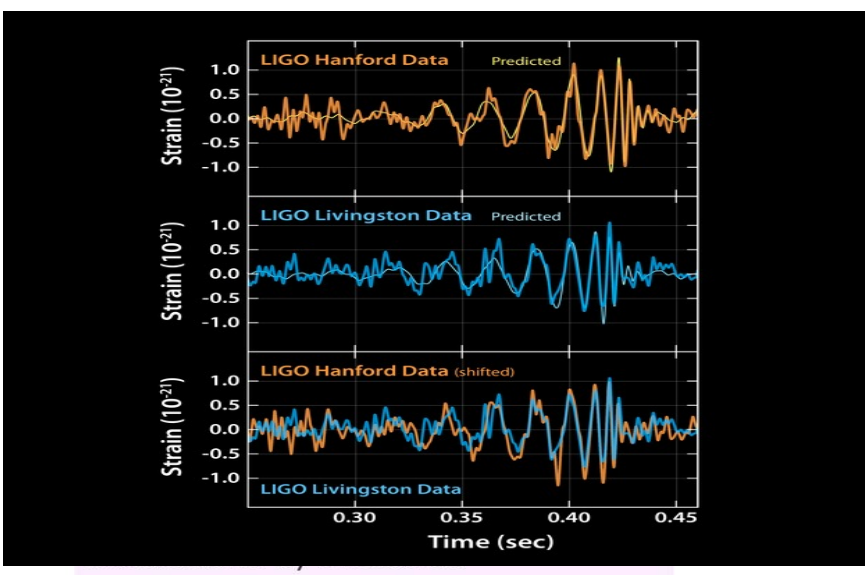


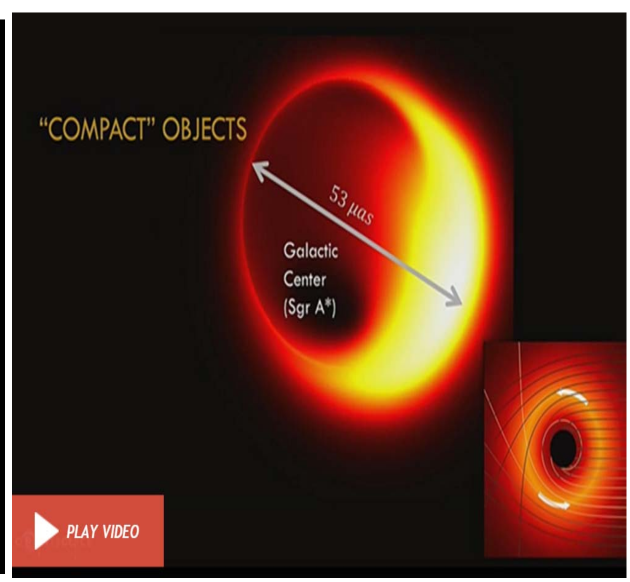
There is substantial observational evidence for both kinds of horizons

black holes



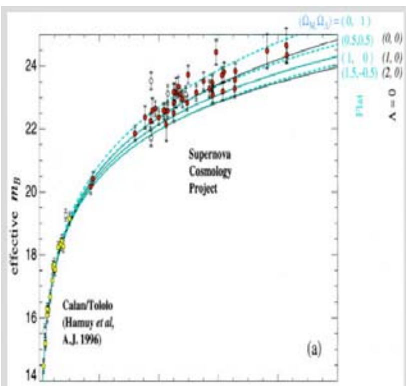
Sensitivity to EM and gravitational radiation closer in

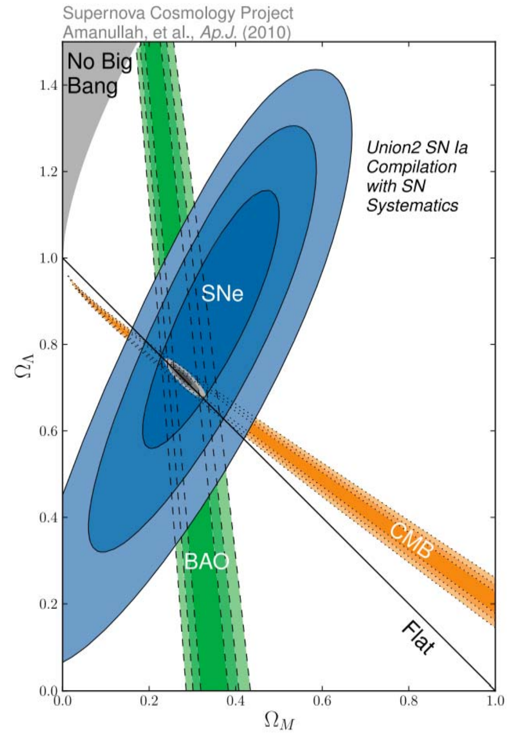


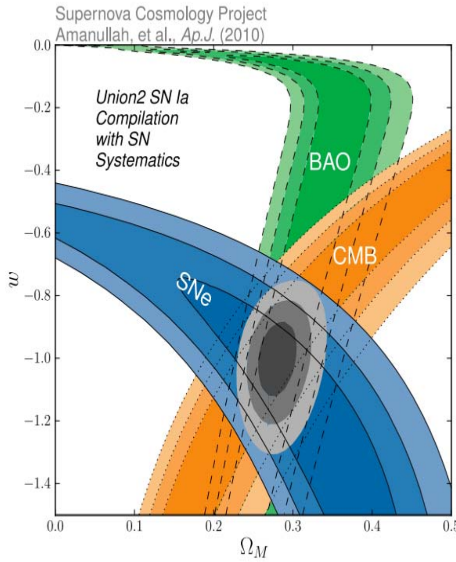


Cosmological horizons

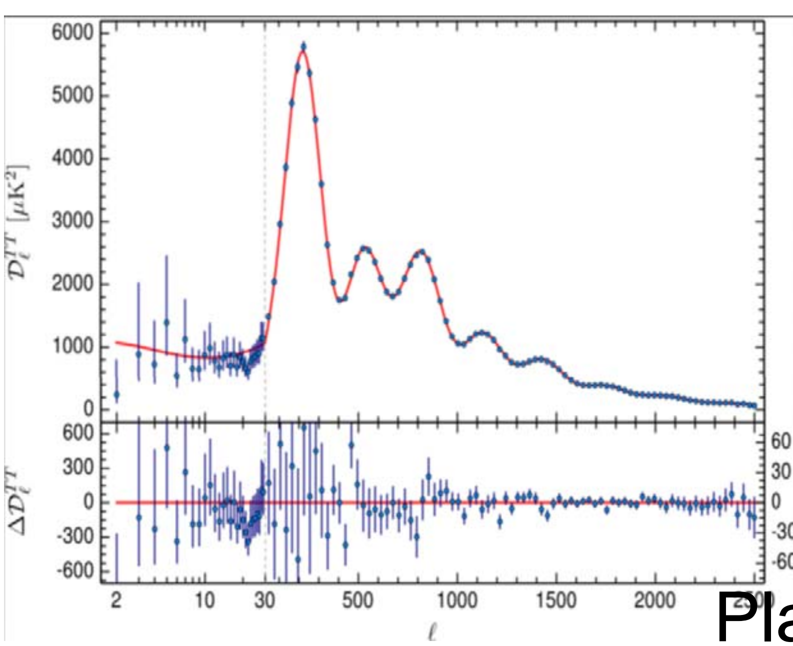
Late universe acceleration

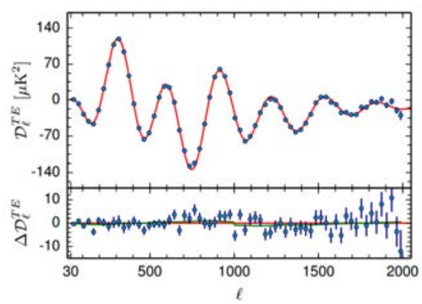


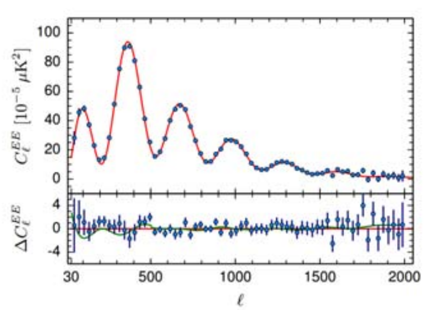




Early U: super-horizon perturbations at recombination, highly nontrivial test of inflationary cosmo







Planck

Curvature is weak (compared to Planck scale) at horizon (it grows large at singularity). As a result, for many purposes, GR and quantum field theory suffice to describe the (coarse grained) physics. However, this description is incomplete in important ways, even at the horizon, far from singularity.

