Magnetogenesis in non-trivial inflationary scenarios

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Plan of the talk

1. Observational evidence for magnetic fields

2. Essential aspects of inflation

3. Generation of magnetic fields in de Sitter inflation

4. Generation of magnetic fields in slow roll inflation

5. Challenges in inflationary models leading to features

6. Circumventing the challenges using two field models

7. Imprints of magnetic fields on the CMB

8. Summary

This talk is based on. . .

◆ S. Tripathy, D. Chowdhury, R. K. Jain and L. Sriramkumar, *Challenges in the choice of the nonconformal coupling function in inflationary magnetogenesis*, Phys. Rev. D **105**, 063519 (2022) [arXiv:2111.01478 [astro-ph.CO]].

◆ S. Tripathy, D. Chowdhury, H. V. Ragavendra, R. K. Jain and L. Sriramkumar, *Circumventing the challenges in the choice of the non-conformal coupling function in inflationary magnetogenesis*, Phys. Rev. D **107**, 043501 (2023) [arXiv:2211.05834 [astro-ph.CO]].

Plan of the talk

Observational evidence for magnetic fields

Magnetic fields are ubiquitous in the universe. They are observed at different strengths over a wide range of scales.

◆ In galaxies, the strength of the observed magnetic fields is $\mathcal{O}(10^{-6}\,\mathrm{G})$, which is coherent over scales of 1–$10\,\mathrm{Kpc}$[1].

◆ In clusters of galaxies, the strength of the magnetic fields is $\mathcal{O}(10^{-7}$–$10^{-6}\,\mathrm{G})$ with a coherent length of $10\,\mathrm{Kpc}$–$1\,\mathrm{Mpc}$[2].

◆ In the intergalactic medium, the strength of the magnetic fields is greater than $10^{-16}\,\mathrm{G}$, which is coherent on scales above $1\,\mathrm{Mpc}$[3].

◆ The observations of the anisotropies in the cosmic microwave background (CMB) constrain the magnetic fields at the scale of $1\,\mathrm{Mpc}$ to be less than $10^{-9}\,\mathrm{G}$[4].

While astrophysical processes may be sufficient to explain the origin of magnetic fields in galaxies and clusters of galaxies, *one may have to turn to cosmological phenomena to explain the magnetic fields observed in voids*.

[1] R. Beck, Space Sci. Rev. **99**, 243 (2001).

[2] See, for instance, T. E. Clarke, P. P. Kronberg and H. Böhringer, Astrophys. J. **547**, L111 (2001).

[3] A. Neronov and I. Vovk, Science **328**, 73 (2010).

[4] Planck Collaboration (P. A. R. Ade *et al.*), Astron. Astrophys. **594**, A19 (2016).
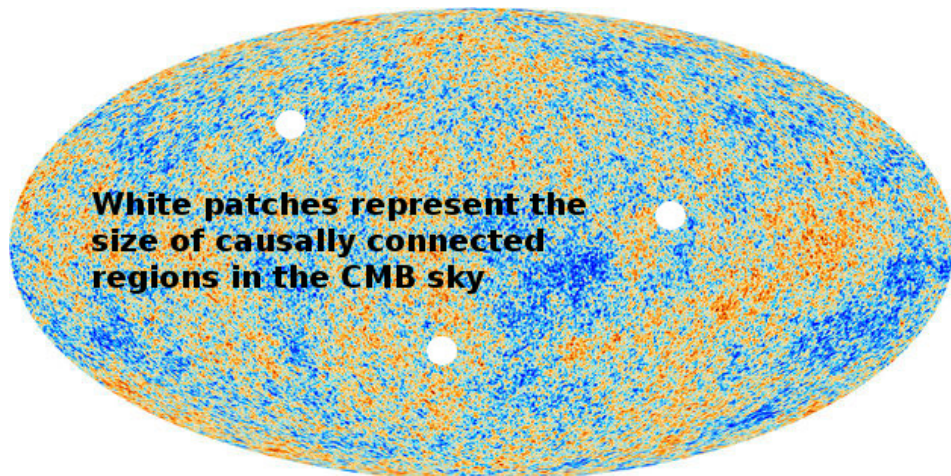
Plan of the talk

The horizon problem

White patches represent the size of causally connected regions in the CMB sky
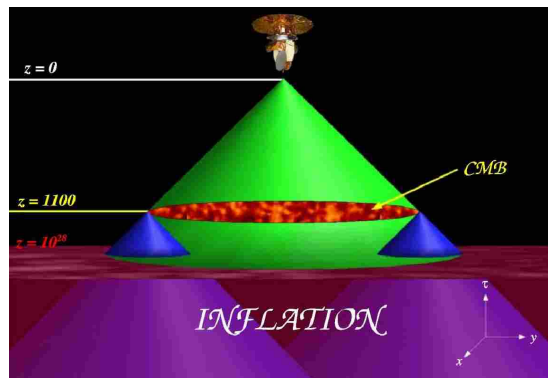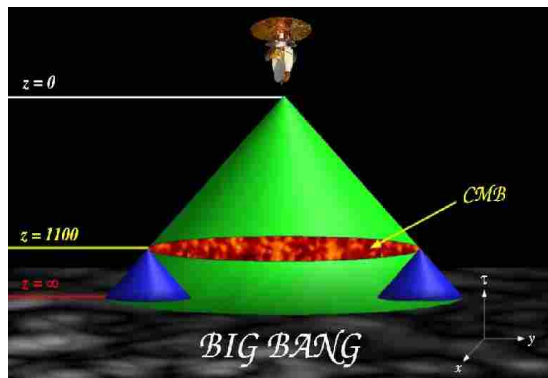
The radiation from the CMB arriving at us from regions separated by more than the Hubble radius at the last scattering surface, which subtends an angle of about $1°$ today, could not have interacted before decoupling.

The resolution of the horizon problem in the inflationary scenario

Another illustration of the horizon problem (on the left), and an illustration of its resolution (on the right) through an early and sufficiently long epoch of inflation[5].

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[5]Images from W. Kinney, astro-ph/0301448.

Behavior of the comoving wave numbers and Hubble radius
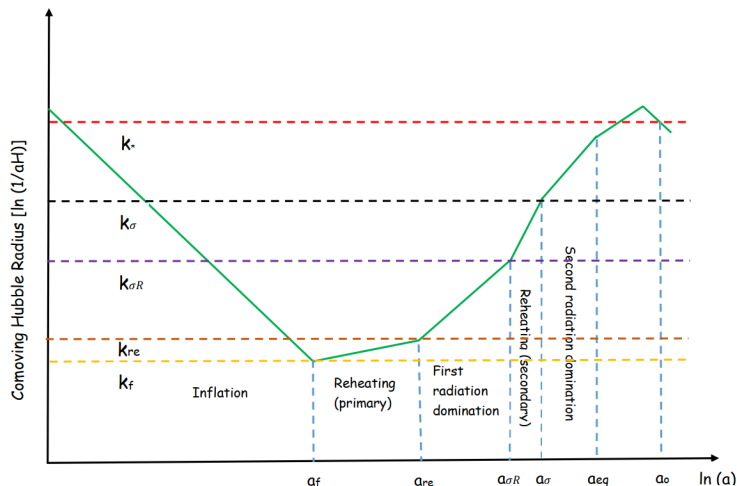
Behavior of the comoving wave number k (horizontal lines in different colors) and the comoving Hubble radius $d_{\mathrm{H}}/a = (a\,H)^{-1}$ (in green) across different epochs[6].

[6] Md. R. Haque, D. Maity, T. Paul and L. Sriramkumar, Phys. Rev. D **104**, 063513 (2021).

The time and duration of inflation
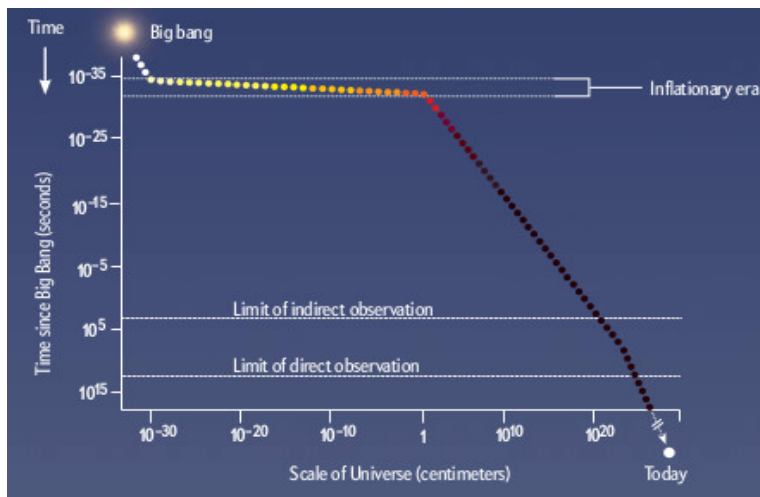
Inflation – a brief period of accelerated expansion – is expected to have taken place during the very stages of the universe[7].

[7]Image from P. J. Steinhardt, Sci. Am. **304**, 18 (2011).

Driving inflation with scalar fields

Inflation can be achieved with scalar fields encountered in high energy physics[8].

[8]Image from P. J. Steinhardt, Sci. Am. **304**, 34 (2011).

