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CMB anisotropies and inflation from non-standard spinors

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ABSTRACT

The apparent alignment of the cosmic microwave background multipoles on large scales challenges the standard cosmological model. Scalar field inflation is isotropic and cannot account for the observed alignment. We explore the imprints, a non-standard spinor driven inflation would leave on the cosmic microwave background anisotropies. We show it is natural to expect an anisotropic inflationary expansion of the Universe which has the effect of suppressing the low multipole amplitude of the primordial power spectrum, while at the same time to provide the usual inflationary features.

Inflation is a successful theory to explain many cosmological puzzles. However, one does not really know what has driven it, since it most probably occurred near the scale of grand unification, hence far beyond the standard model of particle physics.

In this Letter we study the possibility of non-standard spinors to drive inflation and investigate the possible imprints of such spinors on the Cosmic Microwave Background (CMB) anisotropies. In fact, we consider the possibility that such an effect has already been detected, in the form of the Axis of Evil: an apparent alignment of the CMB multipoles on very large scales [1–3]. While a scalar field driven inflationary epoch is naturally isotropic, an anisotropic expansion might occur within a more complex model. This may lead to the existence of a preferred direction in the primordial power spectrum.

Although the statistical significance of such preferred direction is hard to quantify, a variety of models have been put forward to explain this phenomenon [4–12]. These are motivated since the large scale anisotropy claimed by [13] in the CMB quadrupole and octupole seems to be present at several cosmological scales and observations. In particular the quadrupole and octupole seem also to align with the dipole [14]. Recently, there are claims that such alignment even extends to higher multipoles [15]. Furthermore, the polarization of radio galaxies and the optical polarizations of quasars also indicate a preferred direction pointing at the same direction [16]. Finally, there are several indications from the SDSS data that deviations from isotropy and homogeneity are also present at cluster and galactic scales [17]. Hence, there is an entire set of observations that disfavor isotropy at high confidence level.

As for the non-standard Wigner class spinors, we consider a spin one half matter field with mass dimension one, named elko spinors [18]. These spinors are based on the eigenspinors of the charge conjugation operator. The resulting field theory has the unusual property $(CPT)^2 = -\mathbb{I}$.[1] This particular model belongs to a wider class of so-called flagpole spinors [19]. The spinors have mass dimension one and therefore the only power counting renormalizable interactions of this field with standard matter take place through the Higgs doublet or with gravity [18]. Consider the left-handed part ϕ_L of Dirac spinor ψ in Weyl representation, then an elko spinor is defined by [18]

$$\lambda = \begin{pmatrix} \pm\sigma_2\phi_L^* \\ \phi_L \end{pmatrix}, \tag{1}$$

where ϕ_L^* denotes the complex conjugate of ϕ_L. Since the helicities of ϕ_L and $\sigma^2\phi_L^*$ are opposite [18], one has to distinguish the two possible helicity configurations, therefore

$$\lambda_{\{-,+\}} = \begin{pmatrix} \pm\sigma_2\phi_L^{+*} \\ \phi_L^+ \end{pmatrix}, \qquad \lambda_{\{+,-\}} = \begin{pmatrix} \pm\sigma_2\phi_L^{-*} \\ \phi_L^- \end{pmatrix}. \tag{2}$$

The first entry of the helicity subscript $\{-,+\}$ refers to the upper two-spinor while the second to the lower. Let us henceforth denote the helicity subscript by the indices $u, v, \ldots$ and define the elko dual by

$$\bar{\lambda}_u = i\varepsilon_u^{\ v}\lambda_v^\dagger\gamma^0, \tag{3}$$

with the anti-symmetric symbol $\varepsilon_{\{+,-\}}^{\{-,+\}} = -1 = -\varepsilon_{\{-,+\}}^{\{+,-\}}$. Note that due to the double helicity structure of the spinors, these have an

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[1] The original elko field theory is non-local. However, since CPT is an anti-unitary operator, a local field theory in principle must exist. We thank Raymond Streater for elucidating this point.

