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Charged tensor matter fields and Lorentz symmetry violation via spontaneous symmetry breaking

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Abstract. We consider a model with a charged vector field along with a Cremmer–Scherk–Kalb–Ramond (CSKR) matter field coupled to a $U(1)$ gauge potential. We obtain a natural Lorentz symmetry violation due to the local $U(1)$ spontaneous symmetry breaking mechanism triggered by the imaginary part of the vector matter. The choice of the unitary gauge leads to the decoupling of the gauge-KR sector from the Higgs–KR sector. The excitation spectrum is carefully analyzed and the physical modes are identified. We propose an identification of the neutral massive spin-1 Higgs-like field with the massive Z' boson of the so-called mirror matter models.

1 Introduction

Systematic studies of the relations between topological models in $(1+2)$D and the phenomenology of planar theories, like high-T_c superconductivity, have been taken into account since the formulation of the Chern–Simons theory [1,2]. One remarkable characteristic extracted from the dynamics of the high-T_c superconductors is the violation of the P- and the T-symmetries. This fact emphasizes its planar nature. As a matter of fact, topological models originating from a Chern–Simons term are restricted to the description of objects living in $(1+2)$D. This aspect has motivated the study of extensions of planar gauge theories and the mathematical properties underlined, such as the fractional statistics [3]. Relevant extensions are the complex Maxwell–Chern–Simons (MCS*) model in $(1+2)$D [4] and the complex Maxwell–Chern–Simons–Proca theory (MCSP*). The planar scenario and the dynamical mass generation of these two models were largely exploited in [5]. In this context, the physical investigation of topological dynamical aspects of complex matter vector fields may still be better explored, specially in $(1+3)$D high-energy physics as well as in condensed matter systems.

The concept of topological models in $(1+3)$D has for the first time been pointed out by Cremmer and Scherk [6] and Kalb and Ramond [7]; we refer to this class of models as CSKR models. In these works, a topological term in $(1+3)$D is introduced which is an extension to the well-known Chern–Simons term in $(1+2)$D. This topological CSKR term introduces a direct coupling between a 1-form gauge field and another 2-form gauge field without an effective contribution to the energy and the momentum of the model; it however gives a mass contribution at tree level. The CSKR, as a topological model, is another candidate to generate mass without introducing a Higgs scalar field into the Lagrangian. So, gauge symmetry is preserved and a massive spin-1 boson appears. Consequently we can inquire whether a charged spin-1 vector boson could be incorporated into the spectrum of the CSKR model. Incidentally, vector-tensor field models have been largely studied, particularly in the context of $N=2$ supersymmetric models [8]. On the other hand, the contraction of a Chern–Simons term with a fixed vector has been proposed in order to build up a Lorentz violating model to describe astrophysical effects and cosmological new perspectives (possibly from geometrical origin) due to the variation of the universal constants [9]. Indeed, the possibility of spontaneous Lorentz and CPT-violation in string theory have been explored in [10], and more recent works have suggesting models in field theory to study Lorentz violation: Colladay and Kostelecký [11,12] suggested a general Lorentz violation extension of the standard model including CPT-even and CPT-odd terms in $(1+3)$D. They have obtained the result that the extension presents gauge invariance and a conserved energy-momentum tensor while covariance under particle rotations and boosts is broken. Coleman and Glashow [13,14] have also investigated tiny non-invariant terms introduced into the standard model Lagrangian in a perturbative framework. The effects of these perturbations increase rapidly with the energy for a preferred frame which implies Lorentz violation of the system. The occur-

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rence of dynamical breaking of the Lorentz symmetry in Abelian vector field models with the Wess–Zumino interaction have been explored by Andrianov and Soldati [15,16]. On the other hand, Carrol, Field and Jackiw [17] have demonstrated that ordinary Chern–Simons terms, studies previously in $(1+2)$D, can couple to dual electromagnetic tensor to a fixed and external four-vector. The effects of these Chern–Simons terms in Lorentz and CPT-violation have also been treated by Jackiw and Kostelecký [18].

The aim of this work is to study a charged vector-tensor matter-field model based on the complex extension of the CSKR model. We build up a full Lagrangian model where all the possible invariant terms are included. Furthermore, we add up a local $U(1)$ symmetry in order to have an interacting charged vector-tensor field model where a gauge field A_μ mediates the interaction. In this model, the 1- and 2-form fields can coexist and interact with each other by means of a topological term in $(1+3)$D. To check the consistency in a quantum field-theoretic sense, we discuss aspects such as causality and unitarity of the excitation spectrum. To this aim, we take into account the local $U(1)$ interaction formulation and we will analyze the vacuum states in the low-energy limit.

Due to the vector nature of the order parameter of the model, the ground state is identified with a constant four-vector (which we call b_μ) which implies an anisotropy of the vacuum state as a by-product and naturally induces Lorentz symmetry violation. This vacuum anisotropy has recently received much attention in connection with astrophysical phenomena [19]. We study the role played by the vector b_μ and its consequences to the physical degrees of freedom described in the spectrum of the model. We also explore the possible consistent (no ghost and no tachyon) choices of this vector background.

The outline of our paper is as follows: in Sect. 2, we introduce the full global $U(1)$ vector-tensor matter-field model and obtain the equations of motion, the Noether and the topological currents. In Sect. 3, using the hint of accommodating the two kind of currents in a doublet, we compute the propagators, poles and the physical consistency relation that is obeyed. In Sect. 4, we switch on an interacting gauge field and introduce local $U(1)$ symmetry. We study the SSB mechanism and conclude that the potential achieves its minimum for a non-vanishing vacuum expectation value of the charged vector-matter field. We adopt the unitary gauge and, in Sect. 5, we study the spectrum and consistency relations for the gauge-KR sector. In Sect. 6, the Higgs–KR sector is analyzed. Finally, in Sect. 7, we discuss and comment on our results.

2 The global $U(1)$ vector-tensor field model

Based on the CSKR model, we propose to study a full $U(1)$ charged vector-tensor matter-field model [6,7] where we have also included the topological terms. It can be written down as

$$
\begin{aligned}
\mathcal{L} = {} & \frac{1}{3}G^*_{\mu\nu\kappa}G^{\mu\nu\kappa} - \frac{1}{2}F^*_{\mu\nu}F^{\mu\nu} - (\partial_\mu B^\mu)^*(\partial_\nu B^\nu) \\
& +2(\partial_\mu H^{\mu\nu})^*(\partial^\rho H_{\rho\nu}) + \alpha^2 B^*_\mu B^\mu + \lambda(B^*_\mu B^\mu)^2 + \\
minus? \quad & -\beta^2 H^*_{\mu\nu}H^{\mu\nu} + m\epsilon^{\mu\nu\rho\sigma}B^*_\mu\partial_\nu H_{\rho\sigma} \\
& +m\epsilon^{\mu\nu\rho\sigma}B_\mu\partial_\nu H^*_{\rho\sigma},
\end{aligned}
\tag{1}
$$

where B^μ and $H^{\mu\nu}$ are the matter vector and tensor fields respectively; α and β represent the mass term parameters of the fields, λ represent the self-interacting parameter and m the topological mass[1]. The field strengths can be defined by

$$
\begin{aligned}
G_{\mu\nu\kappa} &= \partial_\mu H_{\nu\kappa} + \partial_\nu H_{\kappa\mu} + \partial_\kappa H_{\mu\nu}, \\
F_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu.
\end{aligned}
\tag{2}
$$