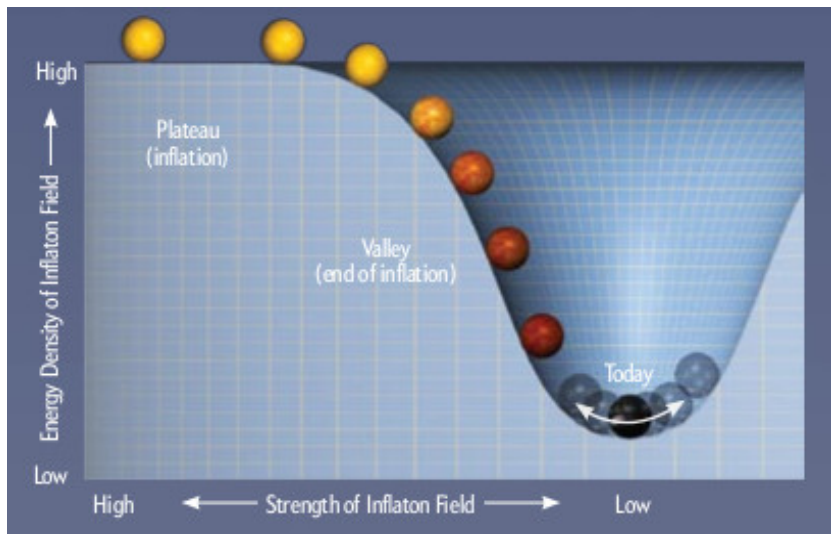# Variety of potentials can drive inflation

Large field
$$V(\phi) = \Lambda^4 (\phi/\mu)^p$$
$$V(\phi) = \Lambda^4 e^{\phi/\mu}$$

Small field
$$V(\phi) = \Lambda^4 [1 - (\phi/\mu)^p]$$
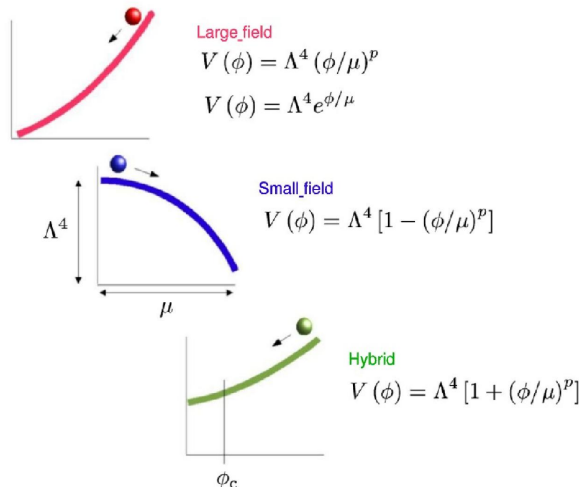
Hybrid
$$V(\phi) = \Lambda^4 [1 + (\phi/\mu)^p]$$

A variety of scalar field potentials have been considered to drive inflation[9]. Often, these potentials are classified as small field, large field and hybrid models.

[9]Image from W. Kinney, astro-ph/0301448.

Scalar spectral index and the tensor-to-scalar ratio

While comparing with the observations, for convenience, one often uses the following power law, template scalar and the tensor spectra:

$$\mathcal{P}_{\mathrm{S}}(k) = A_{\mathrm{S}} \left(\frac{k}{k_*}\right)^{n_{\mathrm{S}}-1}, \qquad \mathcal{P}_{\mathrm{T}}(k) = A_{\mathrm{T}} \left(\frac{k}{k_*}\right)^{n_{\mathrm{T}}},$$

with the spectral indices n_{S} and n_{T} assumed to be constant. The tensor-to-scalar ratio r is defined as

$$r(k) = \frac{\mathcal{P}_{\mathrm{T}}(k)}{\mathcal{P}_{\mathrm{S}}(k)}.$$

CMB angular power spectrum from Planck
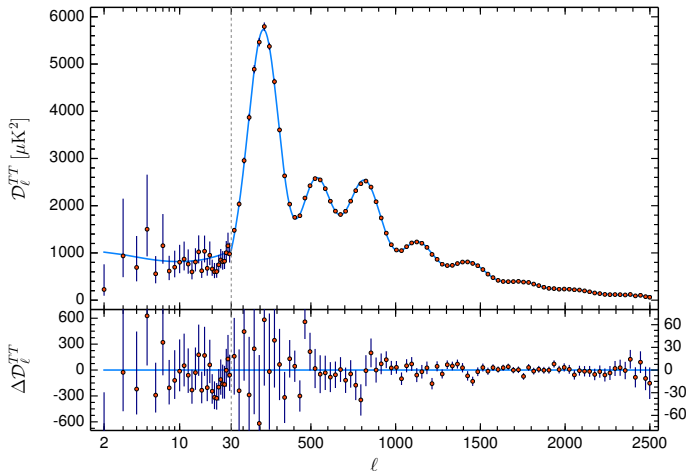
The CMB TT angular power spectrum from the Planck 2018 data (red dots with error bars) and the best fit ΛCDM model with a power law primordial spectrum (solid blue curve)[10]

[10] Planck Collaboration (N. Aghanim *et al.*), Astron. Astrophys. **641**, A6 (2020).

Performance of inflationary models in the n_s-r plane
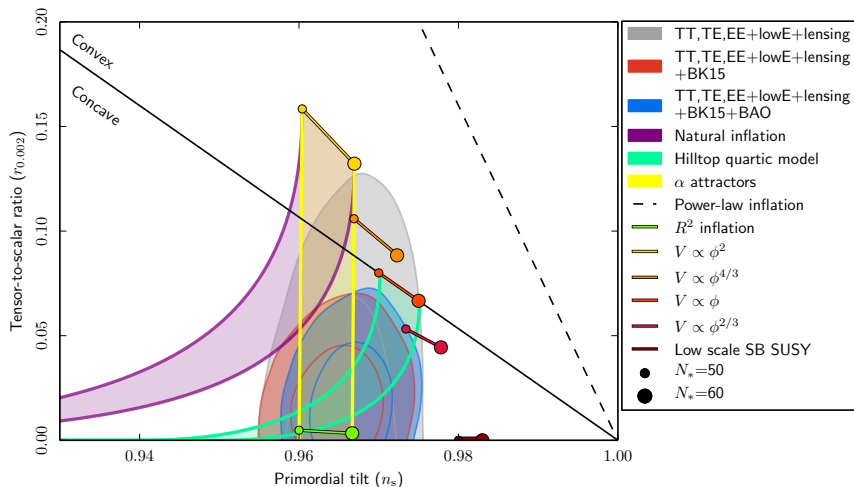
Joint constraints on n_s and $r_{0.002}$ from Planck in combination with other data sets, compared to the theoretical predictions of some of the popular inflationary models[11].

[11] Planck Collaboration (Y. Akrami *et al.*), Astron. Astrophys. **641**, A10 (2020).

Plan of the talk

Non-conformally coupled electromagnetic fields

We shall assume that the electromagnetic field is described by the action[12]

$$S[A^\mu] = -\frac{1}{16\,\pi} \int \mathrm{d}^4x \sqrt{-g}\, J^2(\phi)\, F_{\mu\nu}\, F^{\mu\nu},$$

where $J(\phi)$ denotes the coupling function and the field tensor $F_{\mu\nu}$ is expressed in terms of the vector potential A_μ as $F_{\mu\nu} = (\partial_\mu\, A_\nu - \partial_\nu\, A_\mu)$.

On working in the Coulomb gauge wherein $A_\eta = 0$ and $\partial_i\, A^i = 0$, one finds that the Fourier modes, say, $\bar{A}_k$, describing the vector potential satisfy the differential equation

$$\bar{A}_k'' + 2\,\frac{J'}{J}\,\bar{A}_k' + k^2 \bar{A}_k = 0.$$

If we write $\bar{A}_k = \mathcal{A}_k/J$, then this equation reduces to

$$\mathcal{A}_k'' + \left(k^2 - \frac{J''}{J}\right)\,\mathcal{A}_k = 0.$$

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[12]See, for instance, J. Martin and J. Yokoyama, JCAP **01**, 025 (2008);
K. Subramanian, Astron. Nachr. **331**, 110 (2010).

Quantization and power spectra of the electromagnetic field

For each comoving wave vector $\boldsymbol{k}$, we can define the right-handed orthonormal basis $(\hat{\varepsilon}_1^{\boldsymbol{k}}, \hat{\varepsilon}_2^{\boldsymbol{k}}, \hat{\boldsymbol{k}})$ which satisfy the relations

$$\hat{\varepsilon}_1^{\boldsymbol{k}} \cdot \hat{\varepsilon}_1^{\boldsymbol{k}} = \hat{\varepsilon}_2^{\boldsymbol{k}} \cdot \hat{\varepsilon}_2^{\boldsymbol{k}} = 1, \, \hat{\varepsilon}_1^{\boldsymbol{k}} \cdot \hat{\varepsilon}_2^{\boldsymbol{k}} = 0, \, \hat{\varepsilon}_1^{\boldsymbol{k}} \times \hat{\varepsilon}_2^{\boldsymbol{k}} = \hat{\boldsymbol{k}}, \, \hat{\boldsymbol{k}} \cdot \hat{\varepsilon}_1^{\boldsymbol{k}} = \hat{\boldsymbol{k}} \cdot \hat{\varepsilon}_2^{\boldsymbol{k}} = 0, \, \hat{\varepsilon}_1^{-\boldsymbol{k}} = -\hat{\varepsilon}_1^{\boldsymbol{k}}, \, \hat{\varepsilon}_2^{-\boldsymbol{k}} = \hat{\varepsilon}_2^{\boldsymbol{k}}.$$

On quantization, the vector potential $\hat{A}_i$ can be decomposed as follows:

$$\hat{A}_i(\eta, \boldsymbol{x}) = \sqrt{4\pi} \int \frac{\mathrm{d}^3 \boldsymbol{k}}{(2\pi)^{3/2}} \sum_{\lambda=1}^{2} \hat{\varepsilon}_{\lambda i}^{\boldsymbol{k}} \left[\hat{b}_{\boldsymbol{k}}^{\lambda} \, \bar{A}_k^{\lambda}(\eta) \, \mathrm{e}^{i\,\boldsymbol{k}\cdot\boldsymbol{x}} + \hat{b}_{\boldsymbol{k}}^{\lambda\dagger} \, \bar{A}_k^{\lambda *}(\eta) \, \mathrm{e}^{-i\,\boldsymbol{k}\cdot\boldsymbol{x}} \right],$$

where $\hat{\varepsilon}_{\lambda i}^{\boldsymbol{k}}$ represent components of the polarization vectors and the summation is over to the two orthonormal transverse polarizations.

The power spectra associated with the magnetic and electric fields are defined to be[13]

$$\mathcal{P}_{\mathrm{B}}(k) = \frac{\mathrm{d}\langle 0|\hat{\rho}_{\mathrm{B}}|0\rangle}{\mathrm{d}\ln k} = \frac{J^2(\eta)}{2\pi^2} \frac{k^5}{a^4(\eta)} |\bar{A}_k(\eta)|^2 = \frac{1}{2\pi^2} \frac{k^5}{a^4(\eta)} |\mathcal{A}_k(\eta)|^2,$$

$$\mathcal{P}_{\mathrm{E}}(k) = \frac{\mathrm{d}\langle 0|\hat{\rho}_{\mathrm{E}}|0\rangle}{\mathrm{d}\ln k} = \frac{J^2(\eta)}{2\pi^2} \frac{k^3}{a^4(\eta)} |\bar{A}_k'(\eta)|^2 = \frac{1}{2\pi^2} \frac{k^3}{a^4(\eta)} \left| \mathcal{A}_k'(\eta) - \frac{J'(\eta)}{J(\eta)} \mathcal{A}_k(\eta) \right|^2.$$

[13] J. Martin and J. Yokoyama, JCAP **01**, 025 (2008); K. Subramanian, Astron. Nachr. **331**, 110 (2010).