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imaginary bi-orthogonal norm [18] with respect to the standard Dirac adjoint $\bar{\psi} = \psi^\dagger \gamma^0$. With the dual defined above one finds (by construction)

$$\vec{\lambda}_u(\mathbf{p})\lambda_v(\mathbf{p}) = \pm 2m\delta_{uv}, \tag{4}$$

where $\mathbf{p}$ denotes the momentum.

Notice that the cosmology of such spinors will be different from the ones investigated by Saha and collaborators [20]. Firstly, the scalar field like equations of motion are second order equations opposed to the first order equations for standard spinors. Moreover, only intrinsically massless Dirac spinors are power counting renormalizable, which is one of the main motivations to analyze non-standard spinors. The natural potential is power counting renormalizable and their structure is much richer than that given by standard spinors.

The introduction of elko spinor fields into an arbitrary curved spacetime can be found in [21]. This resulting theory is based on the following matter action

$$S = \frac{1}{2}\int \left(g^{\mu\nu}\nabla_{(\mu}\vec{\lambda}\nabla_{\nu)}\lambda - m^2\vec{\lambda}\lambda + \alpha[\vec{\lambda}\lambda]^2\right)\sqrt{-g}\,d^4x, \tag{5}$$

where m is the mass of the field and α is a coupling constant.

With the aim to understand possible effects of non-standard spinors in cosmology we investigate a quite general metric given by

$$ds^2 = dt^2 - a(t)^2\left(dx^2 + dy^2\right) - b(t)^2\,dz^2, \tag{6}$$

where $a(t)$ and $b(t)$ are two expansion parameters, that define two Hubble parameters by $H_a = \dot{a}/a$ and $H_b = \dot{b}/b$. Note that this reduces to the isotropic FRW metric in the case where $a = b$. The presence of the spin-connection in the matter part leads to additional couplings between the field and the geometry. Hence, such a spinor driven inflationary epoch can naturally result in anisotropic expansion.

We assume that the non-standard spinors only depend on the time coordinate t. Following [21], the cosmological spinors are given by

$$\lambda_{\{-,+\}} = F(t)\xi, \qquad \lambda_{\{+,-\}} = F(t)\zeta, \tag{7}$$

with their respective dual spinors $\vec{\xi}$ and $\vec{\zeta}$, where ξ and ζ are constant non-standard spinors [21] satisfying

$$\vec{\lambda}_{\{-,+\}}\lambda_{\{-,+\}} = \vec{\lambda}_{\{+,-\}}\lambda_{\{+,-\}} = \pm 2F^2. \tag{8}$$

Henceforth we consider the self-dual spinors with $\vec{\xi}\xi = \vec{\zeta}\zeta = +2$. Note that such constant spinors are the ones compatible with homogeneity. Moreover, they also satisfy the condition that an anisotropic expansion requires the initial alignment of spins over the Hubble horizon.

We are interested in finding a solution of the form

$$a(t) = e^{H_a t}, \qquad b(t) = e^{H_b t}, \qquad F(t) = F_0 = \text{const}. \tag{9}$$

Plugging this ansatz into the Einstein field equations, the equations of motion reduce to a system of algebraic equations

$$\begin{aligned}
&H_a^2 + 2H_aH_b = \frac{8\pi}{m_{\text{pl}}^2}\left(-\frac{2}{4}H_a^2 - \frac{1}{4}H_b^2 + m^2 + \alpha F_0^2\right)F_0^2,\\
&-\left(H_aH_b + H_a^2 + H_b^2\right) = \frac{8\pi}{m_{\text{pl}}^2}\left(\frac{1}{4}H_b^2 - m^2 - \alpha F_0^2\right)F_0^2,\\
&-3H_a^2 = \frac{8\pi}{m_{\text{pl}}^2}\left(\frac{2}{4}H_a^2 - \frac{1}{4}H_b^2 - m^2 - \alpha F_0^2\right)F_0^2.
\end{aligned} \tag{10}$$

These three equations can be simultaneously satisfied with $H_b > H_a > 0$, $F_0 > 0$ and $m > 0$, $\alpha > 0$. These rather complicated expressions can be greatly simplified after the following considerations are taken into account.