Effective dimensionless coupling

$$\lambda_G = \frac{\text{Energy}^2}{M_p^2}$$

Effective Field Theory and dangerous irrelevance:

Standard method parameterizing our ignorance of high(er) energy physics

GR breaks down for $\lambda_G \to 1$ (or before)

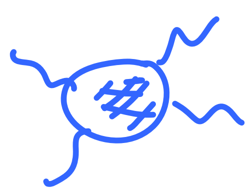
classical & Quantum corrections

$$\mathcal{S} = \int \left(\frac{R}{G_N} - V(\phi)\right)\left(1 + R\left(\frac{c_1}{M_*^2} + \hat{c}_1 G_N\right) + \cdots\right)$$

$$+ \int (\partial\phi)^2 + k_1 \frac{(\partial\phi)^4}{M_*^2} + \cdots$$

$M_*^2 \leftarrow$ scale of "new physics"

with corrections sensitive to short-distance physics



e.g. scalar quantum fields

$$S = \int d^4x \left\{ \dot{\phi}^2 - (\vec{\partial}\phi)^2 - \frac{1}{2}\underbrace{\lambda_2 M_*^2}_{m^2}\phi^2 \right.$$

$$+ \frac{\lambda_4}{4!}\phi^4 + \frac{\lambda_6}{6!\,M_*^2}\phi^6 + \cdots$$

$$\left. - \lambda_{4,2}\frac{(\partial\phi)^2\phi^2}{M_*^2} + \cdots \quad \frac{\lambda\,\mathcal{O}_\Delta}{M_*^{\Delta-4}} \right\}$$

ϕ has dimensions of Energy

$\mathcal{O}_\Delta$ corrections $\sim \left(\dfrac{\text{Energy}}{M_*}\right)^{\Delta-4}$

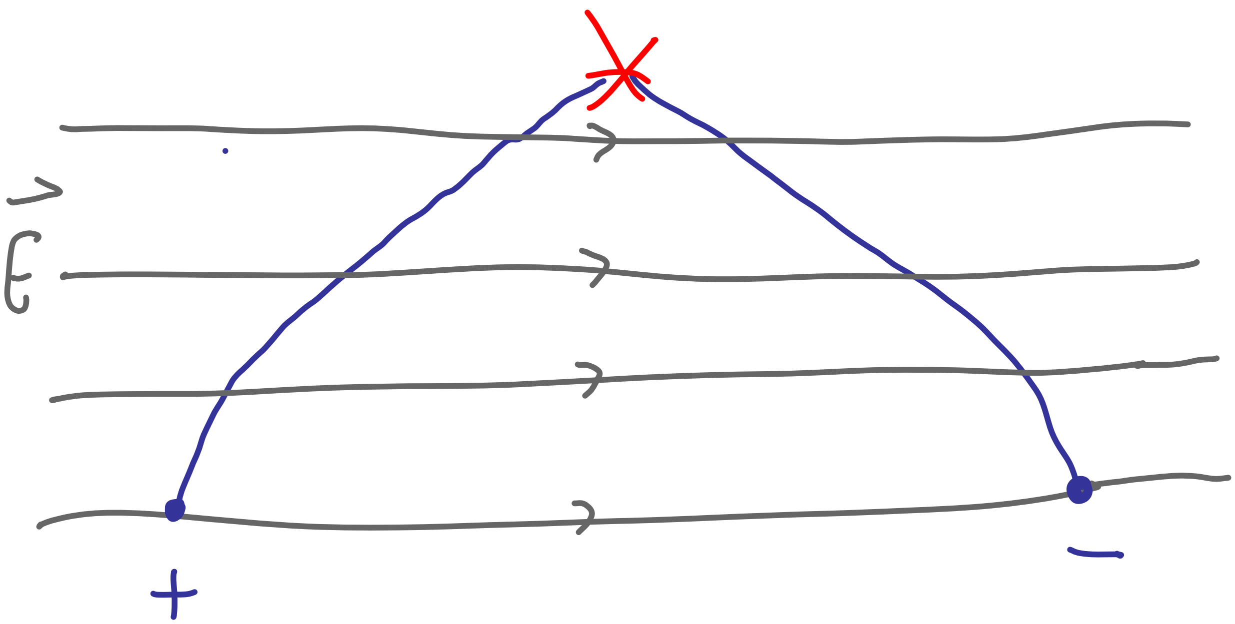
$$\mathcal{O}_\Delta \text{ corrections} \sim \left(\frac{\text{Energy}}{M_*}\right)^{\Delta - 4}$$

There is an infinite sequence of \`irrelevant' perturbations, those with $\Delta > 4$. Since these die out at low Energy, we can often make reliable physical predictions despite our ignorance of this infinite sequence.

However, physics *can* become sensitive to \`UV completion' even in systems with low input energies.

This subtlety arises in the presence of long time evolution and/or large field excursions.

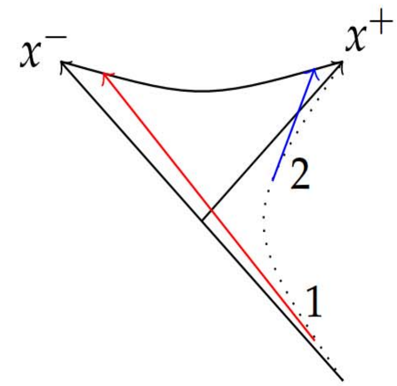
An electromagnetic example:
Consider a weak electric field permeating space, with two charges initially sitting at rest.



The weak field accelerates charges over a long time, producing a large invariant energy

A similar effect occurs in weakly curved geometries with a horizon: evolution of trajectories of (say) two probes sent in with modest energy leads to a large *nonlocal* invariant energy in the near horizon region.

$$ds^2 = -\left(1-\frac{r_s}{r}\right)dt^2 + \left(1-\frac{r_s}{r}\right)^{-1}dr^2$$
$$= -\underbrace{\frac{2r_s}{r}e^{1-r/r_s}}_{r \to r_s}dx^+\,dx^-.$$