We observe that the topological term is a mixed one formed by B_μ and $H_{\mu\nu}$ in 4D as the term studied in [6,7]. Consequently, m is regarded as a topological mass. We also consider a potential term that defines the quadratic mass parameters α^2 and β^2. The conserved matter current J^μ stemming from the global $U(1)$ symmetry is given by

$$
\begin{aligned}
J^\mu = {} & \mathrm{i}(B_\nu F^{\mu\nu *} - B^*_\nu F^{\mu\nu}) - \mathrm{i}(H^*_{\nu\kappa}G^{\mu\nu\kappa} - H_{\nu\kappa}G^{\mu\nu\kappa *}) \\
& +\mathrm{i}m\epsilon^{\mu\nu\kappa\lambda}(B^*_\nu H_{\kappa\lambda} - B_\nu H^*_{\kappa\lambda}) \\
& +\mathrm{i}[B^\mu(\partial_\nu B^\nu)^* - B^{\mu *}(\partial_\nu B^\nu)] + \\
& -\mathrm{i}[(\partial^\rho H_{\rho\nu})^* H^{\mu\nu} - (\partial^\rho H_{\rho\nu})H^{\mu\nu *}].
\end{aligned}
\tag{3}
$$

The coupled Euler–Lagrange equations are

$$
\begin{aligned}
\Box B^\nu &= -\alpha^2 B^\nu - \lambda(B^*_\mu B^\mu)B^\nu - m\epsilon^{\nu\kappa\lambda\rho}\partial_\kappa H_{\lambda\rho}, \\
\Box H^{\nu\kappa} &= -\beta^2 H^{\nu\kappa} + m\epsilon^{\nu\kappa\lambda\rho}\partial_\lambda B_\rho.
\end{aligned}
\tag{4}
$$

According to the symmetry of the Lagrangian (1), we notice the occurrence of a B^4-interaction term that determines an anti-symmetrized identically conserved topological current of type

$$
J^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\kappa\lambda}\partial_\kappa B_\lambda. \tag{5}
$$

Its associated topological-vector charge gives rise to "vector solitons" solution whose value may be regarded as a quantum number. Indeed, this topological current induces directly the non-linear behavior on the vector-matter-field sector. The investigation of the non-linear dynamics and the non-trivial configuration of the fields with anisotropic energy in $(1+2)$D will be explored in a forthcoming work.

3 The spectrum analysis

In order to verify the physical spectrum, we rearrange the Lagrangian (1) in a linearized form, or

$$
\mathcal{L} = \mathbf{V}^{\mathrm{t}}\mathcal{O}\mathbf{V}, \tag{6}
$$

where $\mathcal{O}$ is a unitary wave operator and we represent $\mathbf{V}$ as a vector-tensor duplet, or

$$
\mathbf{V} = \begin{pmatrix} B_\mu \\ H_{\mu\nu} \end{pmatrix}, \tag{7}
$$

[1] We also adopt the metric $(+,-,-,-)$.

To obtain the propagators by means of the usual mechanism, we take the $\mathcal{O}^{-1}$ using the usual product algebra of the ordinary longitudinal, transverse and spin projector operators, which respectively are $\omega_{\mu\nu}$, $\theta_{\mu\nu}$ and $s_{\mu\nu\lambda} = \epsilon^{\gamma}{}_{\mu\nu\lambda}\partial_{\gamma}$. In addition, we also have the anti-symmetric longitudinal and transverse spin four-indexed projector operators written down as follws:

$$\theta_{\mu\nu,\lambda\rho} = \frac{1}{2}(\theta_{\mu\lambda}\theta_{\nu\rho} - \theta_{\mu\rho}\theta_{\nu\lambda}),$$
$$\omega_{\mu\nu,\lambda\rho} = \frac{1}{2}(\eta_{\mu\lambda}\omega_{\nu\rho} - \eta_{\mu\rho}\omega_{\nu\lambda} + \eta_{\nu\rho}\omega_{\mu\lambda} - \eta_{\nu\lambda}\omega_{\mu\rho}),$$
$$\eta_{\mu\nu,\lambda\rho} = \frac{1}{2}(\eta_{\mu\lambda}\eta_{\nu\rho} - \eta_{\mu\rho}\eta_{\nu\lambda}), \quad (8)$$

which implies a closed algebra, such that

	$\theta^{\alpha\beta}{}_{,\lambda\rho}$	$\omega^{\alpha\beta}{}_{,\lambda\rho}$	$s^{\alpha\beta}{}_{\lambda}$
$\theta_{\mu\nu,\alpha\beta}$	$\theta_{\mu\nu,\lambda\rho}$	0	$s_{\mu\nu\lambda}$
$\omega_{\mu\nu,\alpha\beta}$	0	$\omega_{\mu\nu,\lambda\rho}$	0
$s_{\mu\alpha\beta}$	$s_{\mu\lambda\rho}$	0	$-\Box\theta_{\mu\lambda}$

,

where we have obtained the result that $\eta_{\mu\nu,\lambda\rho} = \omega_{\mu\nu,\lambda\rho} + \theta_{\mu\nu,\lambda\rho}$. We obtain the propagators, in momentum space, which can be written down as follows:

$$\langle B_{\mu}^{*}, B_{\nu}\rangle = \frac{\mathrm{i}}{(k^2-\alpha^2)}\omega_{\mu\nu} + \frac{\mathrm{i}(k^2-\beta^2)}{(k^2-\mu_+^2)(k^2-\mu_-^2)}\theta_{\mu\nu},$$
$$\langle B_{\mu}^{*}, H_{\nu\lambda}\rangle = \langle H_{\mu\nu}^{*}, B_{\lambda}\rangle = \frac{\mathrm{i}m}{(k^2-\mu_+^2)(k^2-\mu_-^2)}s_{\mu\nu\lambda},$$
$$\langle H_{\mu\nu}^{*}, H_{\lambda\rho}\rangle = \frac{\mathrm{i}}{(k^2-\beta^2)}\omega_{\mu\nu,\lambda\rho} + \frac{\mathrm{i}(k^2-\alpha^2)}{(k^2-\mu_+^2)(k^2-\mu_-^2)}\theta_{\mu\nu,\lambda\rho}, \quad (9)$$

where

$$\mu_{\pm}^2 = \frac{\alpha^2+\beta^2+2m^2 \pm \sqrt{(\alpha^2+\beta^2+2m^2)^2 - 4\alpha^2\beta^2}}{2}, \quad (10)$$

which can be easily verified to be real and positive. As a consequence, we observe that the poles $k^2 = \alpha^2$, $k^2 = \beta^2$, $k^2 = \mu_+^2$ and $k^2 = \mu_-^2$ indicate the absence of tachyon states. Another point is the positivity of the norm of the states verified from the analysis of the residues of the propagators obtained. To do that, we take the transition amplitudes considering the doublet "vector-tensor current", which can be written down as follows:

$$\mathbf{J} = \begin{pmatrix} J_{\mu} \\ J_{\mu\nu} \end{pmatrix}, \quad (11)$$

where J_{μ} is the usual Noether current, and $J_{\mu\nu}$ is the current given in (3) and (5) which is conserved by definition.