It turns out to be convenient to refer to a fictitious isotropic metric (with expansion parameter $\bar{a}(t)$), defined via an averaged Hubble parameter

$$\bar{H} = \frac{2H_a + H_b}{3}, \tag{11}$$

which can be used to parameterize deviations from isotropy by

$$\epsilon_H = \frac{2}{3}\frac{H_b - H_a}{\bar{H}}. \tag{12}$$

The parameter ϵ_H turns out to be expressed solely in terms of the spinorial part F_0. The function $\epsilon_H(F_0)$ is increasing and vanishes at the origin. Since we are interested in a geometry with only small deviations from isotropy, we assume $F_0^2 \ll 1$. This guarantees the usual post-inflation isotropic expansion, and that non-standard spinors never dominate a cosmological epoch.

Expanding the anisotropy parameter and the mean Hubble parameter for small F_0 yields

$$\bar{H} = mF_0\sqrt{\frac{8\pi}{3}}\left(1 + \frac{\alpha}{m^2}F_0^2\right) + O\left(F_0^5\right), \tag{13}$$

$$\epsilon_H = \frac{8\pi}{3}F_0^2 - 2\left(\frac{8\pi}{3}F_0^2\right)^2 + O\left(F_0^6\right). \tag{14}$$

In order to treat the anisotropy as a perturbation around the background, we furthermore assume that $N_*\epsilon_H \ll 1$, where $N_* = \bar{H}t_*$ is the number of e-folds at the end of inflation which we take to be around 60 as in standard inflation.

Next we verify that the usual inflationary parameters (number of e-folds, near scale invariant spectral index, small non-gaussianities) are also in agreement with the present model. In order to calculate them we express the field equations in terms of the averaged Hubble parameter (11) and the deviation from isotropy (12).

It turns out that the terms linear in ϵ_H vanish identically. Therefore, by neglecting term of the order $O(\epsilon_H^2)$ and higher, we find that the average Hubble parameter and the equation of motion for the spinor field are given by

$$\bar{H}^2 = \frac{8\pi}{3m_{\text{pl}}^2}\left(\frac{1}{2}\partial_t\vec{\lambda}\partial_t\lambda - \frac{3}{8}\bar{H}^2\vec{\lambda}\lambda + V(\vec{\lambda}\lambda)\right), \tag{15}$$

$$\partial_{tt}\lambda + 3\bar{H}\partial_t\lambda - \frac{3}{4}\bar{H}\lambda + V_{\vec{\lambda}}(\vec{\lambda}\lambda) = 0, \tag{16}$$

respectively, where the latter equation is indeed exact. Requiring a power counting renormalizable theory uniquely determines the potential to have the form $V(\vec{\lambda}\lambda) = m^2\vec{\lambda}\lambda + \alpha[\vec{\lambda}\lambda]^2$. Note that in contrast to the scalar field case, the matter part (right-hand side) now also contains the Hubble parameter. One can then (cosmologically) re-interpret the non-standard spinors as a scalar field with a time dependent mass. However, since both Hubble parameters are assumed to be constant throughout inflation this merely leads to a shift of the mass parameter. Therefore, although this model naturally allows for anisotropic inflation it is effectively equivalent to standard single field inflation. This greatly simplifies the interpretation of all equations. Since the expressions for the usual inflationary quantities will be similar in this theory. However, as one will see bellow, there are some cosmological imprints which are very particular to an anisotropic inflationary epoch driven by a non-standard spinor, which are not present in the usual scalar field models.

Eq. (15) can be solved for $\bar{H}$ and yields

$$\bar{H}^2 \simeq \frac{8\pi}{3m_{\text{pl}}^2}\left(\frac{1}{2}\partial_t\vec{\lambda}\partial_t\lambda + V(\vec{\lambda}\lambda)\right), \tag{17}$$

where we neglected terms of the order $\bar{\lambda}\lambda/m_{\rm pl}^2$. Where $m_{\rm pl}$ is the Planck mass.

