Effect of VSR invariant Chern-Simons Lagrangian on photon polarization

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Abstract. We propose a generalization of the Chern-Simons (CS) Lagrangian which is invariant under the SIM(2) transformations but not under the full Lorentz group. The generalized lagrangian is also invariant under a SIM(2) gauge transformation. We study the effect of such a term on radiation propagating over cosmological distances. We find that the dominant effect of this term is to produce circular polarization as radiation propagates through space. We use the circular polarization data from distant radio sources in order to impose a limit on this term.

Keywords: axions, particle physics - cosmology connection, high redshift galaxies

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1 Introduction

Symmetry is the guiding principle of modern physics. The Standard Model (SM), which is based on Lorentz invariance, provides a successful description of nature and has passed all experimental tests. However, it is believed to be the low energy limit of an ultimate theory at the Planck scale energy. At this scale, unification of quantum field theory with gravity leads to possible violation of Lorentz symmetry \([1]\). However it is difficult to probe Lorentz violation at LHC because the signal is suppressed by ratio of electroweak scale \((M_W)\) to Planck mass \((M_{Pl})\), i.e. \(M_W/M_{Pl} \approx 10^{-17}\) \([2]\). In a supersymmetric (SUSY) theory the signal may be suppressed more strongly as \((M_{SUSY}/M_{Pl})^2\), where \(M_{SUSY}\) is the scale of SUSY breaking \([3, 4]\).

Several terrestrial \([5]\), astrometric \([6, 7]\) and astrophysical tests have been conducted on Lorentz violation which impose stringent limits on its violation. Different theoretical models have been proposed on possible depatures from Lorentz invariance \([1, 3, 4, 8–15]\). Colladay and Kostelecky \([2]\) consider the Standard Model Extension (SME) in which Lorentz symmetry is spontaneously violated. Carroll et al. \([16]\) consider a CS term in 3+1 dimension, which is gauge invariant but breaks Lorentz invariance. The authors introduce an external four vector which breaks Lorentz invariance. This term is local and rotates the plane of polarization of photon due to different velocities of left- and right-circularly polarized photon.

We consider SIM(2) invariant CS term which respects SIM(2) gauge invariance but breaks Lorentz invariance. SIM(2) is the proper subgroup of Lorentz group as developed by Cohen and Glashow \([17]\) and termed Very Special Relativity (VSR). Adjoining one of the discrete symmetry such as P, T, CP or CT with Lorentz subgroup enlarges it to the full Lorentz group. In this paper, we introduce nonlocal operator \(\frac{n_\alpha}{n_\beta} \partial\) in the CS term in 3+1 dimension. This operator violates Lorentz invariance while respecting SIM(2) invariance. As we shall see, the modified nonlocal CS term splits the photon into two different polarization states which travel with different velocities. This implies violation of parity and Lorentz invariance in the theory. We find that the dominant effect of this term is to generate circular polarization as the electromagnetic wave travels through space. Furthermore the effect is dominant at low frequencies. Using the circular polarization data of radio sources from the MOJAVE (Monitoring of Jets in active galactic nuclei (AGN) with Very Long Baseline Array (VLBA) Experiments) program \([18]\), we impose a limit on the Lorentz violating parameter.

This paper is organised as follows: in section 2 we present the VSR invariant CS Lagrangian and derive the resulting photon dispersion relation. We then obtain the formulas for the Stokes parameters in this model. In section 3 we extract the Lorentz invariance violating parameter by the standard \(\chi^2\) minimization procedure. Finally, we conclude in section 4.
2 Theory

The well established Maxwell’s theory of electrodynamics is based on gauge and Lorentz invariance. The Lagrangian density for massless photons is given by

\[ \mathcal{L} = -\frac{1}{4} F_{\alpha\beta} F^{\alpha\beta}, \]  

(2.1)

where \( F_{\alpha\beta} = \partial_\alpha A_\beta - \partial_\beta A_\alpha \) is electromagnetic tensor. Eq. (2.1) is invariant under the gauge transformation \( A_\alpha \to A_\alpha + \partial_\alpha \varphi \). A photon mass term breaks the gauge invariance of the Lagrangian and experimental data imposes a stringent limit on this term.