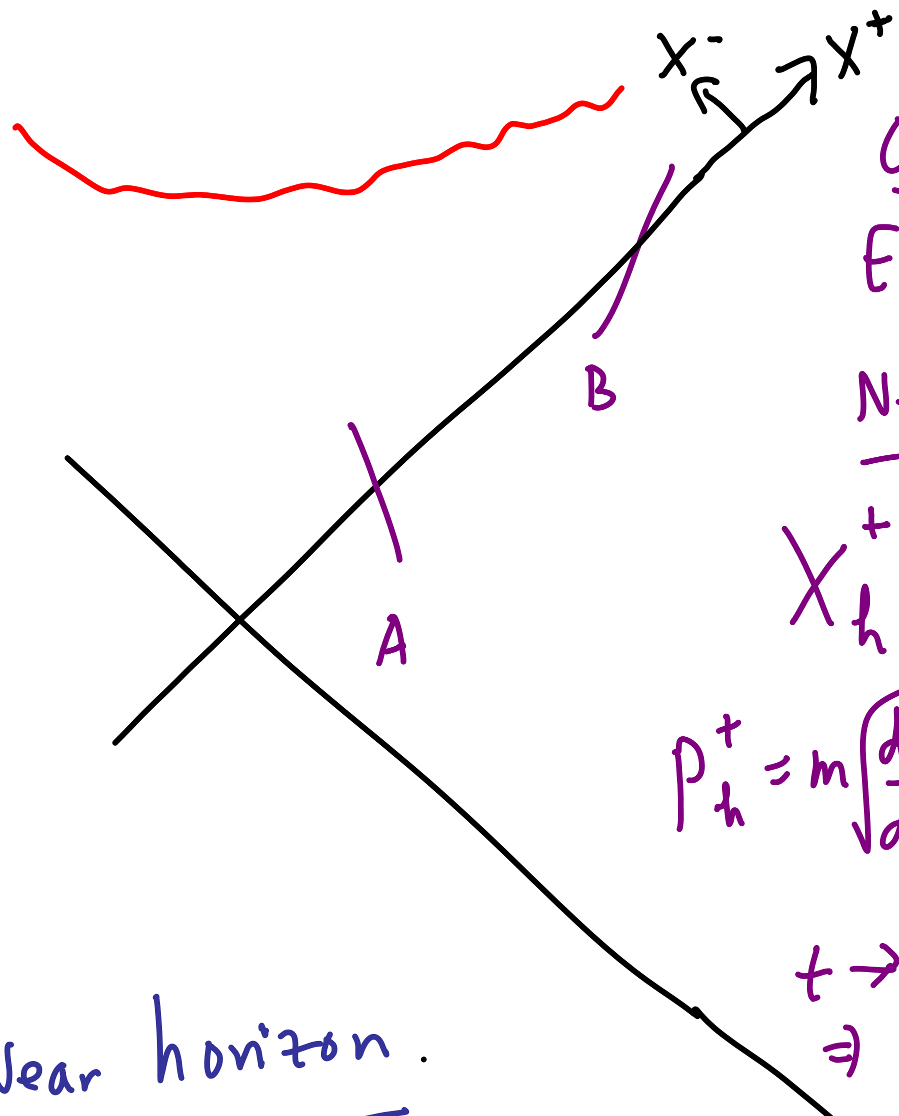


get patch of flat (Minkowski/Rindler) spacetime $ds^2 = -d\hat{t}^2 + d\hat{x}^2$

$$= -2\,dx^+dx^-$$

$x^\pm = \hat{t} \pm \hat{x}$ light cone coordinates

$$ds^2 = -\frac{2r_s}{r}\, e^{1-\frac{r}{r_s}}\, dX^+ dX^- + r^2 d\Omega^2$$



Outside:

E, m fixed

Near Horizon:

$$X_h^+ = 2r_s\sqrt{e}\,\frac{E}{m}\,e^{\eta}$$

$$p_h^+ = m\sqrt{\frac{dX^+}{dX^-}}\Bigg|_h = m e^{\eta}$$

$$t \to t + \Delta t$$

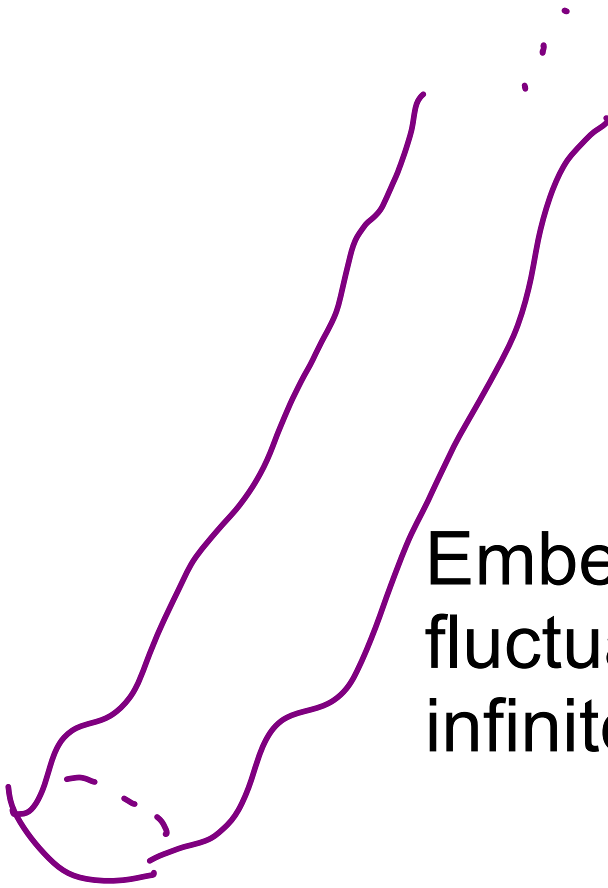
$$\Rightarrow \eta \to \eta + \frac{\Delta t}{2r_s}$$

Near horizon.

$$\left.\begin{aligned} &\cdot\ E^2_{m.h.} \sim 2p^+_{B,h}\, p^-_{A,h} \sim e^{\frac{\Delta t}{2r_s}}\, m^2 \\ &\cdot\ X_B^+ - X_A^+ \propto p_B^+ \propto e^{\frac{\Delta t}{2r_s}} \end{aligned}\right\}$$

Exponentially large, non-local energy in A-B system

<u>String Theory</u> is a good candidate for the theory of quantum gravity, let's consider the question in that framework.



$X^{\mu}(\sigma, \tau)$

Embedding in spacetime fluctuates: formally infinite mean square size

$\langle XX \rangle$

because of high-frequency modes. Need high energy probe to

Near horizon: huge Energy, but separated along X^+.

String Spreading

- Susskind '94
- Brown Polchinski Strassler Tan '06

Light Cone gauge $X^- \sim p^- \tau$,

Constraint determines X^+ in terms of $X^\perp$

$$\langle \psi | (X_\perp - x_\perp)^2 | \psi \rangle = \sum_n^{n_{max}} \frac{1}{n} = \log \frac{n_{max}}{n_0} + O\left(\frac{1}{n_{max}}\right)$$

$$\langle \psi | (X^+ - x^+)^2 | \psi \rangle \simeq \frac{1}{(p^-)^2} \sum^{n_{max}} n \simeq \frac{n_{max}^2}{(p^-)^2}$$

$n_{max} \leftrightarrow$ light cone time resolution

This effect, if not an artifact of a particular gauge choice (i.e. choice of coordinates), leads to an interaction between early and late infallers that goes beyond the predictions of naive effective field theory (i.e. not controlled by the weak curvature). The acceleration of the trajectories exponentially amplifies the energy. However, to get a substantial effect, the late infalling system needs somewhat (but not exponentially) high local energy (e.g. photons very near the horizon).

A similar setup to the black hole appears at six points in tree level flat space string scattering

Exhibits interaction as predicted by longitudinal spreading (M. Dodelson & ES '15, and to appear)

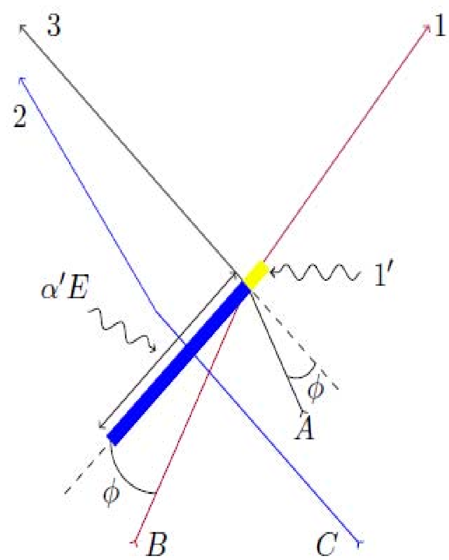
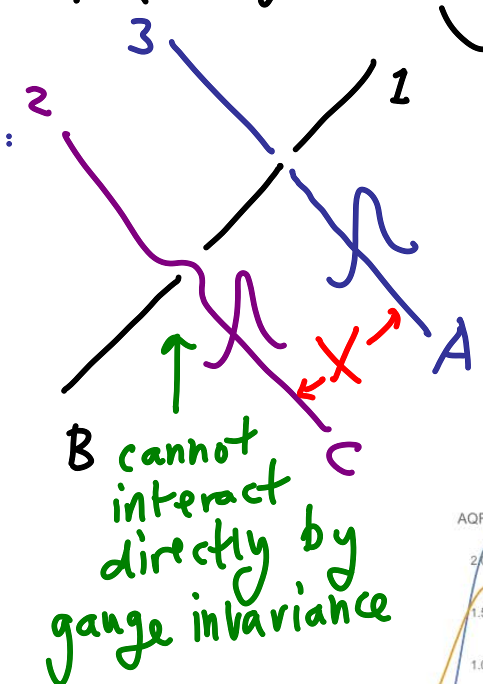


Figure 1: The setup for one version of our process, with peak trajectories shown. The central trajectories of strings A and B collide at $T = X = Y = 0$, producing outgoing strings 1 and 3. We introduce the third incoming string C, with kinematics such that string 1 (and the closely related Pomeron $1'$) has optimal light cone time resolution to detect the longitudinal spreading of strings C and 2 (shown in blue), for any value of the angle ϕ as in (2.15). The amplitude for the process decomposes into the $AB \to 1'3$ and $C1' \to 21$ four-point amplitudes times the $1'$ Pomeron propagator. It has a phase which implies peak support for the $C \to 2$ trajectories propagating through the interaction region a time T_{C*} of order $\alpha' E$ before the center of mass collision of A and B, independently of the value of ϕ, consistently with the longitudinal spreading predictions of [1, 2]. We will refer to this process in the kinematic regime well described by $1'$ exchange as 'case $1'$'.