Initial conditions and normal modes in de Sitter inflation

Let us now consider de Sitter inflation wherein the scale factor is given by $a(\eta) = -1/(H_{\mathrm{I}}\,\eta)$, with H_{I} denoting the constant Hubble parameter.

Typically, the non-conformal coupling function J is assumed to depend on the scale factor as follows:

$$J(\eta) = \left[\frac{a(\eta)}{a(\eta_{\mathrm{e}})}\right]^n = \left(\frac{\eta}{\eta_{\mathrm{e}}}\right)^{-n},$$

where η_{e} denotes the conformal time at the end of inflation.

The Bunch-Davies initial conditions on the electromagnetic modes $\mathcal{A}_k$ can be imposed in the limit $k \gg \sqrt{J''/J}$, which, for the above coupling function, corresponds to the modes being in the sub-Hubble domain at early times.

For the above coupling function, the solution that satisfies the Bunch-Davies initial conditions is given by

$$\mathcal{A}_k(\eta) = \sqrt{-\frac{\pi\,\eta}{4}}\,\mathrm{e}^{i\,(n+1)\,\pi/2}\,H_\nu^{(1)}(-k\,\eta),$$

where $\nu = n + (1/2)$ and $H_\nu^{(1)}(z)$ denotes the Hankel function of the first kind.

Power spectra of electromagnetic fields in de Sitter inflation

The spectra of the electromagnetic fields can be evaluated in the limit $k \ll \sqrt{J''/J}$, which corresponds to the super-Hubble limit for our choice of the coupling function.

In the limit $(-k\,\eta_{\mathrm{e}}) \ll 1$, the spectra of the magnetic and electric fields $\mathcal{P}_{\mathrm{B}}(k)$ and $\mathcal{P}_{\mathrm{E}}(k)$ can be obtained to be[14]

$$\mathcal{P}_{\mathrm{B}}(k) = \frac{H_{\mathrm{I}}^4}{8\,\pi}\,\mathcal{F}(m)\,(-k\,\eta_{\mathrm{e}})^{2\,m+6}, \quad \mathcal{P}_{\mathrm{E}}(k) = \frac{H_{\mathrm{I}}^4}{8\,\pi}\,\mathcal{G}(m)\,(-k\,\eta_{\mathrm{e}})^{2\,m+4},$$

where we the quantities $\mathcal{F}(m)$ and $\mathcal{G}(m)$ are given by

$$\mathcal{F}(m) = \frac{1}{2^{2\,m+1}\cos^2(m\,\pi)\,\Gamma^2(m+3/2)}, \quad \mathcal{G}(m) = \frac{1}{2^{2\,m-1}\cos^2(m\,\pi)\,\Gamma^2(m+1/2)},$$

with

$m = n$ for $n < -1/2$ and $m = -(n+1)$ for $n > -1/2$, in the case of $\mathcal{P}_{\mathrm{B}}(k)$,

$m = n$ for $n < 1/2$ and $m = -(n-1)$ for $n > 1/2$, in the case of $\mathcal{P}_{\mathrm{E}}(k)$.

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[14]See, for instance, J. Martin and J. Yokoyama, JCAP **01**, 025 (2008);
K. Subramanian, Astron. Nachr. **331**, 110 (2010).

Scale invariant spectra of magnetic fields and backreaction

The spectral indices for the magnetic and electric fields, say, n_{B} and n_{E}, can be written as

$$n_{\mathrm{B}} = (2\,n + 6) \text{ for } n < -1/2, \quad n_{\mathrm{B}} = (4 - 2\,n) \text{ for } n > -1/2,$$
$$n_{\mathrm{E}} = (2\,n + 4) \text{ for } n < 1/2, \quad\ \ n_{\mathrm{E}} = (6 - 2\,n) \text{ for } n > 1/2.$$

To be consistent with observations, the spectrum of the magnetic field is expected to be nearly scale invariant and, evidently, this is possible when $n \simeq -3$ or when $n \simeq 2$. In these cases, it is clear that $n_{\mathrm{E}} \simeq -2$ and $n_{\mathrm{E}} \simeq 2$, respectively.

At late times, $n_{\mathrm{E}} \simeq -2$ implies that the energy density in the electric field is significant leading to a large backreaction.

In order to avoid such an issue, one often considers the $n = 2$ case to lead to a scale invariant magnetic field with negligible backreaction due to the electric field. Note that, in this case, the power spectra reduce to the following simple forms (with $k_{\mathrm{e}} = -1/\eta_{\mathrm{e}}$):

$$\mathcal{P}_{\mathrm{B}}(k) = \frac{9\,H_{\mathrm{I}}^4}{4\,\pi^2}, \quad \mathcal{P}_{\mathrm{E}}(k) = \frac{H_{\mathrm{I}}^4}{4\,\pi^2}\,\left(\frac{k}{k_{\mathrm{e}}}\right)^2.$$

Helical electromagnetic fields

The action we had considered earlier can be extended to include a parity violating term as follows[15]:

$$S[A^\mu] = -\frac{1}{16\,\pi} \int d^4x \sqrt{-g} \left[J^2(\phi)\, F_{\mu\nu}\, F^{\mu\nu} - \frac{\gamma}{2}\, I^2(\phi)\, F_{\mu\nu}\, \widetilde{F}^{\mu\nu} \right],$$

where $\widetilde{F}^{\mu\nu} = (\epsilon^{\mu\nu\alpha\beta}/\sqrt{-g})\, F_{\alpha\beta}$, with $\epsilon^{\mu\nu\alpha\beta}$ being the completely anti-symmetric Levi-Civita tensor, and γ is a constant.

In such a case, we can define the orthonormal basis of vectors $(\hat{\varepsilon}_+^{\boldsymbol{k}}, \hat{\varepsilon}_-^{\boldsymbol{k}}, \hat{\boldsymbol{k}})$, where the vectors $\hat{\varepsilon}_\pm^{\boldsymbol{k}}$ are defined as

$$\hat{\varepsilon}_\pm^{\boldsymbol{k}} = \frac{1}{\sqrt{2}}\, (\hat{\varepsilon}_1^{\boldsymbol{k}} \pm i\, \hat{\varepsilon}_2^{\boldsymbol{k}}).$$

One can show that these vectors satisfy the following properties:

$$\hat{\varepsilon}_+^{\boldsymbol{k}} \cdot \hat{\varepsilon}_+^{\boldsymbol{k}*} = 1,\ \hat{\varepsilon}_-^{\boldsymbol{k}} \cdot \hat{\varepsilon}_-^{\boldsymbol{k}*} = 1,\ \hat{\varepsilon}_+^{\boldsymbol{k}} \cdot \hat{\varepsilon}_-^{\boldsymbol{k}*} = 0,\ \hat{\varepsilon}_\pm^{\boldsymbol{k}*} = \hat{\varepsilon}_\mp^{\boldsymbol{k}},\ \hat{\varepsilon}_\pm^{-\boldsymbol{k}} = -\hat{\varepsilon}_\mp^{\boldsymbol{k}},\ i\,\hat{\boldsymbol{k}} \times \hat{\varepsilon}_\pm^{\boldsymbol{k}} = \pm\hat{\varepsilon}_\pm^{\boldsymbol{k}}.$$

[15]M. M. Anber and L. Sorbo, JCAP **10**, 018 (2006);

C. Caprini and L. Sorbo, JCAP **10**, 056 (2014);

D. Chowdhury, L. Sriramkumar and M. Kamionkowski, JCAP **10**, 031 (2018).

Quantization in the helical basis

In the helical basis, on quantization in the Coulomb gauge, the vector potential $\hat{A}_i$ can be Fourier decomposed as follows[16]:

$$\hat{A}_i(\eta, \boldsymbol{x}) = \sqrt{4\pi} \int \frac{\mathrm{d}^3\,\boldsymbol{k}}{(2\pi)^{3/2}} \sum_{\sigma=\pm} \left[\varepsilon_{\sigma i}^{\boldsymbol{k}}\, \hat{b}_{\boldsymbol{k}}^{\sigma}\, \bar{A}_k^{\sigma}(\eta)\, \mathrm{e}^{i\,\boldsymbol{k}\cdot\boldsymbol{x}} + \varepsilon_{\sigma i}^{\boldsymbol{k}\,*}\, \hat{b}_{\boldsymbol{k}}^{\sigma\dagger}\, \bar{A}_k^{\sigma*}(\eta)\, \mathrm{e}^{-i\,\boldsymbol{k}\cdot\boldsymbol{x}} \right],$$

where the Fourier modes $\bar{A}_k^{\sigma}$ satisfy the differential equation

$$\bar{A}_k^{\sigma\,\prime\prime} + 2\,\frac{J'}{J}\,\bar{A}_k^{\sigma\,\prime} + \left(k^2 + \frac{\sigma\,\gamma\,k}{J^2}\,\frac{\mathrm{d}I^2}{\mathrm{d}\eta} \right)\,\bar{A}_k^{\sigma} = 0.$$

Note that $\sigma = \pm 1$, and this causes difference in the evolution of the modes depending on I and γ. In terms of $\mathcal{A}_k^{\sigma} = J\,\bar{A}_k^{\sigma}$, the above equation reduces to
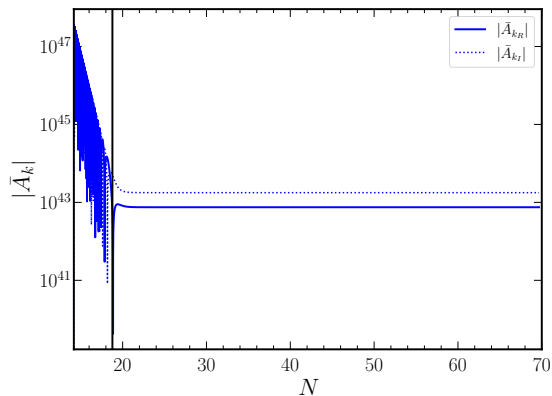
$$\mathcal{A}_k^{\sigma\,\prime\prime} + \left(k^2 + \frac{2\,\sigma\,\gamma\,k\,I\,I'}{J^2} - \frac{J''}{J} \right)\,\mathcal{A}_k^{\sigma} = 0.$$

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[16]See, e.g., C. Caprini and L. Sorbo, JCAP **10**, 056 (2014);
D. Chowdhury, L. Sriramkumar and M. Kamionkowski, JCAP **10**, 031 (2018).