We propose a \( \text{SIM}(2) \) invariant nonlocal CS term which can be written as,

\[ \mathcal{L}'_{\text{cs}} = \frac{\Gamma^2}{2} n_\alpha \partial_\beta \tilde{F}^{\alpha\beta}. \]  

(2.2)

Here \( \tilde{F}^{\alpha\beta} \) is dual electromagnetic tensor, \( \Gamma \) is a parameter of dimension mass and \( n_\alpha = (1, 0, 0, 1) \). This is manifestly not Lorentz invariant [19] but is invariant under \( \text{SIM}(2) \) transformations. The corresponding generators are \( T_1 = K_x + J_y, T_2 = K_y - J_x \), rotations \( (J_z) \) and boosts \( (K_z) \) about z-axis. Here \( J \) and \( K \) are the generators of rotations and boosts respectively. Under a boost along z-axis \( (K_z) \), the preferred vector \( n \) transforms as \( n_\alpha \to e^{\delta} n_\alpha \).

However \( n_\alpha n \cdot \partial \) is homogeneous in \( n \), and hence is \( \text{SIM}(2) \) invariant. We handle the nonlocal term \( 1/n \cdot \partial \) by using the prescription given in [20].

2.1 Gauge invariance

Under the gauge transformation \( A_\alpha \to A_\alpha + \partial_\alpha \varphi \), the variation of \( \text{SIM}(2) \) modified CS term is

\[ \Delta \mathcal{L}'_{\text{cs}} = \frac{\Gamma^2}{2} n_\alpha \partial_\beta \tilde{F}^{\alpha\beta}. \]  

(2.3)

where we have used the fact that \( n_\alpha \) is a constant vector and \( \partial_\beta \tilde{F}^{\alpha\beta} = 0 \). The remaining term is a surface term and would normally give null contribution to the action. However in the present case, this term also involves the inverse of \( n \cdot \partial \). We find that this term vanishes in all cases except if the derivative \( \partial_\beta \) is taken in the direction of \( n_\beta \). In this case the derivative operation cancels with the inverse operator \( 1/n \cdot \partial \) and we obtain a finite contribution. In order to eliminate this violation of gauge invariance we impose the constraint,

\[ n \cdot A = 0. \]  

(2.4)

This constraint is not invariant under Lorentz transformations but is invariant under \( \text{SIM}(2) \) transformations [20]. Hence we can consistently impose it within our framework. We also point out that the vector field \( A_\mu \) does not form an irreducible representation of \( \text{SIM}(2) \) and hence it is not necessary for us to work with the full vector field. In fact the vector potential can be split into four independent one dimensional \( \text{SIM}(2) \) fields [20].

Alternatively we use the modified form of electrodynamics which satisfies \( \text{SIM}(2) \) invariance but not the full Lorentz invariance [21, 22]. Here we briefly review this formalism. The gauge transformation of the vector field gets an additional contribution. It can be expressed as,

\[ \delta A_\mu = \partial_\mu \varphi + \frac{\Gamma^2}{2} n_\mu \left[ (n \cdot \partial)^{-1} \varphi \right]. \]  

(2.5)
The electromagnetic field tensor gets modified to [21, 22],

\[ F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - \frac{\Gamma^2}{2} n_\nu \left[(n \cdot \partial)^{-2}(n \cdot A)_\mu\right] + \frac{\Gamma^2}{2} n_\mu \left[(n \cdot \partial)^{-2}(n \cdot A)_\nu\right]. \tag{2.6} \]

In ref. [21, 22], the authors argue that these extra terms get eliminated by a redefinition of the gauge field,

\[ A_\mu \to A_\mu - \frac{1}{2} \Gamma^2 n_\mu \left[(n \cdot \partial)^{-2}(n \cdot A)\right]. \tag{2.7} \]

Hence these do not lead to any physical consequences. The formalism, however, allows a gauge invariant mass term of the gauge fields. In our analysis we set this mass term to zero. This involves an additional parameter whose physical consequences can be studied independently.

