Minimal coupling of the Kalb-Ramond field to a scalar field

E. Di Grezia and S. Esposito

Dipartimento di Scienze Fisiche, Università di Napoli “Federico II”
and
Istituto Nazionale di Fisica Nucleare, Sezione di Napoli
Complesso Universitario di Monte S. Angelo, Via Cinthia, I-80126 Napoli, Italy

We study the direct interaction of an antisymmetric Kalb-Ramond field with a scalar particle derived from a gauge principle. The method outlined in this paper to define a covariant derivative is applied to a simple model leading to a linear coupling between the fields. Although no conserved Noether charge exists, a conserved topological current comes out from the antisymmetry properties of the Kalb-Ramond field. Some interesting features of this current are pointed out.

I. INTRODUCTION

Space-time non-commutativity is one of the key new ideas which follows from recent developments in string and matrix theory [1]. Non-commutativity implies general covariance and, therefore, it seems likely that a non-commutative Yang-Mills theory is a good candidate for a unified and potentially renormalizable theory of the fundamental interactions including gravity.

Whereas the structure of the space-time becomes non-commutative, we can describe it, in analogy to quantum phase space, in terms of the algebra generated by non-commuting coordinates:

$$[x^\mu, x^\nu] = i\theta^{\mu\nu}(x),$$

with $\theta^{\mu\nu}$ an antisymmetric tensor [2].

Antisymmetric tensor fields are widely used in string models [3], [4], [5] as well as in some supersymmetric theories. For example, they appear naturally in $N = 2$ extended supersymmetry as auxiliary fields (being the highest-dimension fields in a supermultiplet) and in the 11-dimensional formulation of $N = 8$ extended supergravity in which they are dynamical fields. Furthermore, in quantum gravity, antisymmetric tensor fields appear as Lorentz ghost fields [6]. It is, therefore, quite important to study the dynamics of such fields and, especially, their coupling to matter (scalar or fermion particles). The best studied example of an antisymmetric tensor field is the electromagnetic field strength $F_{\mu\nu}$, whose dynamics is very well known (see, for example, [7]). The coupling of the electromagnetic field to matter fields proceeds, usually, through the gauge principle with the aid of the vector potential $A_\mu$ (a massless rank-1 field): interaction is introduced in the theory by requiring local gauge invariance for the matter fields. In the electromagnetic case, the gauge transformations for the vector potential are $A_\mu(x) \rightarrow A_\mu(x) + \partial_\mu \alpha(x)$, where $\alpha(x)$ is an arbitrary $x$-dependent function, and the interaction with a charged scalar field $\varphi$ (with charge $e$) is obtained by replacing the usual derivative $\partial_\mu$ with the covariant derivative $D_\mu$ in the free field Lagrangian:

$$\varphi(x) \rightarrow e^{-ie\alpha(x)}\varphi(x).$$

As a consequence of gauge invariance, by virtue of the Noether theorem, electric charge conservation is obtained. Note, however, that such a minimal coupling prescription is not the only possible one; for example, magnetic moment

*Electronic addresses: digrezia@na.infn.it, salvatore.esposito@na.infn.it
interaction is described by a Lagrangian term involving directly the physical field strength $F_{\mu\nu}$, rather than the gauge-dependent vector potential.

Several papers have appeared [8] in which the generalization of the gauge principle to (abelian) rank-2 antisymmetric fields (Kalb-Ramond fields) is studied. However, the interaction of a scalar or fermion particle with a Kalb-Ramond field usually proceeds through the coupling with the Maxwell field [9]. Indeed, the matter particles interact with the electromagnetic field coupled to a Kalb-Ramond field, so that only an indirect interaction is allowed.

While the coupling between such fields and matter fields can always be introduced by adding an \textit{ad hoc} term in the complete Lagrangian without invoking a gauge principle, the main problem with an interaction generated by an (abelian) gauge group transformation comes from the difficulty to construct a covariant derivative from rank-2 gauge fields in analogy with the electromagnetic case. The present work is aimed to further study such a problem by considering, for simplicity, massless fields.