Behavior of the non-helical and helical electromagnetic modes
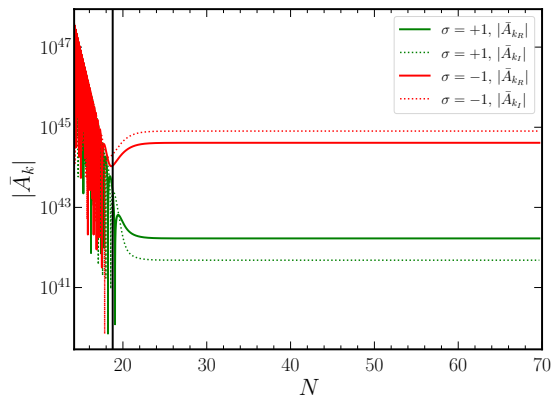
Typical behavior of the real (as solid curves) and imaginary (as dotted curves) parts of the non-helical (in blue, on the left) and helical (in red and green, on the right) electromagnetic modes have been plotted as function of e-folds for a specific wave number. The vertical lines indicates the time when the wave number leaves the Hubble radius[17].

[17] D. Chowdhury, L. Sriramkumar and M. Kamionkowski, JCAP **10**, 031 (2018).

Power spectra of the helical electromagnetic fields

We shall hereafter focus on the case wherein $I = J$.

The power spectra of the magnetic and electric fields can be expressed in terms of the modes $\bar{A}_k^\sigma$ and the coupling function J as follows[18]:

$$\mathcal{P}_{\mathrm{B}}(k) = \frac{k^5}{4\,\pi^2}\,\frac{J^2}{a^4}\,\left[\left|\bar{A}_k^+\right|^2 + \left|\bar{A}_k^-\right|^2\right] = \frac{k^5}{4\,\pi^2\,a^4}\,\left[\left|\mathcal{A}_k^+\right|^2 + \left|\mathcal{A}_k^-\right|^2\right],$$

$$\mathcal{P}_{\mathrm{E}}(k) = \frac{k^3}{4\,\pi^2}\,\frac{J^2}{a^4}\,\left[\left|\bar{A}_k^{+\prime}\right|^2 + \left|\bar{A}_k^{-\prime}\right|^2\right] = \frac{k^3}{4\,\pi^2\,a^4}\,\left[\left|\mathcal{A}_k^{+\prime} - \frac{J'}{J}\mathcal{A}_k^+\right|^2 + \left|\mathcal{A}_k^{-\prime} - \frac{J'}{J}\mathcal{A}_k^-\right|^2\right].$$

[18]See, for example, C. Caprini and L. Sorbo, JCAP **10**, 056 (2014);
   D. Chowdhury, L. Sriramkumar and M. Kamionkowski, JCAP **10**, 031 (2018).

Normal modes and power spectra in de Sitter inflation

For the coupling function J in de Sitter inflation we had considered earlier, the solutions to the electromagnetic modes satisfying the Bunch-Davies initial conditions can be written as follows[19]:

$$\mathcal{A}_k^\sigma(\eta) = \frac{1}{\sqrt{2\,k}}\, \mathrm{e}^{-\pi\,\sigma\,n\,\gamma/2}\, W_{i\,\sigma\,n\,\gamma,\nu}(2\,i\,k\,\eta),$$

where $\nu = n + (1/2)$, and $W_{\lambda,\mu}(z)$ denotes the Whittaker function.

The resulting power spectra can be expressed as

$$\mathcal{P}_{\mathrm{B}}(k) = \frac{H_{\mathrm{I}}^4}{8\,\pi^2}\, \frac{\Gamma^2(|2\,n+1|)}{|\Gamma(\frac{1}{2}+i\,n\,\gamma+|n+\frac{1}{2}|)|^2}\, \frac{\cosh(n\,\pi\,\gamma)}{2^{|2\,n+1|-2}}\, (-k\,\eta_{\mathrm{e}})^{5-|2\,n+1|},$$

$$\mathcal{P}_{\mathrm{E}}(k) = \frac{H_{\mathrm{I}}^4}{4\,\pi^2}\, \frac{\Gamma^2(2\,|n|)}{|\Gamma(|n|+i\,n\,\gamma)|^2}\, \frac{\gamma^2}{1+\gamma^2}\, \frac{\cosh(n\,\pi\,\gamma)}{2^{2\,|n|-2}}\, (-k\,\eta_{\mathrm{e}})^{4-2\,|n|},$$

with the factor $\gamma^2/(1+\gamma^2)$ arising *only* for positive values of n. The spectral indices for the magnetic and electric fields are given by $n_{\mathrm{B}} = 5 - |2\,n+1|$ and $n_{\mathrm{E}} = 4 - 2\,|n|$.

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[19]See, for instance, R. Sharma, K. Subramanian and T. R. Seshadri, Phys. Rev. D **97**, 083503 (2018).

Scale invariant amplitudes and the non-helical limit

When $n = 2$, we find that the spectra of the helical magnetic and electric fields, evaluated in the limit $(-k\,\eta_{\mathrm{e}}) \ll 1$, can be written as[20]

$$\mathcal{P}_{\mathrm{B}}(k) = \frac{9\,H_{\mathrm{I}}^4}{4\,\pi^2}\,f(\gamma),$$

$$\mathcal{P}_{\mathrm{E}}(k) = \frac{9\,H_{\mathrm{I}}^4}{4\,\pi^2}\,f(\gamma)\left[\gamma^2 - \frac{\sinh^2(2\,\pi\,\gamma)}{3\,\pi\,(1+\gamma^2)\,f(\gamma)}\,(-k\,\eta_{\mathrm{e}}) + \frac{1}{9}\,\left(1 + 23\,\gamma^2 + 40\,\gamma^4\right)\,(-k\,\eta_{\mathrm{e}})^2\right],$$

where the function $f(\gamma)$ is given by

$$f(\gamma) = \frac{\sinh(4\,\pi\,\gamma)}{4\,\pi\,\gamma\,(1 + 5\,\gamma^2 + 4\,\gamma^4)}.$$

Note that the non-helical case case corresponds to $\gamma = 0$.

Also, note that the power spectra in the helical case are enhanced when compared to the non-helical case.

[20] S. Tripathy, D. Chowdhury, R. K. Jain and L. Sriramkumar, Phys. Rev. D **105**, 063519 (2022).

Estimating the strengths of the magnetic fields at the present epoch I

In the case of *instantaneous* reheating, the spectrum of the magnetic field today, say, $\mathcal{P}_{\mathrm{B}}^0(k)$, can be related to the spectrum $\mathcal{P}_{\mathrm{B}}(k)$ at the end of inflation as follows:

$$\mathcal{P}_{\mathrm{B}}^0(k) \simeq \mathcal{P}_{\mathrm{B}}(k) \left(\frac{a_{\mathrm{e}}}{a_0}\right)^4,$$

where a_{e} is the scale factor at the end of inflation, while a_0 denotes the scale factor today.

Conservation of entropy post inflation allows us to write

$$\frac{a_0}{a_{\mathrm{e}}} = \left(\frac{g_{\mathrm{s,e}}}{g_{\mathrm{s,0}}}\right)^{1/3} \frac{T_{\mathrm{e}}}{T_0},$$

where $(T_{\mathrm{e}}, g_{\mathrm{s,e}})$ and $(T_0, g_{\mathrm{s,0}})$ denote the temperature and the effective number of relativistic degrees of freedom at the onset of the radiation dominated epoch and today.

Estimating the strengths of the magnetic fields at the present epoch II

If we consider $g_{\mathrm{e}} = 106.75$, since $g_0 = 3.36$ and $T_0 = 2.725\,\mathrm{K}$, we obtain that

$$\frac{a_0}{a_{\mathrm{e}}} \simeq 2.8 \times 10^{28} \left(\frac{H_{\mathrm{I}}}{10^{-5}\,M_{\mathrm{Pl}}} \right)^{1/2} .$$

Given the scale invariant amplitude for the magnetic field at the end of inflation in the $n = 2$, helical case, upon using the above result, we can estimate the present day strength of the magnetic field, say, B_0 (at any scale), to be[21]

$$B_0 \simeq 4.5 \times 10^{-12} \left(\frac{H_{\mathrm{I}}}{10^{-5}\,M_{\mathrm{Pl}}} \right) f^{1/2}(\gamma)\,\mathrm{G}.$$

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[21] S. Tripathy, D. Chowdhury, R. K. Jain and L. Sriramkumar, Phys. Rev. D **105**, 063519 (2022).

Plan of the talk

Coupling to the scalar curvature

Consider the case of the electromagnetic field that is coupled to the scalar curvature R and is described by the action

$$S[A^\mu] = -\frac{1}{16\,\pi}\int \mathrm{d}^4 x\,\sqrt{-g}\,J^2(R)\,F_{\mu\nu}\,F^{\mu\nu}.$$

In slow roll inflation, the scalar curvature behaves in terms of the conformal time as $R \propto \eta^{2\,\epsilon_1}$, where ϵ_1 is the first slow roll parameter.

This suggests that we can possibly work with a non-conformal coupling function of the form

$$J(R) = \left[\frac{R(\eta)}{R(\eta_{\mathrm{e}})}\right]^\alpha \simeq \left(\frac{\eta}{\eta_{\mathrm{e}}}\right)^{2\,\epsilon_1\,\alpha},$$

which reduces to our original form of the coupling function if we choose $\alpha = -n/(2\,\epsilon_1)$.

But, in slow roll inflation ϵ_1 is not a constant!

Behavior of the non-conformal coupling function in slow roll inflation
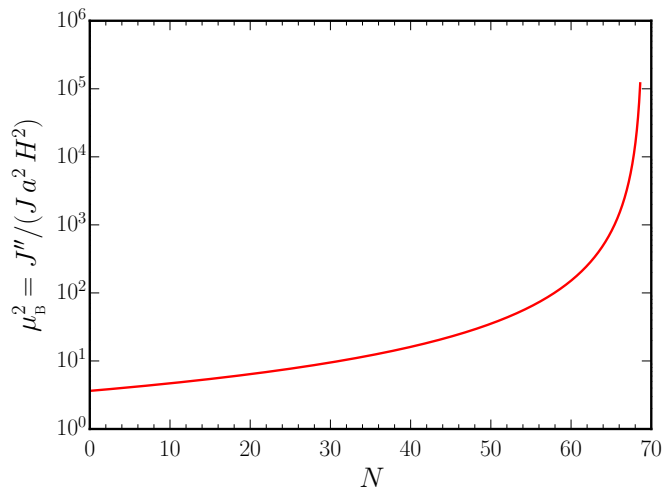
The evolution of the quantity $\mu_{\mathrm{B}}^2 = J''/(J\,a^2\,H^2)$ during slow roll inflation (driven by the popular quadratic potential) in the case wherein J depends on the scalar curvature[22].