Let us now consider the following SIM(2) invariant nonlocal CS term:

\[ L_{cs} = \frac{\Gamma^2}{2} n_\alpha n \cdot \partial A_\beta \tilde{F}^{\alpha\beta} + \partial_\alpha A_\beta \tilde{F}^{\alpha\beta}. \tag{2.8} \]

This Lagrangian is invariant under the gauge transformations given by equation (2.5). These two terms individually are not invariant but the extra terms cancel among the two. We also see that these terms do not change under the redefinition of the gauge field given by equation (2.7). Hence we can ignore the factors proportional to \( \Gamma^2 \) in the formula for field tensor given by equation (2.6). Furthermore the second term on the right hand side of equation (2.8) does not contribute to the equations of motion. Hence, the equations of motion we obtain are same as those obtained from the standard electrodynamics along with the extra term given by eq. (2.2).

We point out that besides the mass term introduced in [21, 22], the term proposed in equation (2.8) is the only term that we can construct at this order which satisfies VSR gauge invariance. All other terms would require higher powers of the non-local operator \( 1/(n \cdot \partial) \) and hence a parameter of mass dimension larger than 2. We do not consider these additional terms in our analysis.

### 2.2 Propagation over cosmological distances

The complete Lagrangian density in presence of conserved current \( J^\alpha \) is

\[ \mathcal{L}_T = -\frac{1}{4} F_{\alpha\beta} F^{\alpha\beta} - J^\alpha A_\alpha + \frac{\Gamma^2}{2} n_\alpha \partial \tilde{F}^{\alpha\beta}. \tag{2.9} \]

As discussed above, the term \( \partial_\alpha A_\beta \tilde{F}^{\alpha\beta} \) in eq. (2.8) does not contribute to the equations of motion and, hence, is not included here. The equation of motion for the Lagrangian density \( \mathcal{L}_T \) is

\[ \partial_\beta F^{\beta\alpha} = J^\alpha - \frac{\Gamma^2}{n \cdot \partial} \frac{n_\lambda}{n \cdot \partial} \tilde{F}^{\lambda\alpha} \tag{2.10} \]

where we keep terms only up to order \( \Gamma^2 \). The modified Maxwell’s equations from eq. (2.10) are given by

\[ \nabla \cdot \mathbf{E} = \rho + \Gamma^2 \frac{1}{n \cdot \partial} \mathbf{n} \cdot \mathbf{B} \tag{2.11a} \]

\[ -\partial_t \mathbf{E} + \nabla \times \mathbf{B} = \mathbf{J} + \Gamma^2 \frac{1}{n \cdot \partial} (\mathbf{B} - \mathbf{n} \times \mathbf{E}). \tag{2.11b} \]
Here \( \mathbf{n} \) is a unit vector along the \( z \) axis. The two homogeneous Maxwell’s equations are

\[
\begin{align*}
\nabla \cdot B &= 0 \quad \text{(2.12a)} \\
\nabla \times E &= -\frac{\partial B}{\partial t} \quad \text{(2.12b)}
\end{align*}
\]

Using eqs. (2.12b) and (2.11), the source free wave equation takes the form

\[
\nabla(\nabla \cdot E) - \nabla^2 E + \frac{\partial^2 B}{\partial t^2} = \Gamma^2 \frac{\partial}{\partial t} \left( \frac{1}{\mathbf{n} \cdot \partial} \mathbf{B} \right) - \Gamma^2 \frac{\partial}{\partial t} \left( \frac{1}{\mathbf{n} \cdot \partial} \mathbf{n} \times \mathbf{E} \right). \quad \text{(2.13)}
\]