C and 1' interact at long range, as predicted by light cone calc.

(paper to appear shortly)

$$A(X) \equiv \int \underbrace{\psi_X(\tilde{p})}_{\text{wavepacket peaked at } X = X_C - X_A} A(\tilde{p})$$



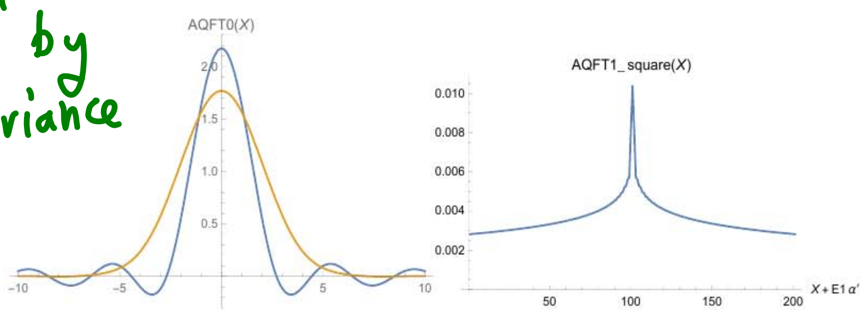


Figure 5: The distribution $A(X)$ for the two QFT comparison models, in the square wave packet for B described in the text. First, we plot the six point contact interaction (our QFT_0 comparison model), plotted for $\sigma = 1/\sqrt{\alpha'}$ and $\sigma = 1/(2\sqrt{\alpha'})$. Its width is $1/\sigma$, the width of the wavepacket, and does not extend out to $\sim \alpha' E$. We plot the numerical integral for the other example with a single massless $1'$ pole (QFT1) for $\sigma = 1/\sqrt{\alpha'}$. The parameters are within the regime described above, with E_1 of order 100.

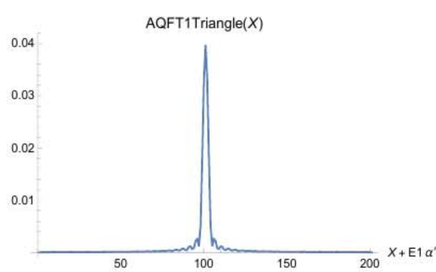


Figure 6: $|A(X)|$ for QFT1, replacing the step wavepacket by the triangular step $\Psi_{B,\mathrm{tri}}$ described in the main text. This wavepacket has a more suppressed power law tail in position space. Comparing this to the second plot in figure (5), we see that $A(X)$ changes consistently with the system scattering on this reduced tail once it takes over from the Gaussian tail of the A and C wavefunctions.

The String Theory results for A(X) are spread out at the expected large scale, and do not track the wavepacket tails.

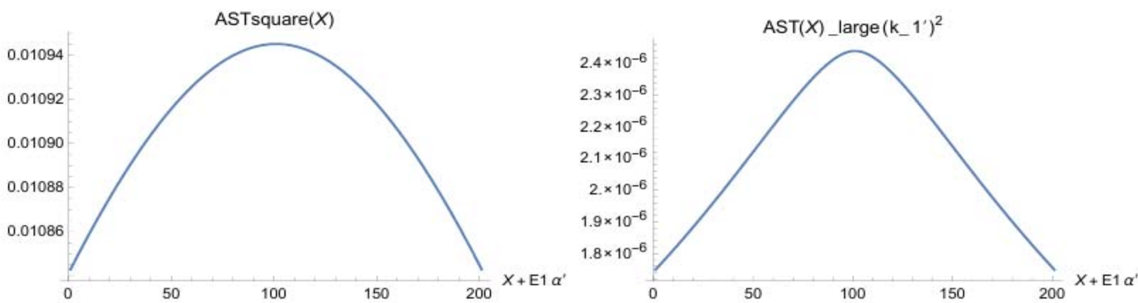


Figure 7: The distribution $A(X)$ for tree-level string theory, plotted for $\sigma = 1/\sqrt{\alpha'}$ and two values of the minimal $k_{1'}^2$, $(1/(10\alpha')$ and $3/\alpha')$. Its width does not increase as we decrease σ, in contrast to QFT, but it does depend on $E/k_{1'}^2$, as expected from $\boxed{2.20}$.

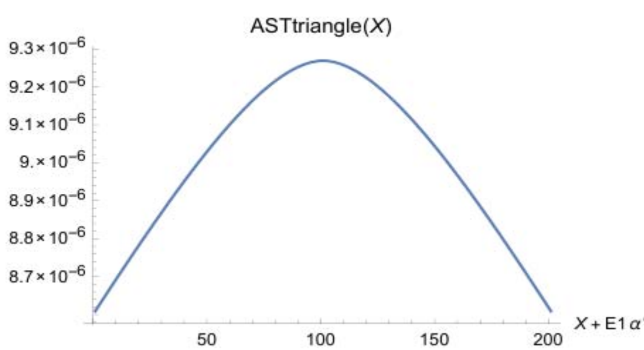


Figure 8: $|A(X)|$ for string theory, replacing the step wavepacket by the triangular step $\Psi_{B,\mathrm{tri}}$ described in the main text. This wavepacket has a more suppressed power law tail in position space. Comparing this to the plots in figure $\boxed{7}$, we see that within the spreading range $X \sim \alpha' E$ depicted, the shape of $A(X)$ is relatively insensitive to the modified B wavefunction tail. (The size of the amplitude is suppressed, reflecting the fact that the support in momentum space is weighted toward larger values of $k_{1'}^2$ in the triangular step wavefunction.) Altogether we find the spread out shape of the string theory amplitude to be insensitive to the variation of the width of the Gaussian wavefunctions for A and C as well as the variation of the B wavefunction. This is consistent with it being dominated by longitudinally spread interaction, rather than scattering locally via the tail of the wavefunction as occurs in the tree level QFT models.

Applied to black hole, this means that a late infalling system can sense early infaller (e.g. matter that formed black hole) in a way that goes beyond effective QFT and GR.

Beyond-GR physics from intrinsic string-theoretic non-locality (open question: what about QCD...?) seems to be able to exploit large energy built up in system of early and late infallers near horizon of black hole. Ongoing work to analyze and apply this to to to thought (and real??) experiments.

There is another subtlety with the low-energy description, known as *dangerous irrelevance*, which comes into early-U cosmology.