[22] S. Tripathy, D. Chowdhury, R. K. Jain and L. Sriramkumar, Phys. Rev. D **105**, 063519 (2022).

Non-conformal coupling function in the quadratic potential

The first model we shall consider is the popular quadratic potential given by

$$V(\phi) = \frac{m^2}{2}\,\phi^2.$$

In such a potential, it is well known that, under the slow roll approximation, the evolution of the field in terms of the e-folds N can be expressed as

$$\phi^2(N) \simeq \phi_{\mathrm{e}}^2 + 4\,(N_{\mathrm{e}} - N)\,M_{\mathrm{Pl}}^2,$$

where $\phi_{\mathrm{e}} \simeq \sqrt{2}\,M_{\mathrm{Pl}}$ denotes the value of the field at the end of inflation at the e-fold N_{e}.

Clearly, we can arrive at the form of $J(N)$ that we desire if we choose $J(\phi)$ to be[23]

$$J(\phi) = \exp\left[-\frac{n}{4\,M_{\mathrm{Pl}}^2}\,(\phi^2 - \phi_{\mathrm{e}}^2)\right].$$

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[23]See, for example, S. Kanno, J. Soda and M. Watanabe, JCAP **12**, 009 (2009);
M. Watanabe, S. Kanno and J. Soda, Phys. Rev. Lett. **102**, 191302 (2009).

Non-conformal coupling function in the Starobinsky model

The other case that we shall discuss is the Starobinsky model described by the potential

$$V(\phi) = V_0 \left[1 - \exp\left(-\sqrt{\frac{2}{3}} \frac{\phi}{M_{\rm Pl}} \right) \right]^2.$$

The evolution of the field in the slow roll approximation is described by the expression

$$N - N_{\rm e} \simeq -\frac{3}{4} \left[\exp\left(\sqrt{\frac{2}{3}} \frac{\phi}{M_{\rm Pl}} \right) - \exp\left(\sqrt{\frac{2}{3}} \frac{\phi_{\rm e}}{M_{\rm Pl}} \right) - \sqrt{\frac{2}{3}} \left(\frac{\phi}{M_{\rm Pl}} - \frac{\phi_{\rm e}}{M_{\rm Pl}} \right) \right],$$

where $\phi_{\rm e}$ is the value of the field at the end of inflation which is determined by the relation $\exp\left[\sqrt{(2/3)}\, \phi_{\rm e}/M_{\rm Pl} \right] \simeq 1 + 2/\sqrt{3}$.

Therefore, to achieve the desired dependence of the coupling function on the scale factor, we can choose $J(\phi)$ in the model to be

$$J(\phi) = \exp\left\{ -\frac{3\,n}{4} \left[\exp\left(\sqrt{\frac{2}{3}} \frac{\phi}{M_{\rm Pl}} \right) - \exp\left(\sqrt{\frac{2}{3}} \frac{\phi_{\rm e}}{M_{\rm Pl}} \right) - \sqrt{\frac{2}{3}} \left(\frac{\phi}{M_{\rm Pl}} - \frac{\phi_{\rm e}}{M_{\rm Pl}} \right) \right] \right\}.$$

Strengths of non-helical electromagnetic fields in slow roll inflation

When $n = 2$, the amplitude of the spectra of the *non-helical* magnetic and electric fields at the pivot scale k_* can be expressed as

$$\frac{\mathcal{P}_{\mathrm{B}}(k)}{M_{\mathrm{Pl}}^4} \simeq \frac{9\,\pi^2}{16}\,(r\,A_{\mathrm{s}})^2\,, \quad \frac{\mathcal{P}_{\mathrm{E}}(k)}{M_{\mathrm{Pl}}^4} \simeq \frac{\mathcal{P}_{\mathrm{B}}(k)}{9\,M_{\mathrm{Pl}}^4}\,\left(\frac{k_*}{k_{\mathrm{e}}}\right)^2 \simeq \frac{\mathcal{P}_{\mathrm{B}}(k)}{9\,M_{\mathrm{Pl}}^4}\,\mathrm{e}^{-100}.$$

In these expressions, $A_{\mathrm{s}} = 2.1 \times 10^{-9}$ denotes the observed amplitude of the scalar power spectrum at the pivot scale and r represents the tensor-to-scalar ratio.

In the two slow roll inflationary models of our interest, viz. the quadratic potential and the Starobinsky model, the tensor-to-scalar ratio can be estimated to be $r \simeq (1.6 \times 10^{-1}, 4.8 \times 10^{-3})$.

The above expressions then suggest that these models will generate non-helical magnetic fields of amplitudes $\mathcal{P}_{\mathrm{B}}(k) \simeq (6.27 \times 10^{-19}, 5.64 \times 10^{-22})\,M_{\mathrm{Pl}}^4$. Moreover, according to expressions above, $\mathcal{P}_{\mathrm{B}}(k) \simeq 10^{-20}\,M_{\mathrm{Pl}}^4$ implies that $\mathcal{P}_{\mathrm{E}}(k) \simeq 10^{-66}\,M_{\mathrm{Pl}}^4$.

Strengths of helical electromagnetic fields in slow roll inflation

In the helical case, when $n = 2$, the amplitude of the spectra of the magnetic and electric fields can be expressed as

$$\frac{\mathcal{P}_{\mathrm{B}}(k)}{M_{\mathrm{Pl}}^4} \simeq \frac{9\,\pi^2}{16}\,(r\,A_{\mathrm{s}})^2\,f(\gamma), \quad \frac{\mathcal{P}_{\mathrm{E}}(k)}{M_{\mathrm{Pl}}^4} \simeq \frac{\mathcal{P}_{\mathrm{B}}(k)}{M_{\mathrm{Pl}}^4}\,\gamma^2.$$

If we need to avoid backreaction due to the helical electromagnetic fields which have been generated, we require that

$$\mathcal{P}_{\mathrm{B}}(k) + \mathcal{P}_{\mathrm{E}}(k) \ll \rho_{\mathrm{I}} = 3\,H_{\mathrm{I}}^2\,M_{\mathrm{Pl}}^2.$$

Since we are considering inflationary models wherein $H_{\mathrm{I}}/M_{\mathrm{Pl}} \lesssim 10^{-5}$, we find that the condition for avoiding backreaction leads to $f(\gamma)\,(1 + \gamma^2) \lesssim 10^{10}$.

This limits the value of γ to be $\gamma \lesssim 2.5$.

Spectra of electromagnetic fields in slow roll inflation
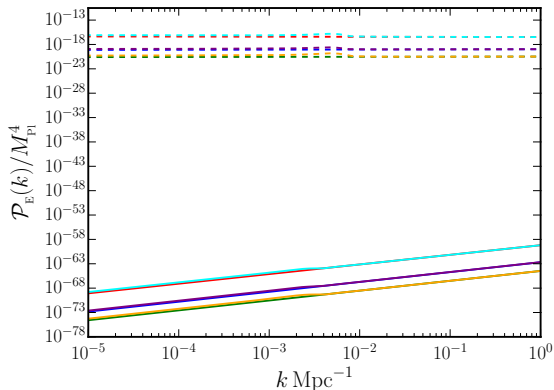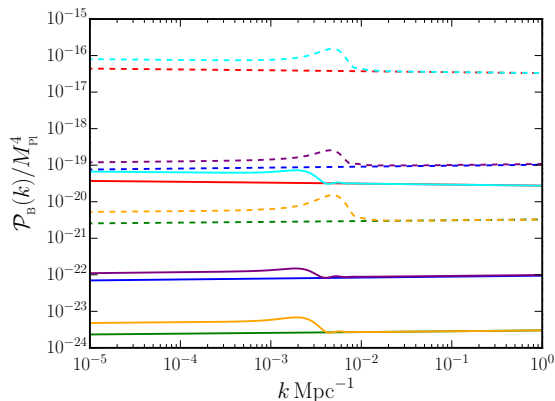
The spectra of the magnetic (on the left) and electric (on the right) fields in the non-helical (solid curves) and helical (dashed curves) cases, arising in three slow roll inflationary models, have been plotted (in red, blue and green) over the CMB scales. We have also plotted the corresponding spectra when a step has been introduced in these potentials (in cyan, purple and orange). We have set $\gamma = 1$ for which $f(\gamma) \simeq 10^3$. ▶ Back to scalar power spectra with features

▶ Back to coupling function in the second Starobinsky model

Plan of the talk

Potentials with a step

Given an inflationary model described by the potential $V(\phi)$, we can introduce a step in the potential as follows[24]:

$$V_{\text{step}}(\phi) = V(\phi) \left[1 + \alpha \tanh \left(\frac{\phi - \phi_0}{\Delta\phi}\right)\right],$$

where, evidently, ϕ_0, α and $\Delta\phi$ denote the location, the height and the width of the step.

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[24] D. K. Hazra, M. Aich, R. K. Jain, L. Sriramkumar and T. Souradeep, JCAP **10**, 008 (2010).

Suppressing power on large scales

In this context, the first example that we shall consider is another model due to Starobinsky, which is governed by the potential[25]

$$V(\phi) = \begin{cases} V_0 + A_+ \, (\phi - \phi_0), & \text{for } \phi > \phi_0, \\ V_0 + A_- \, (\phi - \phi_0), & \text{for } \phi < \phi_0. \end{cases}$$

The second model that we shall consider is the so-called punctuated inflationary model described by the potential[26]

$$V(\phi) = \frac{m^2}{2} \, \phi^2 - \frac{2 \, m^2}{3 \, \phi_0} \, \phi^3 + \frac{m^2}{4 \, \phi_0^2} \, \phi^4.$$

It is easy to show that this potential contains a point of inflection at ϕ_0.

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[25] A. A. Starobinsky, JETP Lett. **55**, 489 (1992).