We observe that there are two dynamical physical poles, both describing massive particles specified by μ_+^2 and μ_-^2. The transverse topological sectors are non-dynamical. We obtained the result that the propagators $\langle B_{\mu}^{*}, B_{\nu}\rangle$ and $\langle H_{\mu\nu}^{*}, H_{\lambda\rho}\rangle$ have the same poles and consequently the same particles. The crossing ones have no dynamics. We can see an order in the spectrum of the model which obeys the relation

$$\mu_+ > \beta > \alpha > \mu_-, \quad (12)$$

resulting in a consistent physical model. We observe that to perform the analysis of the degrees of freedom in 1 + 3 dimensions it was necessary to take the antisymmetric topological current $J_{\mu\nu}$ given in (5) to complete a doublet with the usual vector one (J_{μ}). Indeed, the topological current has induced directly the non-linear behavior of the vector-matter-field sector.

There is a delicate point we should comment on here. It concerns two important issues as far as field-theoretic consistency in under discussion, namely, unitarity and renormalizability. In our specific model, a propagating time-component of a charged vector-matter field is potentially a ghost, in the sense of having a non-positive defined scalar product in the Hilbert space of states. This kind of matter-ghost degrees of freedom must be properly discussed[2]. Indeed, this component is present in the self-interaction quartic term. Nevertheless, the very aim of this work is to study the possible origin of Lorentz symmetry breaking as due to vector and tensor matter fields. For the sake of gauge invariance, we had plugged into the action a dynamical longitudinal component of the B_{μ} field. It is a matter-type field, with the usual $U(1)$-phase transformation, and so this term is allowed, once the space-time derivative is suitable gauge-covariantized. The important question to be worked out further (we shall come back to this question at the end of Sect. 6) is to find a condition to ensure decoupling or non-propagation of the non-physical longitudinal components of B_{μ}. So, we finally state that the task of breaking Lorentz symmetry has been accomplished by the proposed vector and tensor matter fields. Next, the main issue is to ensure quantum-mechanical consistency of the proposed action.

4 The local $U(1)$ theory and SSB

To introduce interactions into the model represented by the matter Lagrangian (1), we take as local the symmetry phase, as the usual method, or

$$B_{\mu}^{'} = \mathrm{e}^{-\mathrm{i}\Lambda(x)}B_{\mu},$$

and

$$H_{\mu\nu}^{'} = \mathrm{e}^{-\mathrm{i}\Lambda(x)}H_{\mu\nu}. \quad (13)$$

In this way the symmetries are restored, introducing the invariant Lagrangian written as

$$\mathcal{L}_{int} = \frac{1}{3}\mathcal{G}_{\mu\nu\kappa}^{*}\mathcal{G}^{\mu\nu\kappa} - \frac{1}{2}\mathcal{F}_{\mu\nu}^{*}\mathcal{F}^{\mu\nu} - \frac{1}{4}f^{\mu\nu}f_{\mu\nu} +$$

[2] Subject of a forthcoming work

$$-(D_\mu B^\mu)^*(D_\nu B^\nu) + 2(D_\mu H^{\mu\nu})^*(D^\rho H_{\rho\nu})$$
$$+m\epsilon^{\mu\nu\kappa\lambda} B^*_\mu D_\nu H_{\kappa\lambda} + m\epsilon^{\mu\nu\kappa\lambda} B_\mu (D_\nu H_{\kappa\lambda})^*$$
$$+\alpha^2 B^*_\mu B^\mu + \lambda(B^*_\mu B^\mu)^2 - \beta^2 H^*_{\mu\nu} H^{\mu\nu}, \quad (14)$$

where $D_\mu = \partial_\mu + \mathrm{i}eA_\mu$ is the covariant derivative. The above Lagrangian describes an interaction model between the matter fields B_μ, $H_{\mu\nu}$ and the gauge field A_μ, where we define

$$\mathcal{G}_{\mu\nu\kappa} = D_\mu H_{\nu\kappa} + D_\nu H_{\kappa\mu} + D_\kappa H_{\mu\nu},$$
$$\mathcal{F}_{\mu\nu} = D_\mu B_\nu - D_\nu B_\mu, \quad (15)$$
$$f_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu.$$

The new conserved current can be written down as follows:

$$\partial_\mu f^{\mu\nu} = \mathcal{J}^\nu = \mathrm{i}(\mathcal{B}_\nu \mathcal{F}^{\mu\nu *} - \mathcal{B}^*_\nu \mathcal{F}^{\mu\nu}) +$$
$$-\mathrm{i}(H^*_{\nu\kappa}\mathcal{G}^{\mu\nu\kappa} - H_{\nu\kappa}\mathcal{G}^{\mu\nu\kappa *})$$
$$+\mathrm{i}m\epsilon^{\mu\nu\kappa\lambda}(B^*_\nu H_{\kappa\lambda} - B_\nu H^*_{\kappa\lambda})$$
$$+\mathrm{i}[B^\mu (D_\nu B^\nu)^* - B^{\mu *}(D_\nu B^\nu)] +$$
$$-\mathrm{i}[(D^\rho H_{\rho\nu})^* H^{\mu\nu} - (D^\rho H_{\rho\nu}) H^{\mu\nu *}], \quad (16)$$

where $\mathcal{J}^\mu$ is the covariant matter current.

As we have seen, the covariant interacting vector-tensor model described by the Lagrangian (14) introduces a $U(1)$ gauge field, A_μ. To explore the behavior of these fields at low-energy phenomenology, we take the mechanism of spontaneous breaking of the gauge symmetry. As the model carries vector and tensor fields as matter degrees of freedom, the discussion of the SSB becomes subtle. The quartic self-interacting non-linear term of the matter field B_μ in the Lagrangian (14) could play a role similar to the Higgs field, but with a vector nature. The $\lambda(B^*_\mu B^\mu)^2$ does not spoil the invariance of the Lagrangian under the group of local $U(1)$ transformations. So the condition of a minimum of the energy (E) can be obtained taking the minimum of the potential energy (V), or

$$\frac{\mathrm{d}E}{\mathrm{d}B_\mu} = \frac{\mathrm{d}V}{\mathrm{d}B_\mu} = \alpha^2 B^*_\mu + 2\lambda(B^*_\nu B^\nu)B^*_\mu = 0, \quad (17)$$

where it is analogous to the B^*_μ term, recalling that α^2 is a mass parameter. In this case, the situation where $\alpha^2 < 0$ and $\lambda > 0$ introduces a non-trivial vacuum, and it follows that the energy is a minimum at

$$B^*_\mu B^\mu = b_\mu b^\mu = b^2 = -\frac{\alpha^2}{2\lambda} u^2, \quad (18)$$

where we observe that, in this case, we require that b^2 be a constant 4-vector parameter such that $-\frac{\alpha^2}{2\lambda} > 0$. In fact, we observe that the VEV for the field B_μ is given by

$$\langle 0|B_\mu|0\rangle = b_\mu = \sqrt{\frac{-\alpha^2}{2\lambda}}\, u_\mu, \quad (19)$$