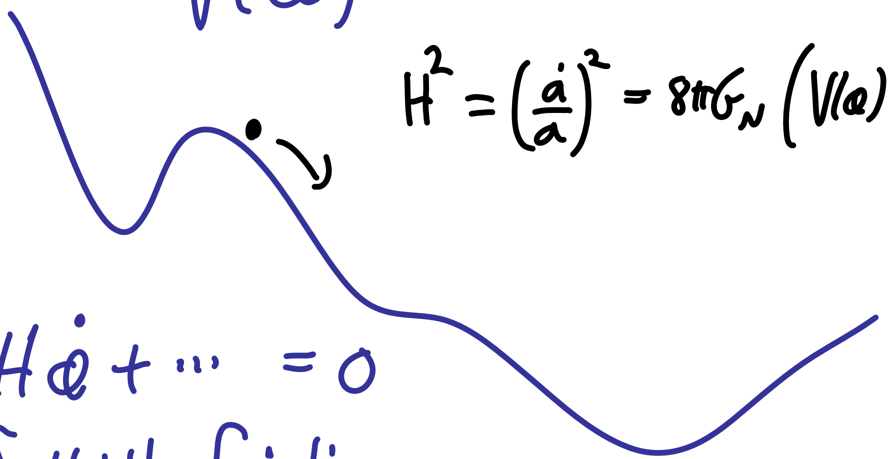
$\mathcal{O}_\Delta$ corrections $\cdot \left(\frac{\text{Energy}}{M_*}\right)^{\Delta-4}$

e.g. $V(\phi) = v_1(\phi-\phi_0) + v_2(\phi-\phi_0)^2$

$$+ \lambda_6 \frac{(\phi-\phi_0)^6}{M_*^2} + \cdots$$

The higher-dimension (`irrelevant') terms can become important if the field ranges over a sufficiently large range in field space. This is automatic in inflationary cosmology:

$V(\phi)$



$$H^2 = \left(\frac{\dot a}{a}\right)^2 = 8\pi G_N \left(V(\phi) + \rho_{\text{kinetic}}(\dot\phi) + \cdots\right)$$

$$\ddot\phi + 3H\dot\phi + \cdots = 0$$

↑ Hubble friction

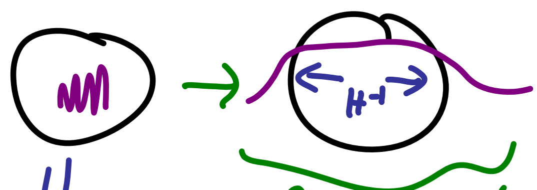
- One can show that $\rho_{kin} \ll V(\phi)$ gives accelerated expansion.
- To solve the flatness, monopole, horizon problems, one needs inflation to last a long time:

$$\frac{a_{end}}{a_{start}} = e^{N_e} \qquad N_e = 60 \text{ for } H \sim 10^{14}\,\text{GeV}$$

Given this, zero-point fluctuations of φ and $g_{\mu\nu}$ (the spacetime geometry) evolve into primordial density perturbations which seed structure:
(cf time-dependent oscillator)

$$\delta\varphi = \int d\vec{k}\, e^{i\vec{k}\cdot\vec{x}}\, \underbrace{\delta\varphi_{\vec{k}}(t)}_{\propto\, \boxed{\frac{e^{i|\vec{k}|t}}{\sqrt{k}}}\ \text{for } \frac{|\vec{k}|}{a} \gg H}$$

$$\delta\ddot{\varphi} + 3H\,\delta\dot{\varphi} + \underbrace{\frac{k^2}{a^2(t)}}_{\substack{\text{physical wavelength} \\ a(t)/k \text{ stretches out}}}\delta\varphi = 0$$



fixed amplitude H

Mukhanov
Bond et al...

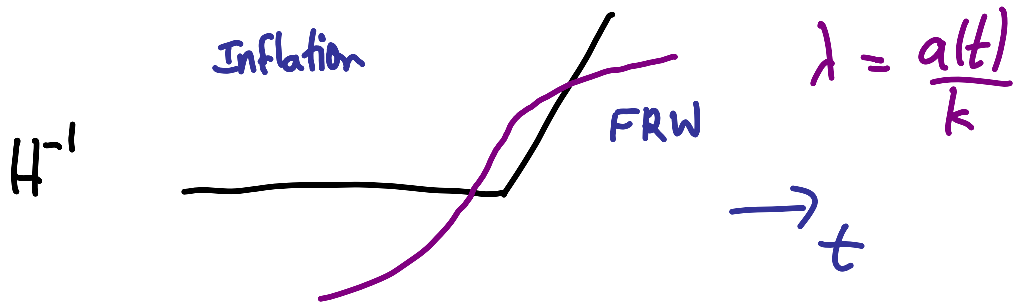
Amplitude: on physical scale L, mode $\delta\varphi_L = X/L^{3/2}$ homogeneous in volume L^3 → $S \sim \frac{1}{2}\int(\dot{X}^2 + \dots)dt$

→ conjugate momentum $P = \dot{X} \sim \Delta X/L$ $\left(\omega \sim \frac{k}{a} \sim \frac{1}{L}\right)$

$$\Delta X\,\Delta P \sim 1 \qquad \frac{(\Delta X)^2}{L} \sim 1 \Rightarrow \Delta X_{min} \sim \sqrt{L},\ \delta\varphi_L \sim \frac{1}{L}$$

$\delta\varphi_L \sim |\delta\varphi_k|\, k^{3/2} \Rightarrow \boxed{|\delta\varphi_k| \sim \frac{1}{a\sqrt{k}}}$. After stretches to $\frac{k}{a} \sim H$, $\boxed{\delta_\varphi(k) \sim |\delta\varphi_k|\, k^{3/2} \sim \frac{k}{a} \sim H}$

When inflation ends, these modes reenter the Hubble horizon during decelerating FRW phase.



Incorporating gravity, normalization of CMB $\Rightarrow$ require

$$P_s \sim \left(\frac{H}{\dot{\varphi}}\right)^2 \langle \delta\varphi^2 \rangle \sim 10^{-10}$$

Also, tensor modes of the graviton are sensitive to scale H:

$$P_{grav} \sim \frac{H^2}{M_p^2} \times \frac{1}{k^3}$$

$$ds^2 = -N^2 dt^2 + h_{ij}\left(dx^i + N^i dt\right)\left(dx^j + N^j dt\right)$$

$$h_{ij} = a^2(t)\left[\underbrace{e^{2\zeta}}_{}\,\delta_{ij} + \underbrace{\gamma_{ij}}_{}\right]$$

$a^2(t)$
$\simeq$
e^{2Ht}

Scalar perturbation
$(\delta\phi = 0$ gauge$)$

remains constant
outside horizon

$\partial_i \gamma_{ij} = 0 = \gamma_{ii}$

tensor perturbations

B modes
Seljak Zaldarriaga ...
Kamionkowski et al
BICEP/SPUD, Spider, ...

Basic Observables

tilt n_s : $P_\zeta = \dfrac{\text{const}}{k^{3+(1-n_s)}}$

tensor/scalar ratio r : $r = \dfrac{P_\gamma}{P_\zeta}$

n_s, r will be observed or constrained at $\simeq 0.01$ level

Non-Gaussianity : $\langle \zeta_{\vec{k}_1} \zeta_{\vec{k}_2} \zeta_{-\vec{k}_1-\vec{k}_2} \rangle$

Bond/Salopek/Komatsu/Spergel, Maldacena ...

Non-Gaussianity: full probability distribution function

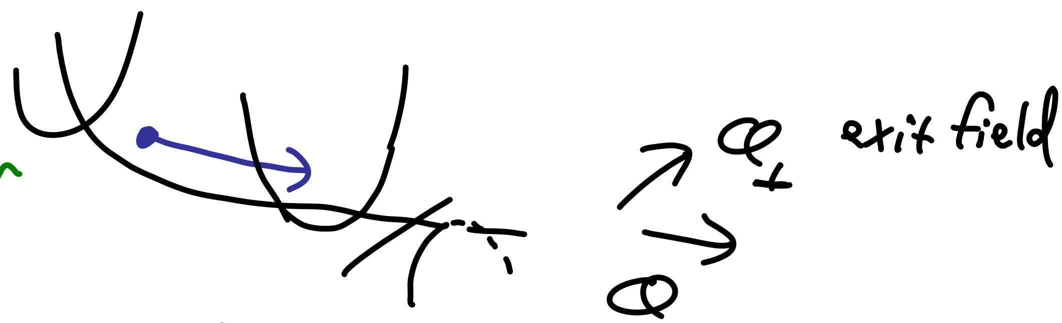
$$P[\zeta(\vec{x}_0)] = |\Psi[\zeta(\vec{x}_0)]|^2$$

This generates N-point correlation functions. Previous searches address certain shapes at low N (=3, sometimes 4). Recently, we found that simple non-adiabatic dynamics involving heavy fields can generate signal/noise that grows with N, requiring analysis of fuller probability distribution. Leading to a new type of NG search.

UV sensitivity of Inflation and dangerous irrelevance.

A seemingly simple way to obtain inflation is to postulate a very flat potential for the inflaton $\varphi(x)$.