In view of some applications considered below in the paper, in the next section we briefly review the dynamics of an antisymmetric tensor field and point out its substantial equivalence (in the massless case) with that of a (real) scalar field. Instead in Section III we give a procedure to couple a scalar particle with a Kalb-Ramond field through a gauge principle with a linear coupling; some comments on quadratic coupling are also reported. Finally, in Section IV, we study the dynamics of the interacting fields, pointing out some peculiar features, while in the last section we give our conclusions and outlook.

II. DYNAMICS OF A FREE KALB-RAMOND FIELD

Let us consider a Maxwell-like gauge theory where the role of the vector potential $A_{\mu}$ is played by an antisymmetric tensor field $\theta_{\mu\nu}$. Its dynamics is described by the following Lagrangian:

$$L_{\theta} = -\frac{1}{12} H_{\lambda\mu\nu} H^{\lambda\mu\nu},$$

(4)

where the field strength $H_{\lambda\mu\nu}$ is defined by:

$$H_{\lambda\mu\nu} = \partial_\lambda \theta_{\mu\nu} + \partial_\mu \theta_{\nu\lambda} + \partial_\nu \theta_{\lambda\mu}.$$

(5)

The equations of motion for the free $\theta_{\mu\nu}$ field are, then, similar to the Maxwell equations and read as follows:

$$\partial_\mu H^{\mu\nu\sigma} = 0.$$

(6)

The Lagrangian in Eq. (4) is invariant under the gauge transformations:

$$\theta_{\mu\nu} \rightarrow \theta'_{\mu\nu} = \theta_{\mu\nu} + \partial_\mu \Lambda_\nu - \partial_\nu \Lambda_\mu,$$

(7)

where $\Lambda_\nu$ is an arbitrary vector field. This gauge freedom can be used to simplify the writing of the field equations. Indeed, for example, by using the "generalized" Lorentz condition $\partial^\mu \theta_{\mu\nu} = 0$ we find that $\theta_{\mu\nu}$ satisfies the ordinary wave equation:

$$\Box \theta_{\mu\nu} = 0.$$

(8)

It is easy to prove that the degrees of freedom of the antisymmetric tensor field $\theta_{\mu\nu}$ are just the same of a scalar field $\theta$. Let us consider the dual vector field (in the sense of Poincaré) $\theta_\mu$ defined by:

$$\theta_\mu = \frac{1}{6} \epsilon_{\mu
u\rho\sigma} H^{\nu\rho\sigma} = \frac{1}{2} \epsilon_{\mu
u\rho\sigma} \partial^\nu \theta^{\rho\sigma},$$

(9)

which satisfies the Bianchi-type identity:

$$\partial^\mu \theta_\mu = 0.$$

(10)

The equations of motion in (6) now become:

$$\partial_\mu \theta_\nu - \partial_\nu \theta_\mu = 0,$$

(11)

pointing out that $\theta_\mu$ has to be the gradient of a scalar field $\theta$: 

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\[ \theta_\mu = \partial_\mu \theta . \] (12)

With these changes of variable, the Lagrangian describing the system can be cast in a form similar to that for a scalar field:

\[ \mathcal{L}_\theta = \frac{1}{2} \partial_\mu \theta \partial^\mu \theta , \] (13)

this showing that the dynamics of a massless antisymmetric tensor field \( \theta_{\mu\nu} \) is completely equivalent (on the classical level\(^1\)) to that of a massless real scalar field \( \theta \). Note, however, that Eq. (12) is a direct consequence of the free field equations of motion (11), so that the mentioned equivalence holds true only in the case of non-interacting fields.

**III. GAUGE PRINCIPLE WITH A LINEAR COUPLING**

Let us consider a charged scalar field \( \phi \) described by the usual Lagrangian:

\[ \mathcal{L}_\phi = (\partial_\mu \phi) \dagger (\partial^\mu \phi) + V(\phi), \] (14)

where \( V(\phi) \) is a given scalar potential including, eventually, a mass term. In this section we study direct interaction of this field with a Kalb-Ramond field \( \theta_{\mu\nu} \) generated by possible gauge group transformations leading to linear coupling with the \( \theta_{\mu\nu} \) field.

In analogy with the electromagnetic case, we introduce a covariant derivative \( D_\mu \) as follows:

\[ iD_\mu = i\partial_\mu + kX_\mu , \] (15)

where \( X_\mu \) is a rank-1 tensor to be defined in terms of the \( \theta_{\mu\nu} \) field and \( k \) is a suitable coupling constant.