The operator \( \frac{1}{\mathbf{n} \cdot \partial} \) becomes,

\[
\frac{1}{\mathbf{n} \cdot \partial} = \frac{1}{\partial_t + \partial_z} = \int dt_+, \quad \text{(2.14)}
\]

where \( t_+ = \frac{t + z}{2} \). Eq. (2.13) further simplifies to

\[
(\omega^2 - k^2) \mathbf{E} + (k \cdot \mathbf{E}) \mathbf{k} = \frac{\Gamma^2}{\omega - k \cos \theta} (k \times \mathbf{E} - \omega \mathbf{n} \times \mathbf{E}). \quad \text{(2.15)}
\]

In the present case, for the source free modified Maxwell’s equation \( k \cdot \mathbf{E} \neq 0 \), so the longitudinal component of the photon polarization (proportional to \( \Gamma^2 \)) is not zero. It acquires a small value compared to the transverse components. Since the vector \( \mathbf{n} \) is along the \( z \)-direction and the wave propagation vector, \( \mathbf{\hat{k}} \), makes an angle \( \theta \) with \( \mathbf{n} \), the electric field can be expressed as,

\[
\mathbf{E} = E_x \mathbf{\hat{x}} + E_{yz} \mathbf{\hat{y}} + E_k \mathbf{\hat{k}}, \quad \text{(2.16)}
\]

where

\[
\mathbf{\hat{p}} = -\mathbf{\hat{y}} \cos \theta + \mathbf{\hat{z}} \sin \theta \quad \text{(2.17a)}
\]

\[
\mathbf{\hat{k}} = \mathbf{\hat{z}} \cos \theta + \mathbf{\hat{y}} \sin \theta \quad \text{(2.17b)}
\]

\[
\mathbf{\hat{n}} = \mathbf{\hat{z}}. \quad \text{(2.17c)}
\]

Comparing the \( x \)-, \( y \)- and \( z \)- component of eq. (2.15), we get

\[
(\omega^2 - k^2) E_x = \frac{\Gamma^2}{(\omega - k \cos \theta)} \left\{ (k - \omega \cos \theta) E_{yz} + \omega \sin \theta E_k \right\} \quad \text{(2.18a)}
\]

\[
\Gamma^2 E_x = (\omega^2 - k^2) \cos \theta E_{yz} - \omega^2 \sin \theta E_k \quad \text{(2.18b)}
\]

\[
\frac{\Gamma^2 k \sin \theta}{\omega - k \cos \theta} E_x = -(\omega^2 - k^2) \sin \theta E_{yz} - \omega^2 \cos \theta E_k. \quad \text{(2.18c)}
\]

Using eq. (2.18), we get the following dispersion relation

\[
\omega^2 - k^2 = \pm i \Gamma^2. \quad \text{(2.19)}
\]

Here the + and − signs correspond to right- and left-handed polarized photons. Hence we find that the modified SIM(2) invariant CS term leads to different dispersion relations for the right and left handed polarizations.
From eq. (2.18), the relation between the components of electric fields for the two solutions in eq. (2.19) are given by

\[
E_x \approx i\frac{(\omega - k \cos \theta)}{(k - \omega \cos \theta)} E_{yz} \tag{2.20a}
\]

\[
E_k \approx \frac{(\omega^2 - k^2) \sin \theta}{\omega^2(1 - \cos \theta)} E_{yz} \tag{2.20b}
\]

where we have kept only the leading order terms in \(\Gamma^2\). Using eq. (2.19), we find that the two eigenmodes of propagation are,

\[
|E_+\rangle \approx \begin{pmatrix} (\omega - k \cos \theta) \\ i \\ i(\omega^2 - k^2) \sin \theta \\ \omega^2(1 - \cos \theta) \end{pmatrix}, \quad |E_-\rangle \approx \begin{pmatrix} (\omega - k \cos \theta) \\ -i \\ -i(\omega^2 - k^2) \sin \theta \\ \omega^2(1 - \cos \theta) \end{pmatrix}. \tag{2.21}
\]

In the limit \(\omega \approx k\) the corresponding eigenvalues are given by

\[
k \approx \frac{i\Gamma^2}{2\omega} \equiv k_\pm \tag{2.22}
\]

where we have used eq. (2.19).