where u_μ is a unitary vector which lies in a fixed direction in space-time. In turn, it breaks the Lorentz symmetry, and due to this arbitrariness we have to choose amongst the possible types of vector: $u^2 = 1$ (time-like), $u^2 = -1$ (space-like) or $u^2 = 0$ (light-like), analogous to the case studied in [21]. As we are interested in non-trivial configurations of the fields, we exclude the light-like possibility. Consequently, we reach a non-trivial vacuum solution for an energy E which breaks spontaneously the $U(1)$ local symmetry and also violates the Lorentz symmetry. We emphasize that the Lorentz violation came along as a by-product of the internal symmetry breaking. The Lorentz violation has received much attention due to possible astrophysical and condensed matter effects [9, 22], which deserves a deeper analysis. In this work, we are going to verify the mass spectrum of this model. To this purpose, we begin by observing that the system under consideration has an infinite set of vacuum states, corresponding to points on a circle of radius given by (19) posed on the complex plane of the field B_μ. So we can decompose the complex fields into components and we shift the field B_μ along the real axis (analogous to the Higgs mechanism). So we have

$$B_\mu \to B_\mu + b_\mu = X_\mu + \mathrm{i}Y_\mu + b_\mu,$$
$$H_{\mu\nu} \to P_{\mu\nu} + \mathrm{i}Q_{\mu\nu}, \quad (20)$$

so, we can express the potential term by

$$V = \lambda(B^*_\mu B^\mu - b^2)^2 - \lambda b^4 - \beta^2 H^*_{\mu\nu} H^{\mu\nu}, \quad (21)$$

which are substituted into the Lagrangian (14), whose expansion we find to be

$$\mathcal{L}_{\text{broken}} = \frac{1}{3} P_{\mu\nu\kappa} P^{\mu\nu\kappa} + \frac{1}{3} Q_{\mu\nu\kappa} Q^{\mu\nu\kappa} - \frac{1}{2} X_{\mu\nu} X^{\mu\nu} +$$
$$-\frac{1}{2} Y_{\mu\nu} Y^{\mu\nu} - (\partial_\mu X^\mu)^2 - (\partial_\mu Y^\mu)^2$$
$$+2(\partial_\mu P^{\mu\nu})(\partial^\rho P_{\rho\nu}) + 2(\partial_\mu Q^{\mu\nu})(\partial^\rho Q_{\rho\nu}) +$$
$$-e^2 b^2 A_\mu A^\mu - \frac{1}{4} f_{\mu\nu} f^{\mu\nu} + 2e(b^\mu A^\nu)(\partial_\mu Y_\nu) +$$
$$-2e(A^\mu b^\nu)(\partial_\mu Y_\nu) + 2e(\partial_\mu Y^\mu)(A_\nu b^\nu)$$
$$+2m\epsilon^{\mu\nu\kappa\lambda} X_\mu \partial_\nu P_{\kappa\lambda} - 2em\epsilon^{\mu\nu\kappa\lambda} b_\mu A_\nu Q_{\kappa\lambda}$$
$$+4\lambda b_\mu b_\nu X^\mu X^\nu - \beta^2 P_{\mu\nu} P^{\mu\nu} - \beta^2 Q_{\mu\nu} Q^{\mu\nu}$$
$$+\text{higher order terms}, \quad (22)$$

where $X_{\mu\nu}$, $Y_{\mu\nu}$, $P_{\mu\nu\kappa}$ and $Q_{\mu\nu\kappa}$, are the field strengths of their respective real components of 1- and 2-form fields written in the definitions (20). We observe that the Lagrangian (22) is non-diagonal, making all this a subtle computation. The terms $2e(b^\mu A^\nu)(\partial_\mu Y_\nu)$, $-2e(A^\mu b^\nu)(\partial_\mu Y_\nu)$ and $2e(\partial_\mu Y^\mu)(A_\nu b^\nu)$ can be absorbed by carrying out the following field re-definitions:

$$A_\mu \to A_\mu - q_\nu(\partial_\mu Y^\nu) + q_\nu(\partial^\nu Y_\mu) + q_\mu(\partial_\nu Y^\nu),$$
$$f_{\mu\nu} \to f_{\mu\nu} + \gamma Y_{\mu\nu} + \Sigma_{\mu\nu}(\partial_\alpha Y^\alpha), \quad (23)$$

where q_ν, γ and $\Sigma_{\mu\nu}$ are operators that can easily be found by manipulating (23) and (22), which can be defined as

$$q^\nu = \frac{b^\nu}{eb^2}, \quad \gamma = q^\alpha \partial_\alpha, \quad \text{and} \quad \Sigma_{\mu\nu} = q_\mu \partial_\nu - q_\nu \partial_\mu. \quad (24)$$

So the resulting Lagrangian (22) is non-gauge invariant because the Lorentz symmetry is broken. It breaks translation due to the presence of the γ operator, and it breaks rotation by virtue of the $\Sigma_{\mu\nu}$ operator present in the new definitions in (23). Nevertheless, we can eliminate the Y_μ field by means of a gauge choice, picking a particular gauge parameter in the $U(1)$ phase transformation; so,

$$X'_\mu = X_\mu - \Lambda Y_\mu,$$
$$Y'_\mu = Y_\mu - \Lambda X_\mu + \Lambda b_\mu, \tag{25}$$

where Λ is an arbitrary gauge parameter. In fact, we can gauge away the Y_μ field choosing a particular gauge, bearing in mind the unitarity condition on the particle spectrum. Then the A_μ field acquires mass due to the presence of the scalar field parameter $\Lambda = \Phi$ in its longitudinal mode. This describes the associated Higgs mechanism for a complex vector. This can be seen through the following re-defined transformations:

$$A_\mu \to A_\mu - \partial_\mu \Phi, \qquad f'_{\mu\nu} = f_{\mu\nu}. \tag{26}$$

Then, the Lagrangian (22) can be rewritten in the absence of the interaction terms as,

$$\begin{aligned}\mathcal{L}_{\text{broken}} = {} & \frac{1}{3}P_{\mu\nu\kappa}P^{\mu\nu\kappa} + \frac{1}{3}Q_{\mu\nu\kappa}Q^{\mu\nu\kappa} - \frac{1}{2}X_{\mu\nu}X^{\mu\nu} + \\ & -(\partial_\mu X^\mu)^2 + 2(\partial_\mu P^{\mu\nu})(\partial^\rho P_{\rho\nu}) \\ & +2(\partial_\mu Q^{\mu\nu})(\partial^\rho Q_{\rho\nu}) - e^2b^2A_\mu A^\mu - \frac{1}{4}f_{\mu\nu}f^{\mu\nu} \\ & +2m\epsilon^{\mu\nu\kappa\lambda}X_\mu\partial_\nu P_{\kappa\lambda} - 2em\epsilon^{\mu\nu\kappa\lambda}b_\mu A_\nu Q_{\kappa\lambda} \\ & +4\lambda b_\mu b_\nu X^\mu X^\nu - \beta^2 P_{\mu\nu}P^{\mu\nu} - \beta^2 Q_{\mu\nu}Q^{\mu\nu}\end{aligned} \tag{27}$$