[26] R. K. Jain, P. Chingangbam, J-O. Gong, L. Sriramkumar and T. Souradeep, JCAP **01**, 009 (2009);

R. K. Jain, P. Chingangbam, L. Sriramkumar and T. Souradeep, Phys. Rev. D **82**, 023509 (2010);

H. V. Ragavendra, D. Chowdhury and L. Sriramkumar, Phys. Rev. D **106**, 043535 (2022).

Scalar power spectra with features over the CMB scales
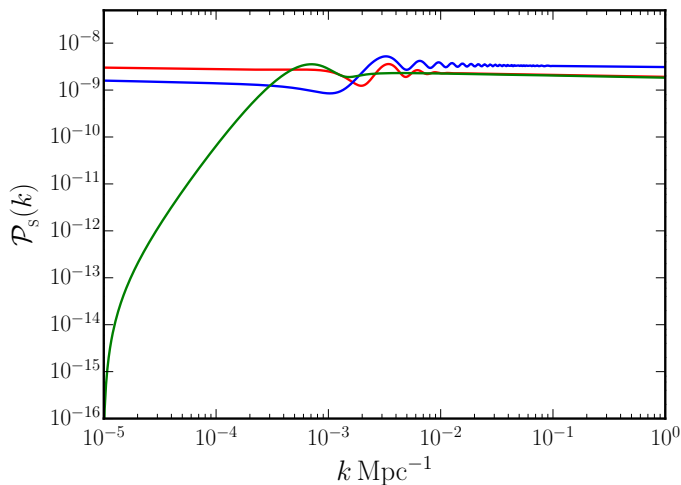
The scalar power spectra with features over the CMB scales in the cases of the quadratic potential with a step (in red), the second Starobinsky model (in blue), and the model of punctuated inflation (in green).

▸ Spectra of electromagnetic fields in slow roll inflation

Potentials admitting ultra slow roll inflation
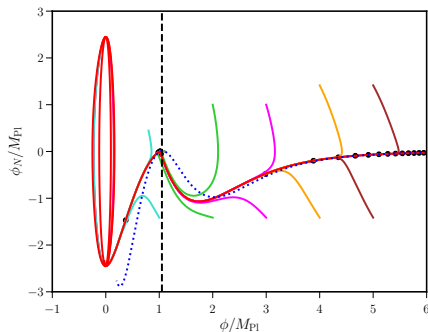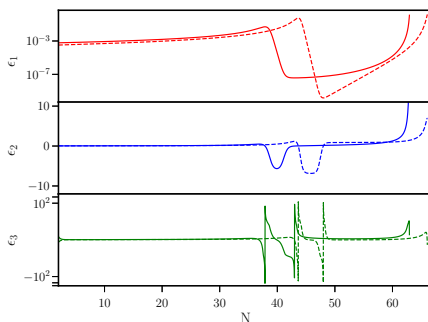
Potentials leading to ultra slow roll inflation (with $x = \phi/v$, v being a constant)[27]:

$$\text{USR1} : V(\phi) = V_0 \frac{6\,x^2 - 4\,\alpha\,x^3 + 3\,x^4}{(1 + \beta\,x^2)^2},$$

$$\text{USR2} : V(\phi) = V_0 \left\{ \tanh\left(\frac{\phi}{\sqrt{6}\,M_{\text{Pl}}} \right) + A \sin\left[\frac{\tanh\left[\phi/\left(\sqrt{6}\,M_{\text{Pl}}\right) \right]}{f_\phi} \right] \right\}^2.$$

[27] J. Garcia-Bellido and E. R. Morales, Phys. Dark Univ. **18**, 47 (2017);
I. Dalianis, A. Kehagias and G. Tringas, JCAP **01**, 037 (2019).

Potentials permitting punctuated inflation
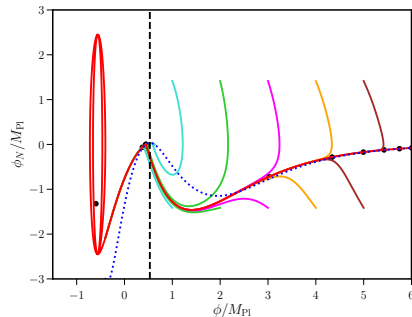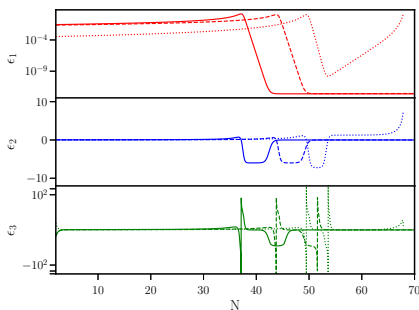
Potentials admitting punctuated inflation[28]:

$$\mathrm{PI1}: V(\phi) = V_0 \left(1 + B\,\phi^4\right), \quad \mathrm{PI2}: V(\phi) = \frac{m^2}{2}\,\phi^2 - \frac{2\,m^2}{3\,\phi_0}\,\phi^3 + \frac{m^2}{4\,\phi_0^2}\,\phi^4,$$

$$\mathrm{PI3}: V(\phi) = V_0 \left[c_0 + c_1\,\tanh\left(\frac{\phi}{\sqrt{6\,\alpha}\,M_{\mathrm{Pl}}}\right) + c_2\,\tanh^2\left(\frac{\phi}{\sqrt{6\,\alpha}\,M_{\mathrm{Pl}}}\right) + c_3\,\tanh^3\left(\frac{\phi}{\sqrt{6\,\alpha}\,M_{\mathrm{Pl}}}\right)\right]^2.$$

[28] D. Roberts, A. R. Liddle and D. H. Lyth, Phys. Rev. D **51**, 4122 (1995);

R. K. Jain, P. Chingangbam, J.-O. Gong, L. Sriramkumar and T. Souradeep, JCAP **01**, 009 (2009);

I. Dalianis, A. Kehagias and G. Tringas, JCAP **01**, 037 (2019).

Scalar power spectra with enhanced power on small scales
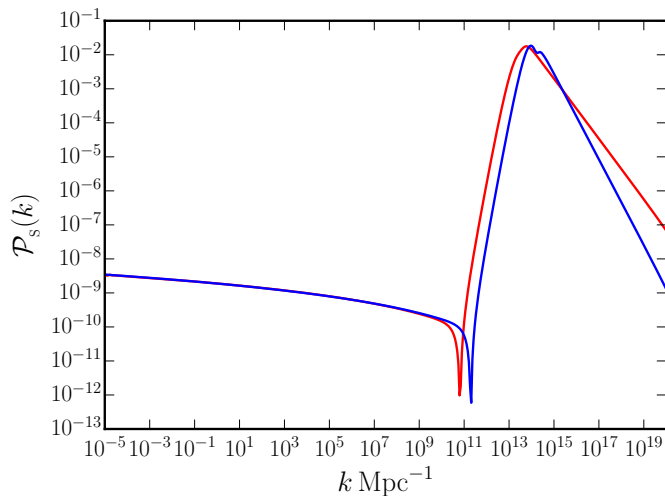
The scalar power spectra with enhanced power on small scales that arise in the inflationary models USR2 (in red) and PI3 (in blue)[29].

[29]H. V. Ragavendra, P. Saha, L. Sriramkumar and J. Silk, Phys. Rev. D **103**, 083510 (2021).

Non-conformal coupling function(s) in the second Starobinsky model

In the second Starobinsky model, on working in the slow roll approximation, the evolution of the scalar field on either side of the transition can be expressed as

$$\phi_+(N) \simeq -\left(\frac{V_0}{A_+} - \phi_0\right) + \left[\left(\phi_i - \phi_0 + \frac{V_0}{A_+}\right)^2 - 2\,M_{\mathrm{Pl}}^2\,N\right]^{1/2},$$

$$\phi_-(N) \simeq -\left(\frac{V_0}{A_-} - \phi_0\right) + \left[\left(\frac{V_0}{A_-}\right)^2 - 2\,M_{\mathrm{Pl}}^2\,(N - N_0)\right]^{1/2}.$$

These solutions allow us to construct the following coupling functions:

$$J_+(\phi) = J_{0+} \exp\left\{-\frac{n}{2\,M_{\mathrm{Pl}}^2}\left[\left(\phi_+ - \phi_0 + \frac{V_0}{A_+}\right)^2 - \left(\phi_i - \phi_0 + \frac{V_0}{A_+}\right)^2\right]\right\},$$

$$J_-(\phi) = J_{0-} \exp\left\{-\frac{n}{2\,M_{\mathrm{Pl}}^2}\left[\left(\phi_- - \phi_0 + \frac{V_0}{A_-}\right)^2 - \left(\frac{V_0}{A_-}\right)^2 - 2\,N_0\,M_{\mathrm{Pl}}^2\right]\right\},$$

where the constants $J_{0\pm}$ are to be chosen suitably so that $J_\pm(\phi_e) = 1$, i.e. the value of J is unity at the end of inflation.

Spectra of the magnetic field in the second Starobinsky model
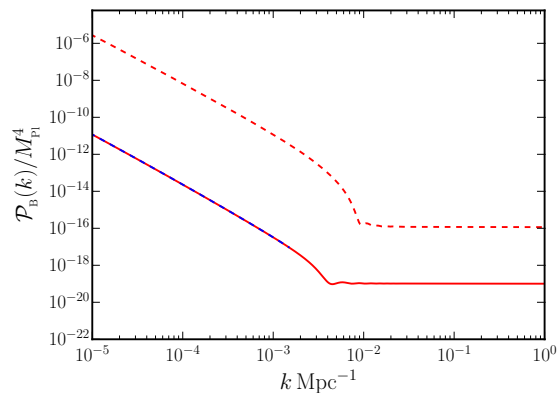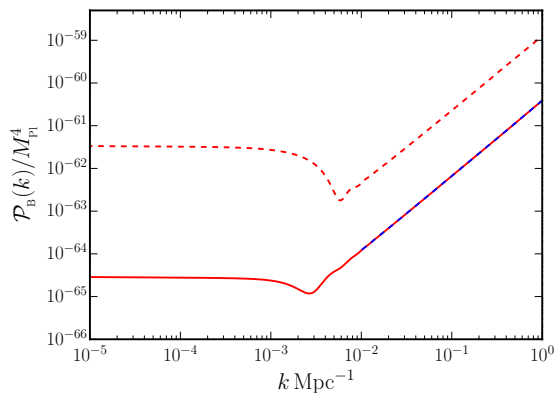
The power spectra of the magnetic field arising in the second Starobinsky model for the two choices of coupling functions $J_+(\phi)$ (on the left) and $J_-(\phi)$ (on the right) have been plotted for $n = 2$ in the non-helical (in solid red) as well as the helical (in dashed red) cases.

Constructing the coupling function in punctuated inflation

Since the solution for the evolution of the scalar field in the first model of punctuated inflation is not easy to obtain analytically, we need to turn to numerics.

We choose a quadratic function of the form

$$N(\phi) = a_1 \left(\frac{\phi^2}{M_{\mathrm{Pl}}^2} \right) + b_1 \left(\frac{\phi}{M_{\mathrm{Pl}}} \right) + c_1$$

to fit the numerical solution we have obtained in the initial slow roll regime.