In the limit, \(\Gamma \rightarrow 0\), the two vectors in eq. (2.21) correspond to the left and right circular polarizations, i.e.

\[
|E_+\rangle = \begin{pmatrix} 1 \\ i \\ 0 \\ 0 \end{pmatrix}, \quad |E_-\rangle = \begin{pmatrix} 1 \\ -i \\ 0 \\ 0 \end{pmatrix}. \tag{2.23}
\]

Eq. (2.22) implies that two polarization modes travel with different velocity which is an indication of parity violation. This model leads to a significant contribution to circular polarization at low frequencies, which we will discuss in the next section.

Let us label the axes along \(\hat{x}, \hat{p}\) and \(\hat{k}\) by \(x_1, x_2\) and \(x_3\) respectively. Hence our wave is propagating along the \(x_3\) direction and the two transverse directions are taken to be along \(x_1\) and \(x_2\). An electric state vector, at any given time, can be written as a linear combination of two state vectors given in eq. (2.21), i.e.

\[
|E(x_3, t)\rangle = \frac{E_+(0)}{\sqrt{2}} \begin{pmatrix} (\omega - k_+ \cos \theta) \\ i \\ i(\omega^2 - k_+^2) \sin \theta \\ \omega^2(1 - \cos \theta) \end{pmatrix} e^{i(k_+ x_3 - wt)} + \frac{E_-(0)}{\sqrt{2}} \begin{pmatrix} (\omega - k_- \cos \theta) \\ -i \\ -i(\omega^2 - k_-^2) \sin \theta \\ \omega^2(1 - \cos \theta) \end{pmatrix} e^{i(k_- x_3 - wt)} \tag{2.24}
\]

Here \(E_3(x_3, t)\) component of the state vector \(|E(x_3, t)\rangle\) is very small and we are interested in determining the change in the photon polarization in the plane perpendicular to the photon propagation. Using eq. (2.24), we obtain,

\[
E_1(x_3, t) = \frac{1}{\sqrt{2}} \left( P_+ E_+(0) e^{ik_+ x_3} + P_- E_-(0) e^{ik_- x_3} \right) e^{-i\omega t} \tag{2.25a}
\]

\[
E_2(x_3, t) = \frac{1}{\sqrt{2}} \left( iE_+(0) e^{ik_+ x_3} - iE_-(0) e^{ik_- x_3} \right) e^{-i\omega t} \tag{2.25b}
\]

where

\[
P_\pm = \frac{\omega - k_\pm \cos \theta}{k_\pm - \omega \cos \theta}. \tag{2.26}
\]
The electric vector rotates in the plane perpendicular to the direction of propagation, when the two photon polarization modes travel with different velocity. Hence, the polarization state of photon changes after propagation over a large distance. It can be determined by calculating the Stokes parameter $I, Q, U, V$.

We assume that the wave is unpolarized at source and calculate its polarization after propagation through a distance $x_3$. The Jones matrix for unpolarized electromagnetic wave is

$$J = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}.$$  \hspace{1cm} (2.27)

Using this as the initial condition, we calculate the Stokes parameters after the photon has travelled a distance $x_3$. These are given by,

$$I = 2 \cosh \left( \frac{\Gamma^2 x_3}{\omega} \right)$$  \hspace{1cm} (2.28a)

$$Q = 0$$  \hspace{1cm} (2.28b)

$$U = \frac{2 \Gamma^2 \cot \left( \frac{\theta}{2} \right) \sinh^2 \left( \frac{\Gamma^2 x_3}{2 \omega} \right)}{\omega^2}$$  \hspace{1cm} (2.28c)

$$V = 2 \sinh \left( \frac{\Gamma^2 x_3}{\omega} \right).$$  \hspace{1cm} (2.28d)