In order to extract the physical content of the Lagrangian (27) one can split it in two sectors,

$$\begin{aligned}\mathcal{L}_{\text{gauge-KR}} = {} & \frac{1}{3}Q_{\mu\nu\kappa}Q^{\mu\nu\kappa} - \frac{1}{4}f_{\mu\nu}f^{\mu\nu} \\ & +2(\partial_\mu Q^{\mu\nu})(\partial^\rho Q_{\rho\nu}) - e^2b^2A_\mu A^\mu + \\ & -2em\epsilon^{\mu\nu\kappa\lambda}b_\mu A_\nu Q_{\kappa\lambda} - \beta^2 Q_{\mu\nu}Q^{\mu\nu},\end{aligned} \tag{28}$$

$$\begin{aligned}\mathcal{L}_{\text{Higgs-KR}} = {} & \frac{1}{3}P_{\mu\nu\kappa}P^{\mu\nu\kappa} - \frac{1}{2}X_{\mu\nu}X^{\mu\nu} - (\partial_\mu X^\mu)^2 \\ & +2(\partial_\mu P^{\mu\nu})(\partial^\rho P_{\rho\nu}) + 2m\epsilon^{\mu\nu\kappa\lambda}X_\mu\partial_\nu P_{\kappa\lambda} \\ & +4\lambda b_\mu b_\nu X^\mu X^\nu - \beta^2 P_{\mu\nu}P^{\mu\nu},\end{aligned} \tag{29}$$

where we can observe from the Lagrangian (28) that the gauge field only interacts, via the topological term, with the imaginary part of the tensor KR field. We can also observe that the mass of the A_μ field depends on the vector b_μ, the broken parameter, and on the topological mass as well. On the other hand, the Lagrangian (29) indicates that the vector b^μ contributes to the mass of the real part of the original neutral meson field. We emphasize that the appearance of this new boson field does not prescribe any new symmetry group in the model. To verify the consistency, we are going to compute the spectral analysis separately.

5 The spectrum of the gauge-KR sector

A remarkable feature of the Lagrangian (28) is that it contains a massive gauge vector field (Proca) that interacts with a 2-form KR field. An analysis of the physical degrees of freedom requires one to deal with the unitary gauge (25). Then, assuming that the fields are well-behaved asymptotically, we can rearrange the Lagrangian considering a mixed doublet defined as $\mathbf{U}^{\text{t}} = (A_\mu, Q_{\mu\nu})$. So, the Lagrangian $\mathcal{L}_{\text{gauge-KR}} = \mathbf{U}^{\text{t}}\mathcal{O}\mathbf{U}$ where $\mathcal{O}$ can be easily written down from the Lagrangian (28). We remark that the re-defined currents preserve the unitarity as in the unbroken case. Consequently, we can obtain the propagators taking the inverse operator, $\mathcal{O}^{-1}$, and using the re-defined current doublet in the momentum space, as follows:

$$\begin{aligned}\langle A_\mu, A_\nu\rangle &= \frac{\text{i}}{(k^2+e^2b^2)}\Lambda_{\mu\nu} \\ &\quad + \frac{\text{i}(k^2-\beta^2)}{(k^2-\tau_+^2)(k^2-\tau_-^2)}\Omega_{\mu\nu}, \\ \langle A_\mu, Q_{\nu\lambda}\rangle &= \langle Q_{\mu\nu}, A_\lambda\rangle \\ &= \frac{\text{i}m}{(k^2-\tau_+^2)(k^2-\tau_-^2)}\Pi_{\mu\nu\lambda}, \\ \langle Q_{\mu\nu}, Q_{\lambda\rho}\rangle &= \frac{\text{i}}{(k^2-\beta^2)}\Lambda_{\mu\nu,\lambda\rho} \\ &\quad + \frac{\text{i}(k^2+e^2b^2)}{(k^2-\tau_+^2)(k^2-\tau_-^2)}\Omega_{\mu\nu,\lambda\rho},\end{aligned} \tag{30}$$

in these expressions above we have used the longitudinal, $\Lambda_{\mu\nu}$, and the transverse, $\Omega_{\mu\nu}$ and $\Pi_{\mu\nu\lambda}$, given by

$$\begin{aligned}\Lambda_{\mu\nu} &= \frac{b_\mu b_\nu}{b^2}, \\ \Omega_{\mu\nu} &= \eta_{\mu\nu} - \frac{b_\mu b_\nu}{b^2}, \\ \Pi_{\mu\nu\lambda} &= \epsilon^{\gamma}{}_{\mu\nu\lambda}b_\gamma,\end{aligned} \tag{31}$$

whose multiplicative table looks as follows now:

	$\Omega^\alpha{}_\nu$	$\Lambda^\alpha{}_\nu$	$\Pi^\alpha{}_{\nu\lambda}$
$\Omega_{\mu\alpha}$	$\Omega_{\mu\nu}$	0	$\Pi_{\mu\nu\lambda}$
$\Lambda_{\mu\alpha}$	0	$\Lambda_{\mu\nu}$	0
$\Pi_{\mu\alpha}{}^\lambda$	$\Pi_{\mu\nu}{}^\lambda$	0	$-b^2\Omega_{\mu\nu}$

We can also define the anti-symmetrized longitudinal and transverse operators as

$$\begin{aligned}\Omega_{\mu\nu,\lambda\rho} &= \frac{1}{2}(\Omega_{\mu\lambda}\Omega_{\nu\rho} - \Omega_{\mu\rho}\Omega_{\nu\lambda}), \\ \Lambda_{\mu\nu,\lambda\rho} &= \frac{1}{2}(\eta_{\mu\lambda}\Lambda_{\nu\rho} - \eta_{\mu\rho}\Lambda_{\nu\lambda} + \eta_{\nu\rho}\Lambda_{\mu\lambda} - \eta_{\nu\lambda}\Lambda_{\mu\rho}), \\ \eta_{\mu\nu,\lambda\rho} &= \frac{1}{2}(\eta_{\mu\lambda}\eta_{\nu\rho} - \eta_{\mu\rho}\eta_{\nu\lambda}),\end{aligned} \tag{32}$$

which implies the closed product algebra:

	$\Omega^{\alpha\beta}{}_{,\lambda\rho}$	$\Lambda^{\alpha\beta}{}_{,\lambda\rho}$	$\Pi^{\alpha\beta}{}_{\lambda}$
$\Omega_{\mu\nu,\alpha\beta}$	$\Omega_{\mu\nu,\lambda\rho}$	0	$\Pi_{\mu\nu\lambda}$
$\Lambda_{\mu\nu,\alpha\beta}$	0	$\Lambda_{\mu\nu,\lambda\rho}$	0
$\Pi_{\mu\alpha\beta}$	$\Pi_{\mu\lambda\rho}$	0	$-b^2\Omega_{\mu\lambda}$

,

where in addition we can observe that $\eta_{\mu\nu,\lambda\rho} = \Lambda_{\mu\nu,\lambda\rho} + \Omega_{\mu\nu,\lambda\rho}$. The mass values, $\tau^2_\pm$, are easily obtained:

$$\tau^2_\pm = \frac{e^2b^2 - \beta^2 \pm \sqrt{(\beta^2 + e^2b^2)^2 + 8m^2e^2b^2}}{2}$$

for the case where $u^2 = 1$,

$$\tau^2_\pm = \frac{-e^2b^2 - \beta^2 \pm \sqrt{(\beta^2 - e^2b^2)^2 - 8m^2e^2b^2}}{2}$$

for the case where $u^2 = -1$. (33)

For the case $u^2 = 1$ (time-like condition), only τ_+ defines a physical excitation. τ_- is a tachyonic excitation. On the other hand, for $u^2 = -1$ (space-like condition), both τ_+ and τ_- are physical excitations for the restricted values of the topological mass,

$$m < \frac{\beta^2 - e^2b^2}{\sqrt{8}eb}, \tag{34}$$

where $b = |b_\mu|$. Observe that β cannot be zero, and $\beta > e|b|$, with the result that the Lorentz symmetry breaking implies an anisotropy of the space-time is realized as a mass generation on the gauge field, similar to the Higgs mechanism whose consistency is ensured by the presence of a mass term for the 2-form $H_{\mu\nu}$ field.