For the slow-roll conditions: $\dot{\lambda}^2/2 \ll V(\bar{\lambda}\lambda)$ and $|\ddot{\lambda}| \ll 3H|\dot{\lambda}|$ we therefore obtain

$$\bar{H}^2 \simeq \frac{8\pi}{3m_{\rm pl}^2} V(\bar{\lambda}\lambda), \qquad 3\bar{H}\partial_t\lambda \simeq -V_{\bar{\lambda}}(\bar{\lambda}\lambda) + \frac{3}{4}\bar{H}^2\lambda. \tag{18}$$

The last term containing $\bar{H}^2$ can again be replaced by the actual expression for $\bar{H}^2$ and results in a rather complicated expression. However, the factors of $m_{\rm pl}$ as before make all additional contributions small.

The spectral index is given in terms of the slow roll parameters $n = 1 - 6\epsilon + 2\eta$. Both the parameters can be calculated straightforwardly from the above equations. The parameter ϵ is given by

$$\epsilon \simeq \frac{m_{\rm pl}^2}{16\pi}\frac{V_{\bar{\lambda}}V_\lambda}{V^2} - \frac{\bar{\lambda}\lambda}{4V} \simeq \frac{m_{\rm pl}^2}{16\pi}\frac{V_{\bar{\lambda}}V_\lambda}{V^2}, \tag{19}$$

as it is usual in scalar field inflation. On the other hand, the parameter η acquires one non-trivial extra term. This term is obtained by differentiating Eq. (18) with respect to t and dividing the resulting equation by $9H^2\partial_t\lambda$, which yields an additional term of $1/12$ to η, then we find

$$\eta \simeq \frac{m_{\rm pl}^2}{8\pi}\frac{V_{\bar{\lambda}\lambda}}{V} - \frac{1}{12}, \tag{20}$$

plus some lower order terms that can be neglected. Hence, we find that η should be smaller for non-standard spinor inflation. Similarly, for the number of e-folds we get

$$N_* = \log\frac{a_f}{a} = \int_t^{t_f} \bar{H}\,dt \simeq \frac{8\pi}{m_{\rm pl}^2}\int_{\bar{\lambda}_f}^{\bar{\lambda}} \frac{V}{V_{\bar{\lambda}}}d\bar{\lambda}.$$

Similarly to the single field inflation scenario, the non-gaussianity parameters within this scenario are given by

$$f_{NL} = \frac{5}{6}(\eta - 2\epsilon), \qquad \tau_{NL} = (\eta - 2\epsilon)^2 = \frac{36}{25}f_{NL}^2, \tag{21}$$

where we neglect the parameter g_{NL} which contains the third derivatives of the potential because of its smallness. For these non-standard spinors the additional contribution of $1/12$ in η will therefore yield a slightly smaller f_{NL} parameter with respect to the usual slow roll inflationary scalar field models

$$f_{NL} = \frac{5}{6}(\eta - 2\epsilon) - \frac{5}{6}\frac{1}{12}. \tag{22}$$

From WMAP3, $-54 \leqslant f_{NL} \leqslant 114$, and the PLANCK satellite's design aim is, among others, to constrain the parameter $|f_{NL}| \leqslant 5$. Hence, we can conclude that non-standard spinor inflation cannot be ruled out by this new data alone.

In standard inflation the primordial power spectrum $P(k)$ only depends on the magnitude of the vector $\mathbf{k}$ which follows from the rotational invariance. An inflationary epoch driven by non-standard spinors results in anisotropic expansion where rotational invariance is broken by a small unit vector $\mathbf{n}$. The imprint of such an anisotropy on the density perturbation power spectrum has the following most general form

$$P'(\mathbf{k}) = P(k)\big(1 + A(k)(\hat{\mathbf{k}}\cdot\mathbf{n})^2\big), \tag{23}$$

where higher powers in $\hat{\mathbf{k}}\cdot\mathbf{n}$ have been suppressed [25,27]. $\hat{\mathbf{k}}$ denotes the unit vector in the direction of $\mathbf{k}$. In leading order in deviations from anisotropy, the rotationally non-invariant part of the power spectrum is characterized by a single function $A(k)$, which is given by

$$A(k) = \frac{9}{2}\epsilon_H \log\left(\frac{k}{\bar{a}(t_*)\bar{H}}\right). \tag{24}$$