Keeping only the leading order in $\Gamma$, we obtain,

$$I = 2, \quad Q = 0, \quad U = 0, \quad V = \frac{2 \Gamma^2}{\omega} x_3.$$  \hspace{1cm} (2.29)

This implies that an initially unpolarized wave acquires circular polarization upon propagation. The polarization state does not depend upon the direction of propagation of photon with respect to the VSR preferred axis. From eq. (2.29), we obtain,

$$\xi \equiv \frac{V}{I} = \frac{\Gamma^2}{\omega} x_3.$$  \hspace{1cm} (2.30)

So far we have confined our analysis to a flat space-time. However we need to compute the change in polarization in an expanding Universe since we are interested in sources located at redshifts comparable to unity. We consider a spatially flat Universe. The propagation of electromagnetic wave follows the same equations as given above but with time replaced by conformal time and the distance $x_3$ replaced by comoving distance [23, 24]. Besides this the overall energy density in the wave decreases due to expansion. However this effect is not relevant for calculation of shift in polarization. Consider a source at a redshift $z$. Its comoving distance $x_3$ is given by [25]

$$x_3 = \frac{1}{a_0 H_0} \int_{\frac{1}{1+z}}^{1} \frac{dx}{x^2 \sqrt{\Omega_\Lambda + \Omega_M x^{-3}}},$$  \hspace{1cm} (2.31)

where $H_0$ is Hubble constant, $a_0$ is scale parameter at present epoch, $\Omega_\Lambda$ is the ratio of vacuum energy density to critical density and $\Omega_M$ is the ratio of non-relativistic matter density to critical density. Hence eq. (2.30) becomes

$$\xi \equiv \frac{V}{I} = \frac{\Gamma^2}{\omega a_0 H_0} \int_{\frac{1}{1+z}}^{1} \frac{dx}{x^2 \sqrt{\Omega_\Lambda + \Omega_M x^{-3}}},$$

$$= \beta \int_{\frac{1}{1+z}}^{1} \frac{dx}{x^2 \sqrt{\Omega_\Lambda + \Omega_M x^{-3}}}.$$  \hspace{1cm} (2.32)
where $\beta = \frac{r^2}{\omega_0 H_0}$ is dimensionless. We find that the circular polarization depends on the redshift of the source. Furthermore the polarization generated increases with decrease in frequency. Hence the effect is dominant at low frequencies, such as, radio waves. In comparison the dominant effect in the case of the local Lorentz violating term, studied in ref. [16] is to produce a frequency independent rotation of linear polarization. We point out that mixing of hypothetical pseudoscalars of very low mass with photons in a background magnetic field also generates circular polarization [26–30]. In this case the effect is found to increase with frequency and is limited by the stringent constraints that have been imposed on the circular polarization which may be generated at optical frequencies [31].

In more generality, the theoretical model can be expressed as,

$$\xi = \beta \int_{1+z}^{1} \frac{dx}{x^2 \sqrt{\Omega_\Lambda + \Omega_M x^{-3}}} + \xi_0$$

(2.33)

where we have added a constant $\xi_0$ to the model in order to allow for the possibility that the mean value of circular polarization over the sample is not zero. In the limit $\beta = 0$, $\xi_0$ is the mean value of circular polarization over the entire sample. For $\beta \neq 0$ we determine both the parameters, $\beta$ and $\xi_0$, by making a $\chi^2$ fit to data.

3 Results

In this section, we determine the best fit parameters by using the $\chi^2$ minimisation procedure for the circularly polarized light emitted from radio jets associated with AGN. Here we use the circular polarization data from the MOJAVE program, which contains 133 bright, mostly compact radio-loud AGN in the northern sky. With VLBA facilities, circular polarization of the AGN jet sample at 15 GHz with flux density greater than 1.5 Jy has been observed. We only consider sources with redshift $z$ greater than 0.25 for which local effects are absent. We find that after this cut only 102 sources remain in the data set.