6 The spectrum of the Higgs–KR sector

Now, we are going to verify the physical spectrum of the Higgs–KR sector. In a similar way as in the previous case, we are also going to analyze the consequences of the symmetry breaking on the spectrum of the Lagrangian (29). We can verify that due to the presence of an anisotropic space-time, the X_μ neutral vector field is deformed, which implies non-defined propagator poles. Therefore, the physical states can only be obtained by considering the projections of the X^μ field along and perpendicular to the b_μ vector. Then defining the W_μ and Z_μ as the parallel and transverse projections of the X_μ field respectively, or

$$W_\mu = \frac{b_\alpha X^\alpha}{b^2}\, b_\mu, \tag{35}$$

$$Z_\mu = X_\mu - \frac{b_\alpha X^\alpha}{b^2}\, b_\mu. \tag{36}$$

We can perform rotations at each point in deformed-space where the doublet $\mathbf{U}^{\mathrm{t}} = (X_\mu, P_{\mu\nu})$ can be taken as

$$\mathbf{U}^{\mathrm{t}} = (X_\mu, P_{\mu\nu}) \Rightarrow \tag{37}$$
$$\left\{\mathbf{U}(W)^{\mathrm{t}} = (W_\mu, P_{\mu\nu})_{\mathrm{l}}|\, W_\mu = \frac{b_\alpha X^\alpha}{b^2}\, b_\mu,\ Z_\mu = 0\right\},$$

$$\mathbf{U}^{\mathrm{t}} = (X_\mu, P_{\mu\nu}) \Rightarrow \tag{38}$$
$$\left\{\mathbf{U}(Z)^{\mathrm{t}} = (Z_\mu, P_{\mu\nu})_{\mathrm{p}}|\, Z_\mu = X_\mu - \frac{b_\alpha X^\alpha}{b^2}\, b_\mu,\ W_\mu = 0\right\},$$

where the sub-indexes l and p mean the longitudinal and the perpendicular projections, respectively. The doublets $\mathbf{U}(W)$ and $\mathbf{U}(Z)$ are orthogonal and may be chosen non-simultaneously. Then we can now rewrite the Lagrangian (29) for each one of the above situations, which for the $\mathbf{U}(W)$ case yields

$$\begin{aligned}\mathcal{L}(W) = &-P_{\mu\nu}(\eta^{\mu\nu,\lambda\rho}\Box)P_{\lambda\rho} - W_\mu(\eta^{\mu\nu}\Box)W_\nu \\ &+4\lambda b^2 W_\mu W^\mu - \beta^2 P_{\mu\nu}P^{\mu\nu} \\ &+m\epsilon^{\mu\nu\rho\sigma}W_\mu\partial_\nu P_{\rho\sigma} - m\epsilon^{\mu\nu\rho\sigma}P_{\rho\sigma}\partial_\nu W_\mu;\end{aligned} \tag{39}$$

hence, we can note that by substituting the re-definition of the W_μ field (35) into (29) the cross massive term $4\lambda b_\mu b_\nu X^\mu X^\nu$ is transformed to $4\lambda b^2 W_\mu W^\mu$ and so is given a mass term form. On the other hand, we can choose the X_μ field perpendicular to b_μ, where one can write (29) in the $\mathbf{U}(Z)$ case. We have

$$\begin{aligned}\mathcal{L}(Z) = &-P_{\mu\nu}(\eta^{\mu\nu,\lambda\rho}\Box)P_{\lambda\rho} + Z_\mu(\eta^{\mu\nu}\Box)Z_\nu + \\ &-\beta^2 P_{\mu\nu}P^{\mu\nu} + m\epsilon^{\mu\nu\rho\sigma}Z_\mu\partial_\nu P_{\rho\sigma} + \\ &-m\epsilon^{\mu\nu\rho\sigma}P_{\rho\sigma}\partial_\nu Z_\mu,\end{aligned} \tag{40}$$

where the mass term $4\lambda b_\mu b_\nu X^\mu X^\nu$ is gauged away. In order to verify the degrees of freedom of these two cases, we are going to deal with the Lagrangian expressions (39) and (40), and we use the analogous method of the gauge-KR sector. So we now have two operators: $\mathcal{O}(W)$, $\mathcal{O}(Z)$ and their respective inverses are straightforwardly obtained. The associated to $\mathcal{O}(W)^{-1}$ propagator is obtained as

$$\begin{aligned}\langle W_\mu, W_\nu\rangle &= \frac{\mathrm{i}}{(k^2-\bar{\alpha}^2)}\omega_{\mu\nu} \\ &\quad+\frac{\mathrm{i}(k^2-\beta^2)}{(k^2-\bar{\mu}_+^2)(k^2-\bar{\mu}_-^2)}\,\theta_{\mu\nu}, \\ \langle W_\mu, P_{\nu\lambda}\rangle &= \langle P_{\mu\nu}, W_\lambda\rangle \\ &= \frac{\mathrm{i}m}{(k^2-\bar{\mu}_+^2)(k^2-\bar{\mu}_-^2)}\, s_{\mu\nu\lambda}, \\ \langle P_{\mu\nu}, P_{\lambda\rho}\rangle &= \frac{\mathrm{i}}{(k^2-\beta^2)}\omega_{\mu\nu,\lambda\rho} \\ &\quad+\frac{\mathrm{i}(k^2-\bar{\alpha}^2)}{(k^2-\bar{\mu}_+^2)(k^2-\bar{\mu}_-^2)}\,\theta_{\mu\nu,\lambda\rho},\end{aligned} \tag{41}$$

where we have defined

$$\bar{\mu}^2_\pm = \frac{\bar{\alpha}^2 + \beta^2 + 2m^2 \pm \sqrt{(\bar{\alpha}^2 + \beta^2 + 2m^2)^2 - 4\bar{\alpha}^2\beta^2}}{2} \tag{42}$$

and $\bar{\alpha} = 4\lambda b^2$. In the same way, for the perpendicular wave operator $\mathcal{O}(Z)^{-1}$ this results in

$$\langle Z_\mu, Z_\nu\rangle = \frac{\mathrm{i}}{k^2}\omega_{\mu\nu} + \frac{\mathrm{i}(k^2-\beta^2)}{k^2(k^2-\xi^2)}\theta_{\mu\nu},$$