Since we assume around 60 e-folds before the end of inflation, we find that $\log(k/(\bar{a}(t_*)\bar{H}))$ is of the order -60 (the minus sign is present due to $\bar{a}_*$ in the denominator) for a wide range of scales which are cosmologically relevant today. CMB measurements probe $k/\bar{a}(t_*)$ up to 10^3. Hence, one can roughly assume that $A(k)$ is k-independent at the astrophysical scales of interest. However, the additional effect on $A(k)$ is a decrease with k. Therefore, anisotropies will predominantly suppress the low multipoles. Together with the explicit solutions given by Eq. (14) in the lowest order we obtain

$$A(k) \sim A_* \approx -720\pi F_0^2. \tag{25}$$

This is consistent with our above approximations and, as one will see, it also allows sufficiently large values of A_* to account for the quadrupole anomaly.

The effects of a preferred direction, $\mathbf{n}$, in the primordial power spectrum will affect the CMB temperature anisotropies by (see e.g. [22–27])

$$\frac{\Delta T}{T}(\mathbf{n}) = \int \mathrm{d}\mathbf{k}\sum_l \left(\frac{2l+1}{4\pi}\right)P_l(\hat{\mathbf{k}}\cdot\mathbf{n})\delta(\mathbf{k})\Theta_l(k), \tag{26}$$

where P_l is the Legendre polynomial. $\Theta_l(k)$ encompasses the transfer functions of the usual isotropic post-inflationary epochs. Hence, it is a function of the magnitude of the wavevector $\mathbf{k}$ only. The CMB power spectra can then be obtained decomposing it into the usual isotropic part plus a primordial anisotropic piece which is of first order in $A(k)$,

$$\langle a_{lm}a^*_{l'm'}\rangle = \langle a_{lm}a^*_{l'm'}\rangle_{\rm iso} + \varphi(lm;l'm'), \tag{27}$$

where the sought-after perturbation is given by

$$\varphi(lm;l'm') = \Xi_{lm;l'm'} \times \int_0^\infty \mathrm{d}k\,k^2 P(k)A(k)\Theta_l(k)\Theta_{l'}(k), \tag{28}$$

where

$$\begin{aligned}\Xi_{lm;l'm'} = {}& \frac{4\pi}{3}\int \mathrm{d}\Omega_{\mathbf{k}}\, Y_l^m(\hat{\mathbf{k}})\big(Y_{l'}^{m'}(\hat{\mathbf{k}})\big)^* \\ & \times \big(n_+Y_1^1(\hat{\mathbf{k}}) + n_-Y_1^{-1}(\hat{\mathbf{k}}) + n_0Y_1^0(\hat{\mathbf{k}})\big)^2.\end{aligned} \tag{29}$$

The constants $\Xi_{lm;l'm'}$ are purely geometric, and n_+, n_0, n_- are the spherical components of the vector that defines the preferred direction. Those are given in [25].

Taking into account only the astrophysical scales of interest for us today ($A(k)$ becomes roughly k-independent) we have $A(k) = A_*$, then we find

$$\begin{aligned}&\frac{\varphi(lm;lm)}{\langle a_{lm}a^*_{lm}\rangle_{\rm iso}} \\ &\quad = \frac{A_*}{2}\left[\sin^2\theta_* + \big(3\cos^2\theta_* - 1\big)\left(\frac{2l^2+2l-2m^2-1}{(2l-1)(2l+3)}\right)\right].\end{aligned} \tag{30}$$

It is interesting to notice that within this scenario one gets a low quadrupole. The multipole spectrum is described by

$$Q_l = \sqrt{\frac{1}{2\pi}\frac{l(l+1)}{(2l+1)}\sum_{m=-l}^{l}\langle a_{lm}a^*_{lm}\rangle_{\rm iso}\left[1 + \frac{\varphi(lm;lm)}{\langle a_{lm}a^*_{lm}\rangle_{\rm iso}}\right]}. \tag{31}$$