We extract the Lorentz violating parameter ($\beta$) by minimizing $\chi^2$ for the model given in eq. (2.33). The resulting $\chi^2$ is compared with the null model $\xi = \xi_0$ in order to determine the significance of the fit. For the null model we find that the mean value $\xi_0 = 0.38$ with $\chi^2 = 347.6$. The Lorentz violating model, eq. (2.33), yields $\beta = 0.067 \pm 0.12$, $\xi_0 = 0.40$ with minimum value of $\chi^2 = 347.1$. Hence we do not find a significant signal of Lorentz violation in the data. The corresponding best fit is shown in figure 1. Using the extracted value of $\beta$ we find that the one sigma limit on the Lorentz violating parameter $\Gamma$ is $13 \times 10^{-29}$ GeV.
4 Conclusion

We have proposed a modified Chern-Simons term which breaks the full Lorentz invariance but is invariant under SIM(2) transformations. The term is non-local and depends on a preferred vector, $n^a$. The term also respects gauge invariance. The non-local CS term changes the dispersion relation of photon. Hence it changes the polarization of the electromagnetic waves travelling over large distances. We find that at leading order, an initially unpolarized picks up circular polarization upon propagation. We test the predicted signal by using the circular polarization data from distant radio galaxies. We do not find a significant signal of violation of Lorentz invariance and impose a stringent limit on the Lorentz violation parameter.

References

[1] J. Collins, A. Perez, D. Sudarsky, L. Urrutia and H. Vucetich, Lorentz invariance and quantum gravity: an additional fine-tuning problem?, *Phys. Rev. Lett.* **93** (2004) 191301 [gr-qc/0403053] [INSPIRE].

[2] D. Colladay and V.A. Kostelecky, Lorentz violating extension of the standard model, *Phys. Rev. D* **58** (1998) 116002 [hep-ph/9809524] [INSPIRE].

[3] S. Groot Nibbelink and M. Pospelov, Lorentz violation in supersymmetric field theories, *Phys. Rev. Lett.* **94** (2005) 081601 [hep-ph/0404274] [INSPIRE].

[4] P. Jain and J.P. Ralston, Supersymmetry and the Lorentz fine tuning problem, *Phys. Lett. B* **621** (2005) 213 [hep-ph/0502106] [INSPIRE].

[5] R.J. Kennedy and E.M. Thorndike, Experimental Establishment of the Relativity of Time, *Phys. Rev.* **42** (1932) 400 [INSPIRE].

[6] Z. Bay and J.A. White, Radar astronomy and the special theory of relativity, *Acta Phys. Acad. Sci. Hungar.* **51** (1981) 273.

[7] J. Müller and M.H. Soffel, A Kennedy-Thorndike experiment using LLR data, *Phys. Lett. A* **198** (1995) 71.

[8] V.A. Kostelecky and M. Mewes, Electrodynamics with Lorentz-violating operators of arbitrary dimension, *Phys. Rev. D* **80** (2009) 015020 [arXiv:0905.0031] [INSPIRE].

[9] R. Bluhm and V.A. Kostelecky, Spontaneous Lorentz violation, Nambu-Goldstone modes and gravity, *Phys. Rev. D* **71** (2005) 065008 [hep-th/0412320] [INSPIRE].

[10] M.S. Berger and V.A. Kostelecky, Supersymmetry and Lorentz violation, *Phys. Rev. D* **65** (2002) 091701 [hep-th/0112243] [INSPIRE].

[11] S.M. Carroll, J.A. Harvey, V.A. Kostelecky, C.D. Lane and T. Okamoto, Noncommutative field theory and Lorentz violation, *Phys. Rev. Lett.* **87** (2001) 141601 [hep-th/0105082] [INSPIRE].