Linde '93
hybrid inflation



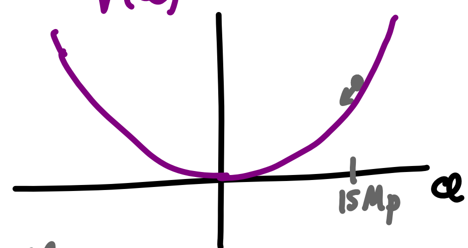
$$\varepsilon \equiv \frac{M_p^2}{2}\left(\frac{V'}{V}\right)^2 << 1 \qquad \eta \equiv M_p^2\left|\frac{V''}{V}\right| << 1$$

However, corrections from the UV physics can generate substructure in $\mathcal{L}(\varphi, \partial\varphi)$:

$$\frac{V(\varphi-\varphi_0)^2}{M_p^2} \rightarrow \Delta\eta \sim 1$$

This UV Sensitivity is greatest in the case of "chaotic inflation" A. Linde '83 where the inflaton ϕ ranges over more than a distance M_p

e.g. $V(\phi) = \frac{1}{2}m^2\phi^2$



$$\left.\begin{matrix} \varepsilon = \frac{1}{2}\left(\frac{V'}{V}M_p\right)^2 \\ \eta = M_p^2\left|\frac{V''}{V}\right| \end{matrix}\right\} \sim \left(\frac{M_p}{\phi}\right)^2 \Rightarrow \phi \sim 15\, M_p$$

In General:

$$\boxed{\star \text{ "Lyth Bound"}: \quad \frac{\Delta\phi}{M_p} \sim \left(\frac{r}{0.01}\right)^{\frac{1}{2}}}$$

UV sensitive if $\gtrsim 1$

observable

→ Control with approximate shift symmetry (Wilsonian 'natural')

UV Sensitivity of Inflation

① Terms of order

$$V \cdot \frac{(\varphi-\varphi_0)^2}{M_p^2} \quad \text{(dimension 6)}$$

in the effective action can ruin inflation

② $$\frac{\Delta\varphi}{M_p} \simeq r^{\frac{1}{2}} \frac{N_e}{\sqrt{24}} \quad \text{(Lyth)}$$

GUT-scale inflation (with observable tensor modes) $\Leftrightarrow$ $\Delta\varphi > M_p$

③ General Single-field inflation involves higher derivative terms which affect solution & perturbations

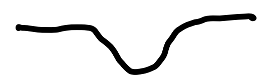
cf Non-Gaussianity

④ $g^2\varphi^2\chi^2$ couplings $\Rightarrow$ temporarily light fields affect evolution. . . .

(5a) (mass $> H$) Dong et al '10 'Flattening'

<u>Heavy</u> fields affect results: they adjust in response to inflationary potential energy. QFT toy model

$$V(\varphi_L, \varphi_H) = g^2 \varphi_L^2 \varphi_H^2 + m^2 (\varphi_H - \varphi_0)^2$$



$$\frac{\partial V}{\partial \varphi_H} \equiv 0 \Rightarrow V = \frac{g^2 \varphi_L^2}{g^2 \varphi_L^2 + m^2} m^2 \varphi_0^2$$

($\dot{\varphi}_H^2$ term subdominant)

flatter: energetically favorable.

- UV completion of gravity (e.g. string theory) can introduce ϕ_H (e.g. 'moduli' scalar fields).

↳ $V \propto \phi^n \to V \propto \phi^{p<n}$ in examples.

(5b) Time-dependent masses (from coupling to inflaton) can affect perturbations even if <u>minimal</u> mass $\underline{\mu} \gg H$:

w/Flauger, Mirbabayi, Senatore

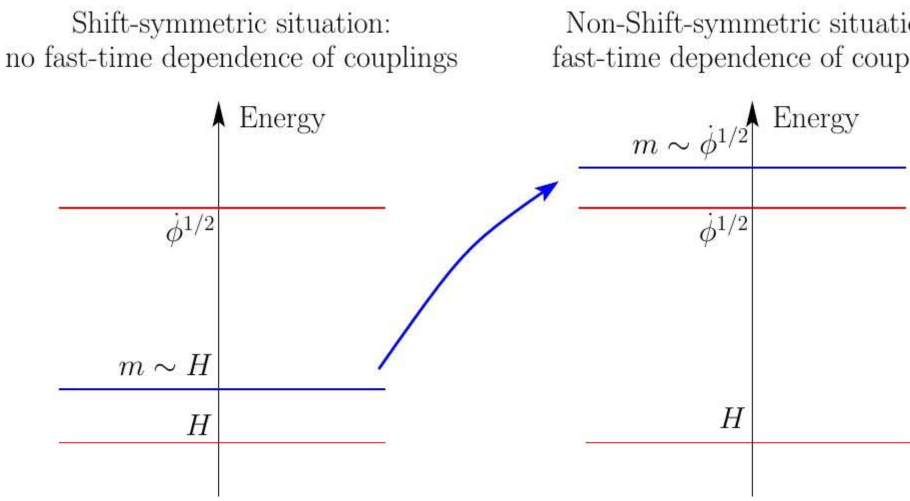
$$\left[g^2(\phi-\phi_n)^2+\mu^2\right]\chi_n^2$$

$$\Rightarrow \langle n_\chi \rangle \sim (g\dot{\phi})^{3/2} e^{-\frac{\pi\mu^2}{g\dot{\phi}}}$$

sources $\langle\zeta\zeta\rangle$, $\langle\zeta\zeta\zeta\rangle$, ...

$\langle\gamma\gamma\rangle$, ...

Relevant for refining searches for oscillatory (non-)Gaussianity



→ $O(.1-1)$ couplings to inflaton (preserving inflation overall)

Figure 1: Pictorial representation of our findings: in an inflationary theory with an approximate continuous shift symmetry for the inflaton, only particles that are not much heavier than the Hubble scale H are relevant for the dynamics of the fluctuations. However, as we will see, if the continuous shift symmetry is broken, e.g. to a discrete shift symmetry, heavier particles can become relevant as depicted on the right. In the scenarios studied in this work, the new scale is set by $\dot\phi$. The basic estimate $\exp(-\pi m^2/\dot\phi) \sim 1/\sqrt{N_{\rm modes}}$ suggests observational sensitivity to these massive particles, which we confirm in a detailed analysis.

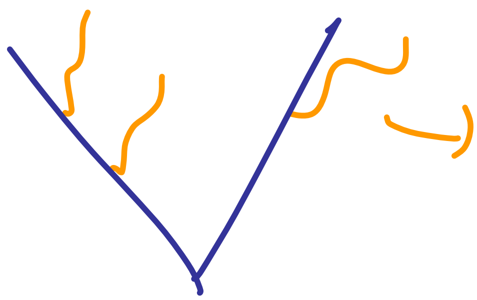
EFT of inflationary perturbations contains arbitrary functions of t even at single-field level.

Precision of data means we can't safely integrate out heavy ($m_\chi \gtrsim \dot\phi^{\frac{1}{2}}$) fields

Inflaton ϕ coupled to

heavy $\chi \rightarrow m_\chi(\phi(t))$

time-dependent



$\rightarrow$ $\delta\phi$ radiation

non-adiabatic χ production

Kofman Linde
Starobinski Traschen
Brandenberger,
Chung et al, Green
Horn ES Senatore...