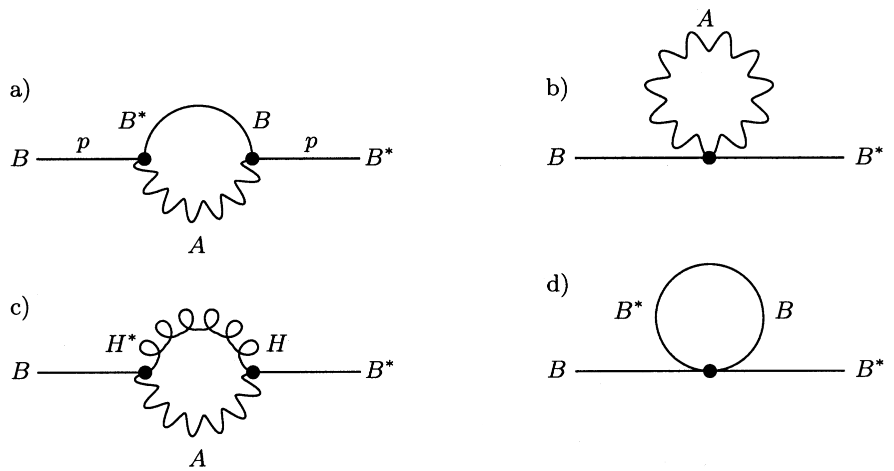

Fig. 1. One-loop corrections to the B-field self-energy. p stands for the external momentum carried by B

$$\langle Z_\mu, P_{\nu\lambda}\rangle = \langle P_{\mu\nu}, Z_\lambda\rangle = \frac{m}{k^2(k^2-\xi^2)} s_{\mu\nu\lambda}, \tag{43}$$

$$\langle P_{\mu\nu}, P_{\lambda\rho}\rangle = \frac{\mathrm{i}}{(k^2-\beta^2)}\omega_{\mu\nu,\lambda\rho} + \frac{\mathrm{i}}{(k^2-\xi^2)}\theta_{\mu\nu,\lambda\rho},$$

where $\xi^2 = \beta^2 + 2m^2$. In order to guarantee unitarity, we must have the real part of the current–current amplitude greater than zero. We can observe from the propagators (41) that the longitudinal sectors of the fields W and P exhibit no tachyons and they are not dynamical. On the other hand, the transverse degrees of freedom of the non-mixing terms are physical as far as we assume in the expression (42) that $\bar{\mu}_+^2 > \bar{\mu}_-^2$, $\bar{\mu}_+^2 > \bar{\beta}^2$ and $\bar{\alpha}^2 > \bar{\mu}_-^2$. As a consequence, the $\mathcal{L}(W)$ has a pole $k^2 = \bar{\mu}_+^2$ for the propagator $\langle W_\mu, W_\nu\rangle^{\mathrm{T}}$, and a pole $k^2 = \bar{\mu}_-^2$ for the propagator $\langle W_{\mu\nu}, W_{\lambda\rho}\rangle^{\mathrm{T}}$; these are dynamical physical excitations. As already mentioned at the end of Sect. 3, we should state more clearly our argument concerning the consequences of the quartic B_μ-coupling which yields non-unitarity due to the coupling of its longitudinal component. Indeed, let us present below the expression giving the power-counting, δ, for the primitively divergent graphs:

$$\delta = 4 - E_H - E_B - E_A - V_{BHA} - I_{BH}, \tag{44}$$

where E_H, E_B and E_A stand for the number of external lines of the H, B and A fields respectively, V_{BHA} expresses the number of super-renormalizable topological vertices and I_{BH} counts the number of mixed matter propagators, whose asymptotic behavior goes as $\frac{1}{k^3}$. The self-energy contribution for the B-field has a logarithmic divergence associated to the kinetic part (the external momentum, p, comes from the momentum-space loop integral, lowering the divergence from quadratic to a logarithmic one). Its mass correction is, however, quadratically divergent. Based on this power-counting argument and on the Ward identities derived for the model, we can safely state that the loop corrections to the kinetic term $|\partial_\mu B^\mu|^2$ are suppressed by a logarithmic factor of the cut-off. Indeed, in view of the massive propagators inside the loops and due to the dimensional regularization technique, there appear logarithms involving the masses. Therefore, if one does not introduce that kinetic term at tree level, and if one considers the model much below its energy cut-off, non-physical dynamical modes are suppressed. So, if working in this safe region of the model, the B_μ-field accomplishes its task of breaking the Lorentz symmetry without the price of violating unitarity. Taking the viewpoint of considering the model below its cut-off, we understand that the low-energy ghost modes do not show up through radiative corrections and renormalizability is not jeopardized as long as we keep far below the energy cut-off. The model should then be viewed as an effective theory valid up to a cut-off below which Lorentz symmetry is spontaneously broken.

7 Conclusions

In this paper, we have presented a charged vector-tensor (CSKR) matter field model which shows a connection between a vector (1-form) field B_μ and an antisymmetric tensor (2-form) field $H_{\mu\nu}$ via a topological interaction in (1+3)D. Furthermore, it presents a global $U(1)$ symmetry. We have shown that the introduction of a self-interacting B^4-type potential in the Lagrangian (1) is necessary to define the topological 2-form (tensor) current (5). Thus, we propose a vector-tensor current which is a doublet where it accommodates the ordinary Noether current J_μ along with the topological 2-form current $J_{\mu\nu}$. With this procedure, we can obtain the physical spectrum of the model in a direct way, where we verify the existence of two distinct simple physical (non-tachyonic and non-ghost) poles with masses μ_+ and μ_- for the transverse sectors. We emphasize that, in spite of the peculiar form of the model, the longitudinal ones decouples. The propagator poles include topological (m^2) and Proca (α^2 and β^2) mass parameters which implies that the very same physical degrees of freedom are present in the propagators of the 1-form B_μ (KR) field and the 2-form $H_{\mu\nu}$ (KB) field. In the charged case, we introduce a gauge interacting field A_μ and we explore the low-energy dynamics of the model, by studying the spontaneous symmetry breaking (SSB) mechanism. The quartic form of the potential energy of the vector-tensor matter field indicates that the SSB mechanism could take place for this field. Only the vector (1-form) B_μ (KR) field can reach the very minimum of energy of the model, consequently it is responsible for the SSB mechanism. On the other hand, the vacuum energy value of this field naturally violates the Lorentz symmetry which fixed a vector deviation to the minimum and implies contributions to the splitting of the mass spectrum. We observe that the possibilities of the fixed vector has to be of physical consistency preventing ghost and tachyon degrees of freedom, and so following in an analogous way the classification obtained in [21]. The effects of Lorentz violating terms have recently been subject of study due to effects in astrophysical and condensed matter physics [9,22]. We have studied the contributions of the fixed vector b_μ to the mass at tree level, where to this purpose we find it suitable to introduce the unitary gauge, and we redefine the fields and parameters in this effective resulting model. We remark that the field Y_μ is gauged away and we consequently recover the usual gauge transformation of the dynamical elements of the model. We obtain two independent dynamical systems: one describes a model where a KB-field interacts with a gauge field, and another one a neutral KB-field interacting with the "Higgs" vector field. We emphasize that this model presents a new neutral vector particle without the introduction of an "extra" $U(1)$ group as it has been largely explored in the literature. Indeed, many phenomenological works have been proposed with the aim of suggesting an extra symmetry to account for discrepancies in the standard model, particularly Z–Z' mixing [30]. In astrophysical models coming from high energy considerations, it has been suggested that non-baryonic dark matter could have an exotic astrophysical origin where a possible mirror matter described by extra symmetries [31, 32] emerges as an interesting possibility. In our case, an extra and/or mirror matter has a topological origin as it was suggested in [31], and the Lorentz violation could imply optical astrophysical effects that influence the red-shift deviations and the anisotropy of the cosmic radiation background, and that possibly mask the observational effects of topological defects. Finally, we analyze the spectrum of the two independent dynamical models obtained after the SSB mechanism. In the gauge-KR sector, we obtain the condition on the mass parameters and the fixed vector so as to have a physically consistent condition where the mass term to the 2-form matter field is crucial. We obtain that the longitudinal part has no dynamics. The result is that this sector could represent the effective dynamics of a charged spin-1 particle. On the other hand, the Higgs–KR sector can represent the dynamics of a massive neutral particle that, due to the Lorentz violation, can only be analyzed in the longitudinal and transverse projected degrees of freedom. Indeed, from another perspective this particular feature of Lorentz violation is emphasized in its classical electromagnetic version [33] and in perturbation studies of the model [34]. We obtain the conditions on the mass parameters (topological and not) to get physically consistent degrees of freedom at tree level. In this sector, the fixed vector b_μ dictates a preferred direction in space. We can conclude to the given perspective to compute a dimensional reduction of the KR model from $(1+3)$D to $(1+2)$D, where we can conjecture the existence of charged "vector solitons" derived off topological solutions in $(1+2)$D, which is the study of a forthcoming work [20].