[12] D. Colladay and V.A. Kostelecky, Cross-sections and Lorentz violation, *Phys. Lett. B* **511** (2001) 209 [hep-ph/0104300] [INSPIRE].

[13] D. Colladay and V.A. Kostelecky, CPT violation and the standard model, *Phys. Rev. D* **55** (1997) 670 [hep-ph/9703464] [INSPIRE].

[14] A.V. Kostelecky and N. Russell, Classical kinematics for Lorentz violation, *Phys. Lett. B* **693** (2010) 443 [arXiv:1008.5062] [INSPIRE].
[15] S.R. Coleman and S.L. Glashow, *High-energy tests of Lorentz invariance*, Phys. Rev. D 59 (1999) 116008 [hep-ph/9812418] [inSPIRE].

[16] S.M. Carroll, G.B. Field and R. Jackiw, *Limits on a Lorentz and Parity Violating Modification of Electrodynamics*, Phys. Rev. D 41 (1990) 1231 [inSPIRE].

[17] A.G. Cohen and S.L. Glashow, *Very special relativity*, Phys. Rev. Lett. 97 (2006) 021601 [hep-ph/0601236] [inSPIRE].

[18] D.C. Homan and M.L. Lister, *Mojave: monitoring of jets in agn with vlba experiments. 2. first-epoch 15-ghz circular polarization results*, Astron. J. 131 (2006) 1262 [astro-ph/0511838] [inSPIRE].

[19] A.G. Cohen and S.L. Glashow, *A Lorentz-Violating Origin of Neutrino Mass?*, hep-ph/0605036 [inSPIRE].

[20] J. Vohanka, *Gauge Theory and SIM(2) Superspace*, Phys. Rev. D 85 (2012) 105009 [arXiv:1112.1797] [inSPIRE].

[21] J. Alfaro and V.O. Rivelles, *Non Abelian Fields in Very Special Relativity*, Phys. Rev. D 88 (2013) 085023 [arXiv:1305.1577] [inSPIRE].

[22] J. Alfaro, P. González and R. Ávila, *Electroweak standard model with very special relativity*, Phys. Rev. D 91 (2015) 105007 [arXiv:1504.04222] [inSPIRE].

[23] S.M. Carroll and G.B. Field, *The Einstein equivalence principle and the polarization of radio galaxies*, Phys. Rev. D 43 (1991) 3789 [inSPIRE].

[24] P. Tiwari, *New limit insertion on pseudoscalar-photon mixing from WMAP Observations*, Phys. Rev. D 86 (2012) 115025 [arXiv:1207.0606] [inSPIRE].

[25] S. Weinberg, *Cosmology*, Oxford University Press, (2008).

[26] P. Sikivie, *Experimental Tests of the Invisible Axion*, Phys. Rev. Lett. 51 (1983) 1415 [Erratum ibid. 52 (1984) 695] [inSPIRE].

[27] L. Maiani, R. Petronzio and E. Zavattini, *Effects of Nearly Massless, Spin Zero Particles on Light Propagation in a Magnetic Field*, Phys. Lett. B 175 (1986) 359 [inSPIRE].

[28] G. Raffelt and L. Stodolsky, *Mixing of the Photon with Low Mass Particles*, Phys. Rev. D 37 (1988) 1237 [inSPIRE].

[29] P. Jain, S. Panda and S. Sarala, *Electromagnetic polarization effects due to axion photon mixing*, Phys. Rev. D 66 (2002) 085007 [hep-ph/0206046] [inSPIRE].

[30] S. Das, P. Jain, J.P. Ralston and R. Saha, *Probing dark energy with light: Propagation and spontaneous polarization*, JCAP 06 (2005) 002 [hep-ph/0408198] [inSPIRE].

[31] A. Payez, J.R. Cudell and D. Hutsemekers, *New polarimetric constraints on axion-like particles*, JCAP 07 (2012) 041 [arXiv:1204.6187] [inSPIRE].