$$\left(\frac{\dot\phi^{1/2}}{H}\right)^3 e^{-\frac{\pi\mu^2}{\dot\phi}} \sim \frac{1}{\sqrt{N_{pixel}}} \sim 10^{-3}$$

sensitivity to heavy χ fields

Novel shape $\langle \delta\phi_{k_1} \cdots \delta\phi_{k_N} \rangle$

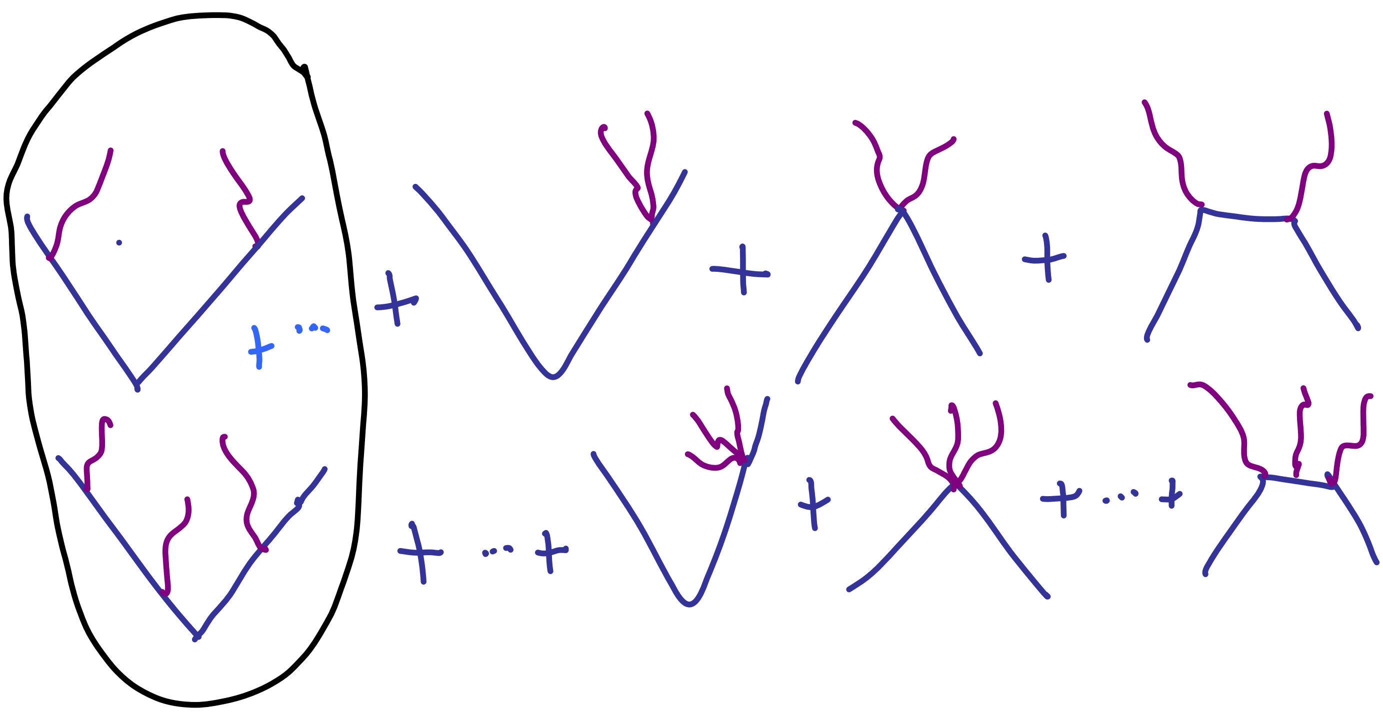
and $(S/N)_{N+1} / (S/N)_N > 1$ possible for a range of N!

General Calculation

$$\langle in|\bar{T}\exp(i\int_{-\infty(1+i\epsilon)}^{t} dt_1\mathcal{H}_{int})\delta\phi_{k_1}(t)\ldots\delta\phi_{k_N}(t)T\exp(-i\int_{-\infty(1-i\epsilon)}^{t} dt_2\mathcal{H}_{int})|in\rangle$$

$$\mathcal{H}_{int} = \chi^2 m_\chi^2(\phi_0(t)+\delta\phi)$$

$$|\mathcal{N}|^2\langle out|e^{\int_q \frac{\beta_q^*}{2\alpha_q^*}a_q^2}\bar{T}\exp(i\int_{-\infty(1+i\epsilon)}^{t} dt_1\mathcal{H}_{int})\delta\phi_{k_1}(t)\ldots\delta\phi_{k_N}(t)T\exp(-i\int_{-\infty(1-i\epsilon)}^{t} dt_2\mathcal{H}_{int})e^{\int_k \frac{\beta_k}{2\alpha_k}a_k^{\dagger 2}}|out\rangle$$



Novel shape for $\mathcal{L}_3 \sim \delta\phi\,\chi\chi\sin\omega t$

Probability Distribution

In general this is the functional

$$P[\delta\phi^0(\mathbf{x})] = \int D\chi^0 |\Psi[\delta\phi^0(\mathbf{x}), \chi^0(\mathbf{x})]|^2$$

$$\Psi[\delta\phi^0(\mathbf{x}), \chi^0(\mathbf{x})] = \int D\delta\phi(\mathbf{x}, t)|_{\delta\phi(t=t_C)=\delta\phi^0(\mathbf{x})} D\chi(\mathbf{x}, t)|_{\chi(t=t_C)=\chi^0(\mathbf{x})} e^{i\mathcal{S}[\delta\phi(\mathbf{x},t),\chi(\mathbf{x},t)]}$$

The histogram of temperature fluctuations in the map would in general be given by

$$N_{\delta\hat{\phi}} = \int d\mathbf{x}' \int D\delta\phi^0(\mathbf{x}) P[\delta\phi^0(\mathbf{x})] H\delta(\delta\phi(\mathbf{x}') - \delta\hat{\phi})$$

This can introduce non-Gaussian features that we may be able to constrain in a more model-independent way than searching for specific N-point functions.

For factorized contributions

$$N_{\delta\hat{\phi}} \sim \delta(0)\bar{n}_\chi \int d\mathbf{x}' \sum_n \eta_n^{-3}\delta(\tilde{H}_n(\mathbf{x}') - \frac{\delta\hat{\phi}}{H})$$

$$\tilde{H}_n(\mathbf{y}) = \int d\mathbf{k}_j e^{i\mathbf{k}_j\cdot\mathbf{y}} \frac{\hat{h}(k_j\eta_n)}{k_j^3}$$

(then convolved with Gaussian).
For one Fourier mode of mass,

$$N_{\delta\hat{\phi}} \simeq 4\pi\delta(0)\bar{n}_\chi N_e^{data} \frac{\omega}{2\pi H} \frac{1}{\sqrt{1 - (\frac{\delta\hat{\phi}}{H})^2(\frac{\omega}{Hc_b})^2}}$$

Moreover, we can implement this for general $m_\chi(\phi)$ and get a much more model independent constraint.

String theory and Inflation

Despite the complications of the landscape of string compactifications, clear mechanisms for inflation emerge, some of which are falsifiable based on CMB data (e.g. primordial gravitational waves or shapes of non-Gaussianity).

These led to a more systematic low energy effective field theory analysis of inflation and signatures, helping determine which features follow from low energy considerations and which don't.

A rare opportunity to do some traditional science with String theory

Variety of inflationary mechanisms in string theory

Many contributors (Baumann/McAllister book)

$V(\phi)$ flat

$\Delta\phi \le M_p$

Accidental 1% cancellation

KKLMMT, Roulette, LARGE Fibre

[cosmic strings]

Self-interactions slowing ϕ

$V(\phi)$ steep

e.g. DBI

λ

Trapped ...

NG

$V(\phi)$ flat, $\Delta\phi > M_p$

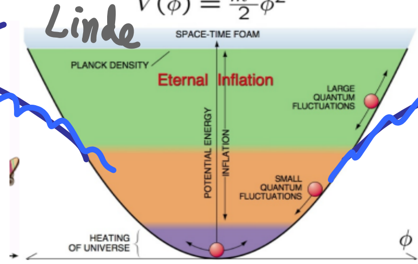
Monodromy

L (impact parameter)

Linde

$p_0 = 2$

$p<2$



N-axion Alignment

Gauge fields, ...

[drifting $\boxed{r}$, oscillations]

I'll focus on Axion Monodromy, but other scenarios also interesting

- KKLMMT / DBI ~ illustrates range of inflation & Non-Gaussianity; Relativity on field space

  (KKLMMT: pioneered careful assessment of Planck-suppressed contributions)

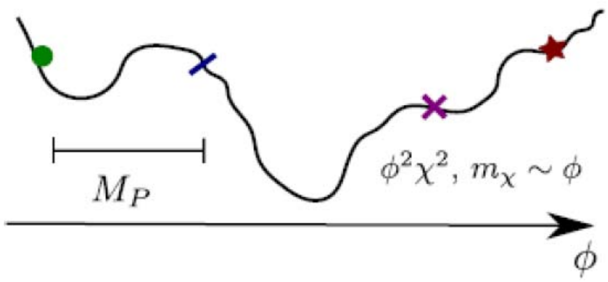
- Fibre, Roulette, ...