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References

1. A. Linde, Rep. Progr. Phys. **42**, 389 (1979)
2. S. Deser, R. Jackiw, S. Templenton, Phys. Rev. Lett. **48**, 975 (1982)
3. F. Wilczek, Phys. Rev. Lett. **48**, 1144 (1982)
4. J. Schonfeld, Nucl. Phys. B **185**, 157 (1985)
5. O.M. Del Cima, F.A.B. Rabelo de Carvalho, Int. J. Mod. Phys. A **10**, 1641 (1995)
6. E. Cremmer, J. Scherk, Nucl. Phys. B **72**, 117 (1974)
7. M. Kalb, P. Ramond, Phys. Rev. D **9**, 227 (1974)
8. N. Dragon, S.M. Kuzenko, Phys. Lett. B **420**, 64 (1998); I. Buchbinder, A. Hindawi, B.A. Ovrut, Phys. Lett. B **413**, 79 (1997); P. Claus, B. de Wit, M. Faux, P. Termonia, Nucl. Phys. B **491**, 201 (1997); P. Claus, B. de Wit, M. Faux, B. Kleijn, R. Siebelink, P. Termonia, Phys. Lett. B **373**, 81 (1996)
9. G. Amelino-Camelia, The three perspectives on the quantum-gravity problem and their implications for the fate of Lorentz symmetry, gr-qc/0309054; O. Bertolami, R. Lehnert, R. Potting, A. Ribeiro, Phys. Rev. D **69**, 083513 (2004) ; O. Bertolami, Threshold Effects and Lorentz Symmetry, hep-ph/0301191; R. Lehnert, Phys. Rev. D **68**, 085003 (2003); V. A. Kostelecký, R. Lehnert, M. J. Perry, Phys. Rev. D **68**, 123511 (2003); E.C. Vagenas, JHEP **0307**, 046 (2003); J.W. Moffat, Int. J. Mod. Phys. D **12**, 1279 (2003); S. Das, G. Kunstatter, Class. Quant. Grav. **20**, 2015 (2003); J.D. Barrow, Phys. Lett. B **564**, 1 (2003); P.C.W. Davies, T.M. Davis, C.H. Lineweaver, Nature **418**, 602 (2002); V. A. Kostelecký, R. Lehnert, Phys. Rev. D **63**, 065008 (2001); K. Greisen, Phys. Rev. Lett. **16**, 748 (1966); G.T. Zatsepin, V.A. Kuzmin, Pisma Zh. Eksp. Teor. Fiz. **4**, 114 (1966)
10. V.A. Kostelecký, R. Potting, Nucl. Phys. B **359**, 545 (1991); V.A. Kostelecky, S. Samuel, Phys. Rev. D **39**, 683 (1989)
11. D. Colladay, V.A. Kostelecký, Phys. Rev. D **58**, 116002 (1998)
12. D. Colladay, V.A. Kostelecký, Phys. Rev. D **55**, 6760 (1997)
13. A. Coleman, S.L. Glashow, Phys. Lett. B **405**, 249 (1997)
14. A. Coleman, S.L. Glashow, Phys. Rev. D **59**, 116008 (1999)
15. A.A. Andrianov, R. Soldati, Phys. Rev. D **51**, 5961 (1995)
16. A.A. Andrianov, R. Soldati, Phys. Lett. B **435**, 449 (1998)
17. S. Carrol, G. Field, R. Jackiw, Phys. Rev. D **41**, 1231 (1999)
18. R. Jackiw, V.A. Kostelecký, Phys. Rev. Lett. **82**, 3572 (1999)
19. D. Hutsemékers, H. Lamy, Astron. Astrophys. **367**, 381 (2001); **332**, 410 (1998); M. Goldhaber, V. Trimble, J. Astrophys. Astron. **17**, 17 (1996)

20. L.P. Colatto, A.L.A. Penna, W.C. Santos, in preparation
21. H. Belich, Jr., M.M. Ferreira, Jr., J.A. Helaÿel-Neto, M.T.D. Orlando, Phys. Rev. D **67**, 125011 (2003); D **68**, 025005 (2003); A.P. Baeta Scarpelli, H. Belich, J.L. Boldo, L.P. Colatto, J.A. Helayël-Neto, A.L.M.A. Nogueira, Nucl. Phys. Suppl. **27**, 105 (2004); H. Belich, J.L. Boldo, L.P. Colatto, J.A. Helayël-Neto, A.L.M.A. Nogueira, Phys. Rev. D **68**, 065030 (2003)
22. H. Muller, C. Braxmaier, S. Herrmann, A. Peters, C. Laemmerzahl, Phys. Rev. D **67**, 056006 (2003)
23. C.N. Yang, R.L. Mills, Phys. Rev. **96**, 191 (1954)
24. P.W. Higgs, Phys. Lett. **12**, 132 (1964)
25. P.W. Higgs, Phys. Rev. Lett. **13**, 508 (1964)
26. P.W. Higgs, Phys. Rev. **145**, 1156 (1966)
27. R. Jackiw, K. Johnson, Phys. Rev. D **8**, 2386 (1973)
28. J.M. Cornwall, R.E. Norton, Phys. Rev. D **8**, 3338 (1973)
29. S. Coleman, E. Weinberg, Phys. Rev. D **7**, 1688 (1973)
30. S. Baek, N.G. Deshpande, X.G. He, P. Ko, Phys. Rev. D **64**, 055006 (2001); R. Foot, X.G. He, H. Lew, R.R. Volkas, Phys. Rev. D **50**, 4571 (1994); R. Foot, X.G. He, Phys. Lett. B **267**, 509 (1991)
31. R. Foot, Experimental implications of mirror matter-type dark matter, astro-ph/0309330
32. R. Foot, Phys. Rev. D **69**, 036001 (2004); A.Y. Ignatiev, R.R. Volkas, Mirror matter, hep-ph/0306120; R. Foot, S. Mitra, Astropart. Phys. 19, 739 (2003); R. Foot, R.R. Volkas, Phys. Rev. D **68**, 021304 (2003); A.Y. Ignatiev, R.R. Volkas, Phys. Rev. D **68**, 023518 (2003); G. Sanchez-Colon, R. Huerta, A. Garcia, J. Phys. G **29**, L15 (2003); R.N. Mohapatra, S. Nussinov, V.L. Teplitz, Phys. Rev. D **66**, 063002 (2002); Z. Berezhiani, D. Comelli, F.L. Villante, Phys. Lett. B **503**, 362 (2001); N.F. Bell, Phys. Lett. B **479**, 257 (2000)
33. L.P. Colatto, J.A. Helayël-Neto, M. Hott, W.A. Moura-Melo, in preparation
34. H. Belich, J.L. Boldo, J.A. Helayël-Neto, A.L.M.A. Nogueira, to appear