To evaluate the spectra of the electromagnetic fields, we shall work with a coupling function of the form

$$J(\phi) = \exp \left\{ n \left[a_1 \left(\frac{\phi^2 - \phi_{\mathrm{e}}^2}{M_{\mathrm{Pl}}^2} \right) + b_1 \left(\frac{\phi - \phi_{\mathrm{e}}}{M_{\mathrm{Pl}}} \right) \right] \right\}.$$

Spectra of electromagnetic fields in punctuation inflation
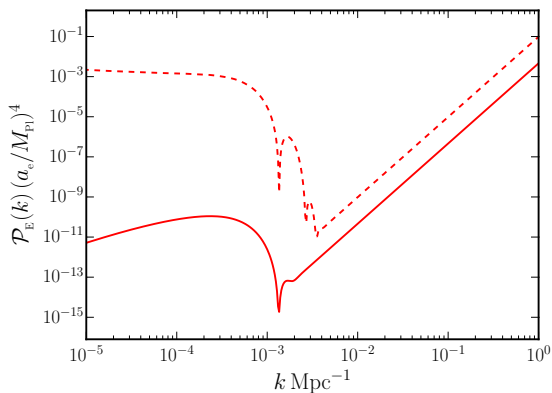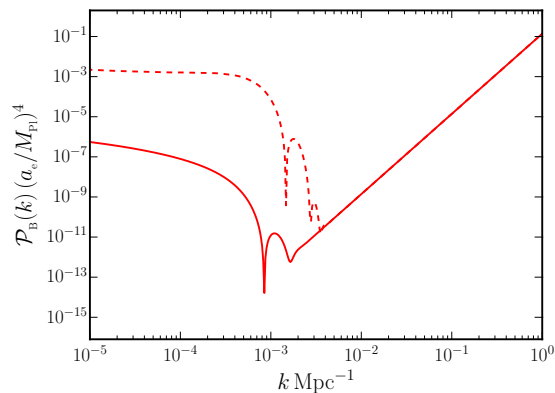
The spectra of the magnetic (on the left) and electric (on the right) fields arising in the case of the first punctuated inflation model have been plotted for both the non-helical (as solid curves) and helical (as dashed curves) cases. ▸ Models leading to suppression of power on large scales

Behavior of the coupling function in models permitting ultra slow roll
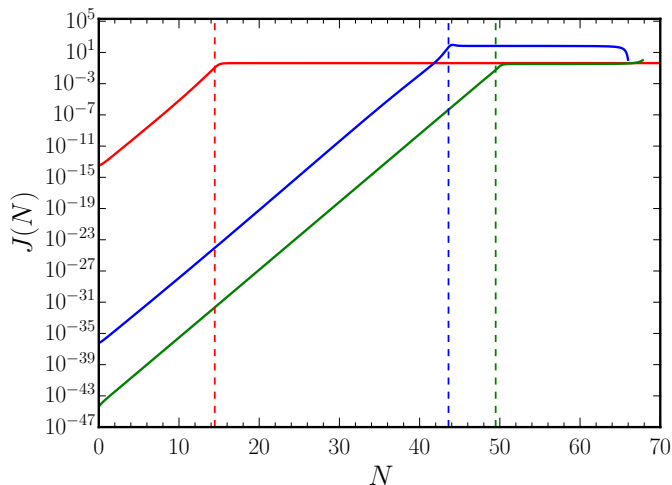
The evolution of the non-conformal coupling function J in inflationary models leading to ultra slow roll. Note that the coupling function does not change appreciably once ultra slow roll sets in (indicated by the vertical lines). ▸ Back to spectra of electromagnetic fields in two field models

Resulting spectra of the electromagnetic fields
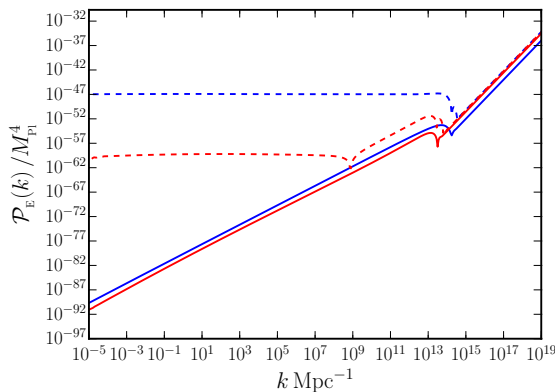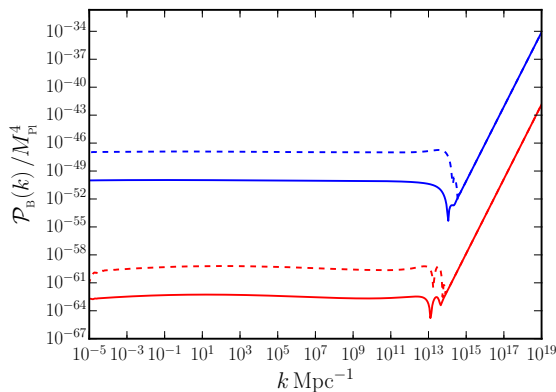
The spectra of the magnetic (on the left) and electric (on the right) fields arising in two inflationary models permitting a period of ultra slow roll (at late times) have been plotted in the non-helical (as solid lines) and helical (as dashed lines) cases[30].

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[30] S. Tripathy, D. Chowdhury, R. K. Jain and L. Sriramkumar, Phys. Rev. D **105**, 063519 (2022).

Can the features in the electromagnetic spectra be ironed out?
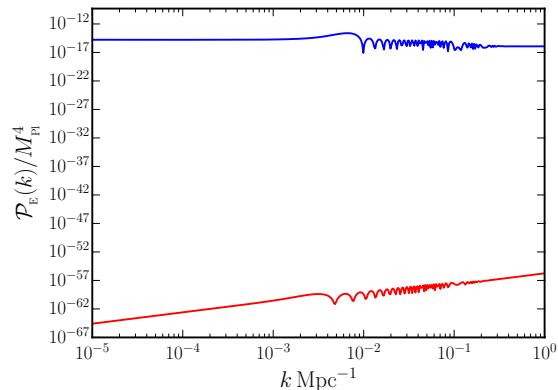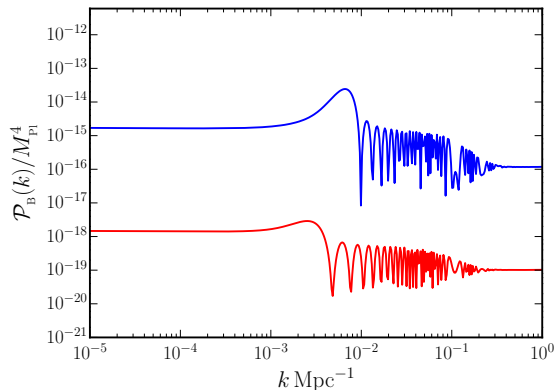
Spectra of electromagnetic fields obtained for the following coupling function in the second Starobinsky model[31]:

▸ Coupling function in the second Starobinsky model

$$J(\phi) = \frac{J_1}{2\,J_{0+}}\left[1 + \tanh\left(\frac{\phi - \phi_0}{\Delta\phi_1}\right)\right]J_+(\phi) + \frac{J_1}{2\,J_{0-}}\left[1 - \tanh\left(\frac{\phi - \phi_0}{\Delta\phi_1}\right)\right]J_-(\phi).$$

[31] S. Tripathy, D. Chowdhury, R. K. Jain and L. Sriramkumar, Phys. Rev. D **105**, 063519 (2022).

Plan of the talk

The two field models of interest

It has been shown that two scalar fields ϕ and χ governed by the action

$$S[\phi, \chi] = \int \mathrm{d}^4x \sqrt{-g} \left[-\frac{1}{2} \partial^\mu \phi \, \partial_\mu \phi - \frac{f(\phi)}{2} \partial^\mu \chi \, \partial_\mu \chi - V(\phi, \chi) \right]$$

described by *separable* potentials and the non-canonical functions $f(\phi) = \exp(2\,b\,\phi)$ or $f(\phi) = \exp(2\,b\,\phi^2)$ can lead to features in the scalar power spectrum.

While the potential[32]

$$V(\phi, \chi) = \frac{1}{2} m_\phi^2 \phi^2 + V_0 \frac{\chi^2}{\chi_0^2 + \chi^2}$$

leads to a suppression in power on large scales, the potential[33]

$$V(\phi, \chi) = V_0 \frac{\phi^2}{\phi_0^2 + \phi^2} + \frac{1}{2} m_\chi^2 \chi^2$$

generates enhanced power on small scales.

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[32]M. Braglia, D. K. Hazra, L. Sriramkumar and F. Finelli, JCAP **08**, 025 (2020).

[33]M. Braglia, D. K. Hazra, F. Finelli, G. F. Smoot, L. Sriramkumar and A. A. Starobinsky, JCAP **08**, 001 (2020).

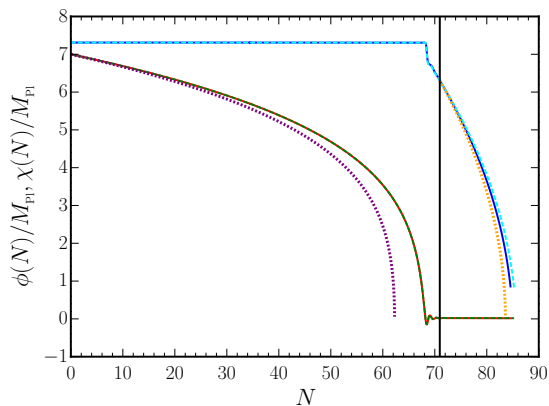# Behavior of the scalar fields and the first slow roll parameter I
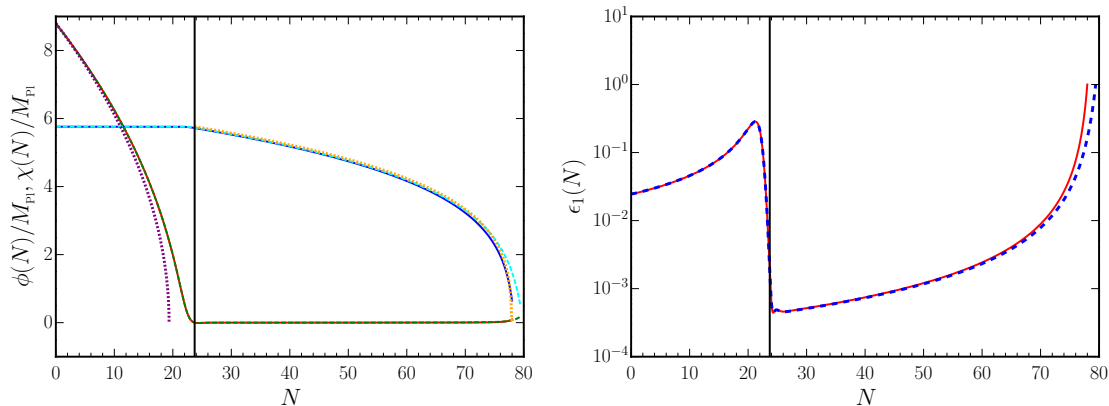
Behavior of the two scalar fields ϕ and χ (in blue and red, on the left) and the first slow roll parameter ϵ_1 (on the right) in the first two field model[34]. Note that there arises a turn in the field space around the e-fold $N = 24$, when the first slow roll parameter begins to decrease before increasing again, leading to the termination of inflation.