  (Fibre: Exponential potential ⟷ Starobinsky)

- Multiple-field effects, connections to 'weak gravity conjecture' etc.

- ...

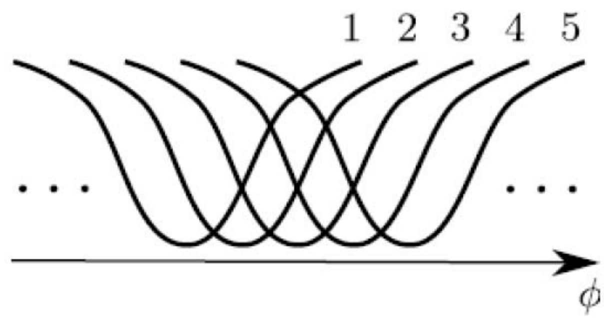
Parameterized ignorance of quantum grav.



New degrees of freedom each $\Delta\Phi \sim M_P$

No continuous global symm. in QG

String Theory axions (and duals)



From ubiquitous Axion-Flux couplings

Discrete shift symm., f<<M_p

[cf Chaotic Infl.(Linde), Natural Infl. (Freese et al)]

$$\int d^D x \sqrt{G} \sum_q \Big| \overbrace{F_q - C_{q-3} \wedge H + F_{q-2n}\, B \wedge \cdots \wedge B}^{\tilde{F}_q \text{ Gauge-invar.}} \Big|^2$$

This generalizes Stueckelburg couplings in electromagnetism

$$S = \int d^4x \left\{ F^2 - \rho^2 (\partial\theta - A)^2 \right\}$$

Gauge symmetry $\quad A_\mu \to A_\mu + \partial_\mu \Lambda$
$\qquad\qquad\qquad\qquad\; \theta \to \theta - \Lambda$

In string theory, the string sources a <u>2-index</u> gauge potential B_{MN} analogously to how a charged particle sources A_μ in Electromagnetism

axions $= B_{MN}$ - modes
(and duals)

Is there a corresponding unbroken phase?

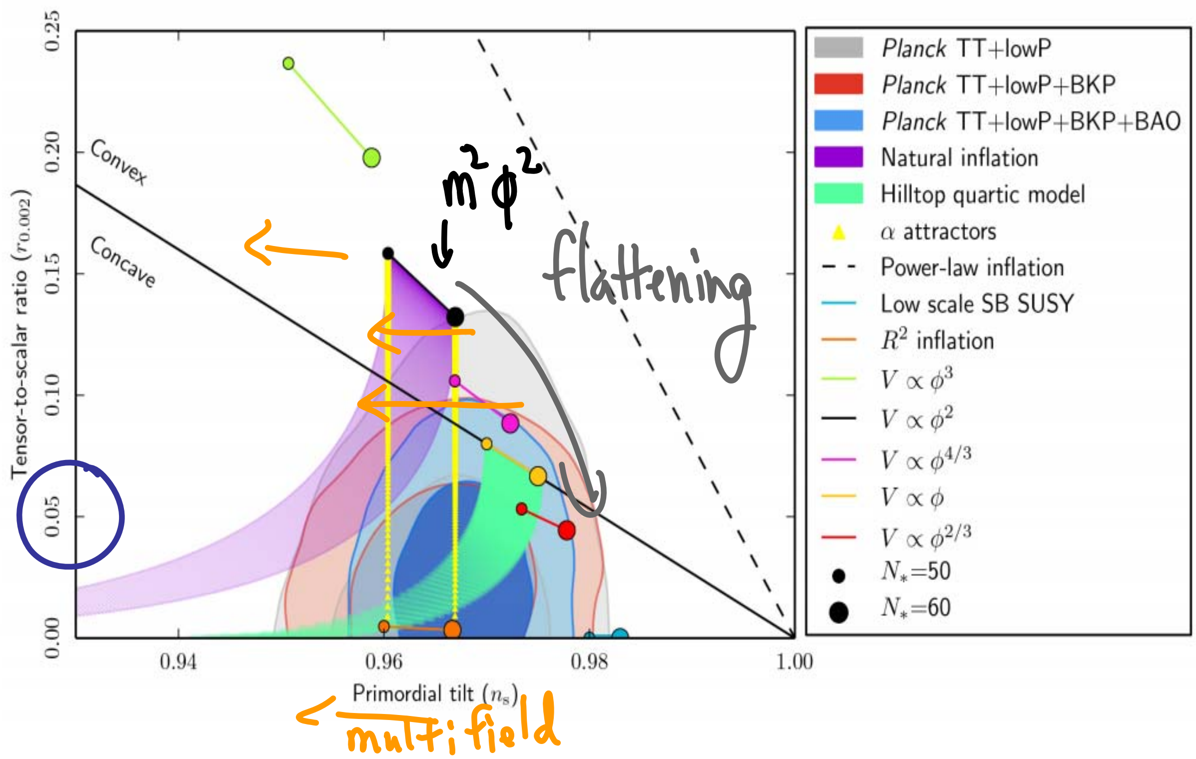


Fig. 54. Marginalized joint 68 % and 95 % CL regions for n_s and $r_{0.002}$ from *Planck* alone and in combination with its cross-correlation with BICEP2/Keck Array and/or BAO data compared with the theoretical predictions of selected inflationary models.

Axion monodromy remains viable only because of the flattening effect. (Tested soon by r)

Constraints are tightening...

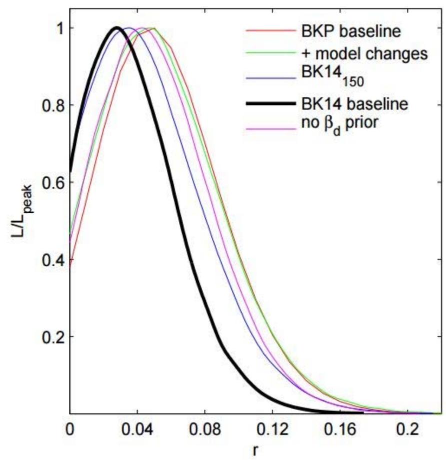


FIG. 5. Likelihood results on r for several intermediate steps between the BKP (previous) and BK14 (current) analyses. See text for details.

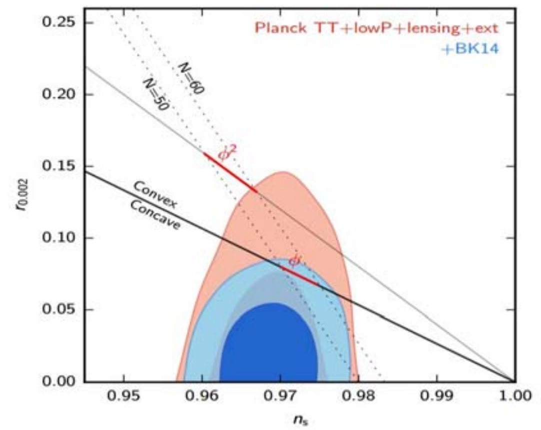


FIG. 7. Constraints in the r vs. n_s plane when using *Planck* plus additional data, and when also adding BICEP2/*Keck* data through the end of the 2014 season including new 95 GHz maps—the constraint on r tightens from $r_{0.05} < 0.12$ to $r_{0.05} < 0.07$. This figure is adapted from Fig. 21 of Ref. [2]—see there for further details.

Multiple projects aiming to detect or constrain this as well as observing interesting astrophysical sources, include:

(Planck+) BICEP/Keck, SPIDER, ACT, SPT, CMBS4, LiteBird,...

Other projects going for large-scale structure probes of cosmology, e.g. Sphere-X proposal

Inflationary cosmology and related CMB (and large-scale structure) horizon observables are formally `UV sensitive', requiring modeling with control of QG corrections (especially with respect to primordial GW).

Ongoing program to analyze and test string theoretic inflation mechanisms, although can't turn it around given limited data.

Regardless of their fate, this has led to much more systematic understanding of inflationary theory and motivates specific searches for new physics.

Summary:

Horizons are predicted by general relativity and near-horizon physics is seen in some observations.

Even though weakly curved, they exhibit sensitivity to physics beyond GR and quantum field theory, because of long-time processes and, the `dangerous irrelevance' of large field excursions, and their microphysical entropy counts.

This is a fruitful playground for thought experiments and even some real ones.