The observed value of this is $Q_2^{\rm obs} \approx 5.72 \times 10^{-3}$, while the standard concordance model predicts $Q_2^{\Lambda\rm CDM} \approx 13 \times 10^{-3}$ [29]. It has been suggested in previous works that this discrepancy could also be explained by an ellipsoidality of the universe [31], by inhomogeneous cosmological magnetic fields [32], or a dark energy component with an anisotropic equation of state [33]. This would require that the anisotropy of the background is suitably oriented with respect to the intrinsic quadrupole and cancels its power to a sufficient amount. For any orientation then, we should have $Q_2 \lesssim 19.7 \times 10^{-3}$ to be consistent with observations taking into account the cosmic variance. Inserting the values predicted by the concordance model $a_{2m} = \sqrt{\pi/3} \cdot 13 \times 10^{-6}$ into Eq. (31) one obtains the following quadrupole moment

$$Q_2^{\text{model}} = 13\sqrt{1 + A_*/3} \times 10^{-3}, \tag{32}$$

which must be compared with the observed values for a_{2m} from the cleaned SILC400 (a), WILC3YR (b) and TCM3YR (c) maps, see [28–30], which lead to the following observed quadrupoles

$$\frac{Q_2^{(a)}}{10^{-3}} = 6.08, \qquad \frac{Q_2^{(b)}}{10^{-3}} = 5.77, \qquad \frac{Q_2^{(c)}}{10^{-3}} = 5.30.$$

In order to have agreement between the value predicted by an anisotropic model and the actually observed value, it now becomes clear that models in which the anisotropy is treated as a small quantity can indeed explain the low quadrupole moment we observe. From Eq. (32) we find that A_* should be around $A_* \approx -2.41$. This in turn fixes the spinorial part of the model, namely F_0 should be of the order of $F_0 \approx 0.033$ which in turn leads to $F_0^2 \approx 1.1 \times 10^{-3}$ which clearly is in agreement with our above assumption $F_0^2 \ll 1$.

Taking into account the rather reasonable assumption that for the isotropic background we can assume the modulus of a_{lm} to be equal for all modes. In that case, we can give an explicit expression of the corrected power spectrum

$$Q_l = \sqrt{l(l+1)C_l/2\pi(1 + A_*/3)}. \tag{33}$$

Hence, within this model multipole moments are suppressed by the factor $(1 + A_*/3)$ where we neglected variations of $A(k)$. Such feature might result in a better agreement between the observational data and the theoretical models, since for the low l multipoles there are mild discrepancies between the prediction of the power spectrum from the ΛCDM model and the actually observed values.

Finally a question which remains to be answered is: How does inflation end and reheating occur? In spite of its differences, scalar field inflation and spinor field inflation, have however several similar features. Just like in scalar field inflation, it is natural to define the end of inflation when the inflation parameters ϵ and η are of order unity. The interpretation of slow-roll inflation stays intact. The spinor field slowly rolls down its potential and near the minimum inflation ends. Once inflation is over, oscillations around the minimum of the potential begin. Next, one includes the decay of the spinorial inflaton particles into either fermion or bosons. A rapid decay, i.e. preheating, would require a decay into bosons since the Pauli exclusion principle would prevent decays into the same energy state. Once the non-relativistic particles decay into relativistic ones, the universe becomes radiation dominated, reheating. A detailed investigation of these processes is however being left for a future study.

In resume, non-standard spinors are a candidate to drive anisotropic inflation. The presence of the spin-connection in the matter part leads to additional couplings between the field and the geometry. Hence, inflation naturally becomes an anisotropic expansion, yielding a preferred direction which might have been detected as the axis of evil. Our derivation of the anisotropy corrected power spectrum is valid for all models in which the anisotropy can be treated as a perturbation around an isotropic background. Remarkably, while the usual inflationary features are obtained (low non-gaussianities, 60 e-folds, etc.), one finds that non-standard spinor driven inflation naturally results into a suppression of the lower multipoles of the CMB. This, in particular, cures the quadrupole anomaly that puzzles today's cosmological observations.

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