[34]M. Braglia, D. K. Hazra, L. Sriramkumar and F. Finelli, JCAP **08**, 025 (2020).

Behavior of the scalar fields and the first slow roll parameter II
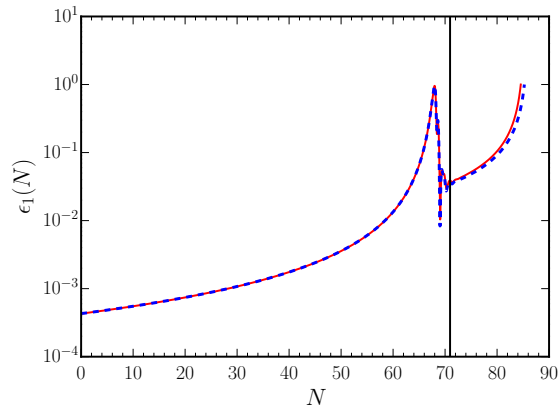
Behavior of the two scalar fields and the first slow roll parameter in the second two field model has been plotted as in the case of the first model[35]. Evidently, barring the time of the turn in field space, the fields broadly behave in the same manner as in the previous example.

[35] M. Braglia, D. K. Hazra, F. Finelli, G. F. Smoot, L. Sriramkumar and A. A. Starobinsky, JCAP **08**, 001 (2020).

Scalar and tensor power spectra in two field models
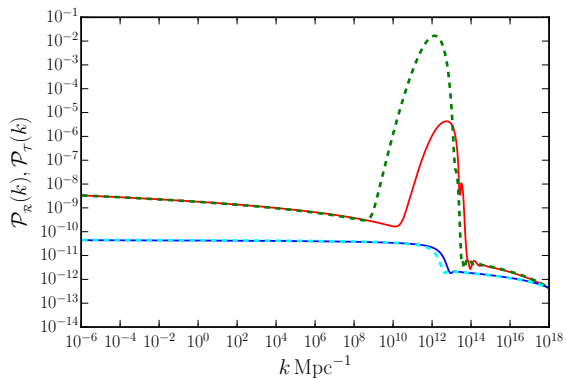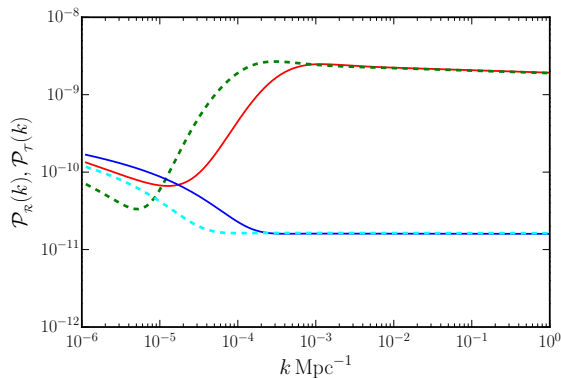
The spectra of curvature (in solid red and dashed green) and tensor (in solid blue and dashed cyan) perturbations, viz. $\mathcal{P}_{\mathcal{R}}(k)$ and $\mathcal{P}_{\mathcal{T}}(k)$, have been plotted for the two field inflationary models that we have considered.[36]

[36] S. Tripathy, D. Chowdhury, H. V. Ragavendra, R. K. Jain and L. Sriramkumar, Phys. Rev. D **107**, 043501 (2023).

Circumventing the challenge in two field models
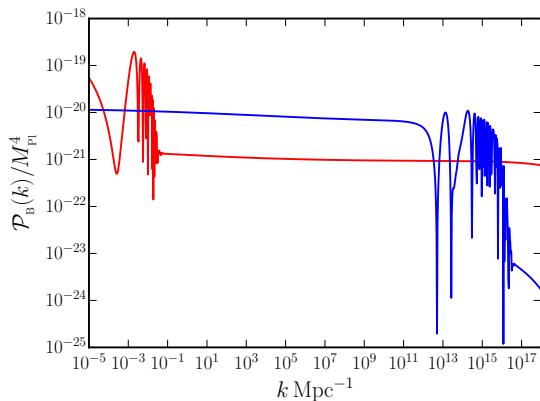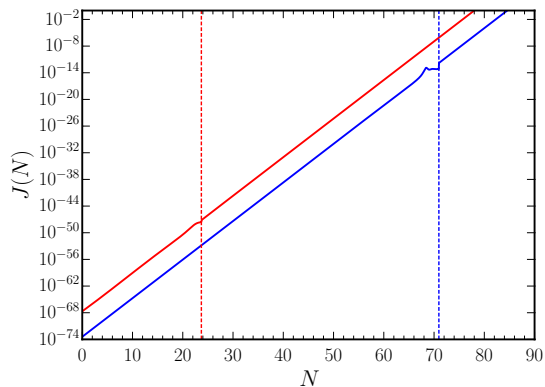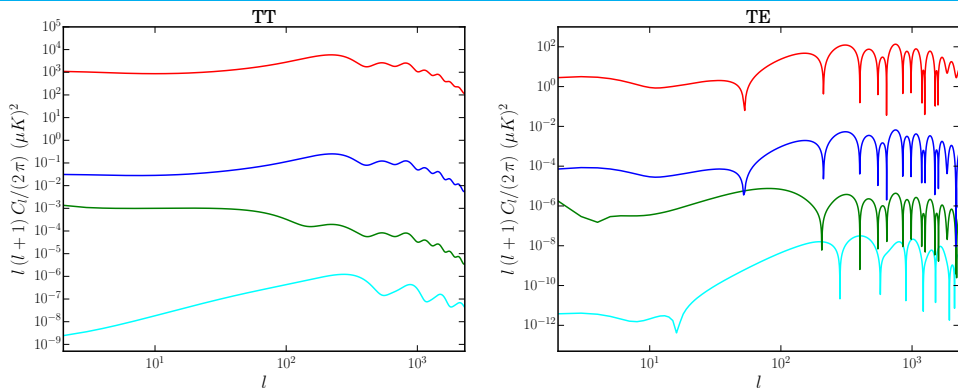

The evolution of the non-conformal coupling function J (on the left) and the corresponding spectrum of magnetic field (on the right) arising in the two field inflationary models leading to features on large (in red) and on small (in blue) scales[37]. The vertical lines (on the left) indicate the time when the turn in the field space takes place. ▸ Evolution of J in ultra slow roll inflation

[37]S. Tripathy, D. Chowdhury, H. V. Ragavendra, R. K. Jain and L. Sriramkumar, Phys. Rev. D **107**, 043501 (2023).

Plan of the talk

# Effects of magnetic fields on the CMB anisotropies I

Contributions of the magnetic modes to the temperature and polarization angular power spectra of the CMB due to the total (i.e. scalar plus tensor) passive (in green) and the total compensated (in cyan) modes in the second two field model. We have arrived at these quantities using MagCAMB corresponding to a magnetic field with smoothed strength of $B_{1\,\mathrm{Mpc}}^0 = 2.05 \times 10^{-1}\,\mathrm{nG}$ *today* and a spectral index of $n_{\mathrm{B}} = -0.0112$[38].

[38] A. Zucca, Y. Li and L. Pogosian, Phys. Rev. D **95**, 063506 (2017).

Effects of magnetic fields on the CMB anisotropies II
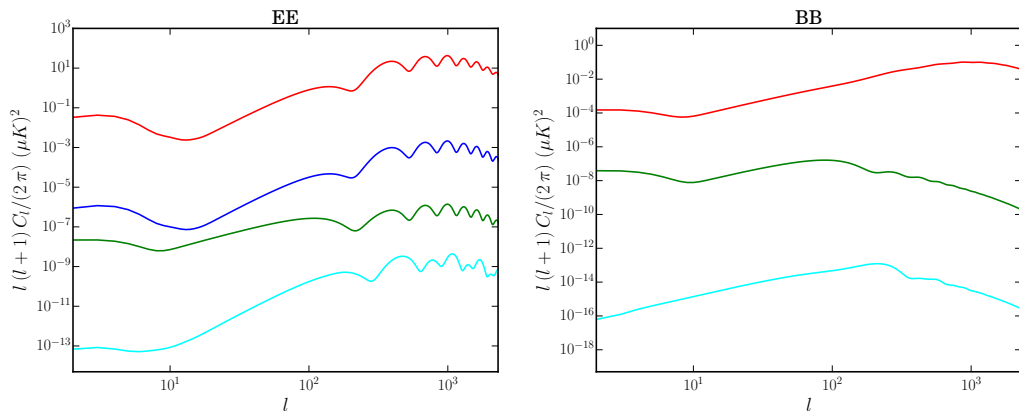
In addition, using CAMB, we have plotted the standard angular power spectra of the CMB (in red) induced by the primary scalar and tensor perturbations. Moreover, we have also presented the contribution due to the curvature perturbation induced by the magnetic field during inflation (in blue), which we have computed using CAMB[39].

[39] S. Tripathy, D. Chowdhury, H. V. Ragavendra, R. K. Jain and L. Sriramkumar, Phys. Rev. D **107**, 043501 (2023).

Plan of the talk

Summary

✦ In single field models of inflation, substantial departures from slow roll inflation can lead to strong features in the spectra of magnetic fields.

✦ Strong departures from slow roll inflation can also suppress the strength of the magnetic field on large scales.

✦ Some of these challenges can be overcome in two field models of inflation. However, there always seems to be an element of fine tuning involved.

✦ It seems necessary to examine the behavior of additional quantities such as the three-point functions and the corresponding observables to arrive at constraints on the nature and form of the non-conformal coupling function.

Collaborators

Sagarika Tripathy   Debika Chowdhury   H. V. Ragavendra   Rajeev Kumar Jain

Thank you for your attention