The Lagrangian describing the scalar field \( \phi \) interacting with \( \theta_{\mu\nu} \) is, therefore:

\[ \mathcal{L}_{\text{int}} = (D_\mu \phi) \dagger (D^\mu \phi) + V(\phi) . \] (16)

Obviously the Lagrangian in Eq. (16) must be invariant under the gauge transformation in Eq. (7), so that the transformation properties of \( X_\mu \) (and the corresponding ones for the field \( \phi \)) are crucial in the identification of \( X_\mu \) itself.

Note, however, that the field \( X_\mu \) cannot be written as a gradient of a given function \( \alpha(x) \), in order to have a "genuine" interaction Lagrangian in Eq. (16). Indeed, let us assume that:

\[ X_\mu = \partial_\mu \alpha , \] (17)

and consider a local phase transformation for the scalar field \( \phi \):

\[ \phi \rightarrow \Phi = \phi e^{ik\alpha} . \] (18)

Substitution into Eq. (16) immediately leads to:

\[ \mathcal{L}_{\text{int}} = (D_\mu \phi) \dagger (D^\mu \phi) + V(\phi) = (\partial^\mu \Phi) \dagger \partial_\mu \Phi) + V(\phi) , \] (19)

and since \( \Phi \) and \( \phi \) represent the same physical object, we conclude that the interaction introduced by means of the \( X_\mu \) field in Eq. (17) is fictitious, because it can be re-absorbed by a phase redefinition of the scalar field.

Starting from the rank-2 tensor \( \theta_{\mu\nu} \), the simplest linear choice for \( X_\mu \) is the following:

\[ X_\mu = \partial^\nu \theta_{\mu\nu} . \] (20)

Such a definition for \( X_\mu \) is, however, useful only when the "generalized" Lorentz condition is not fulfilled, since in this case \( X_\mu \) is identically zero and no physical interaction appears. Disregarding this gauge choice, the gauge group for the scalar field \( \phi \) is easily obtained. Indeed, by applying the gauge transformation in Eq. (7) to the Lagrangian in (16), we find that it remains unchanged if the scalar field \( \phi \) transforms as follows:

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\(^1\)For the quantum equivalence see, for example, Ref. [10]
\[ \phi \rightarrow \phi' = \phi e^{ik\eta}, \]

with \( \eta = \partial^\nu \Lambda_\nu \). Note that the gauge freedom for the \( \theta_{\mu\nu} \) field is not altered if in Eq. (7) we choose a divergence-less gauge function \( \Lambda_\nu \). In this case \( \eta \) is zero and Eq. (21) reduces to the unit transformation, so that the gauge group underlying the choice in Eq. (20), when it is applicable, acts as the identity on the scalar field \( \phi \). As a consequence, by means of the Noether theorem, no conserved charge comes out.

The only alternative for \( X_\mu \), which is linear in the \( \theta_{\mu\nu} \) field, is the following

\[ X_\mu = \theta_\mu = \frac{1}{2} \epsilon_{\mu\rho\sigma} \partial^\rho \theta^{\rho\sigma}. \]

Now we have no limitation on the gauge condition to be used and is a simple task to show that the gauge group associated to the scalar field is, again, the identity. Indeed, it is easy to convince ourselves that \( X_\mu \) in Eq. (22) is invariant under the gauge transformation in Eq. (7), and thus the Lagrangian in Eq. (16) is automatically invariant when \( \phi \) is unchanged:

\[ \phi \rightarrow \phi' = \phi . \]

We remark that the choice (22) for \( X_\mu \) is not a trivial one, since \( X_\mu \) does not satisfy the constraint (17). In fact, as pointed out in the previous section, the dual field \( \theta_\mu \) can be written as the gradient of \( \theta \) (see Eq. (16)) only for a massless non-interacting Kalb-Ramond field, which is not the present case. It is remarkable that in the limit of no interaction the \( X_\mu \) field becomes unphysical.

A final remark on the physical dimensions of the coupling constant is in order. We immediately see that, for the cases considered above, \( k \) has the dimensions of the inverse of a mass, so that the corresponding theory is non-renormalizable. We point out that such a property is a peculiar feature of Quantum Gravity, which is a natural framework, however, for the theory developed here.

### A. Non-linear coupling

For completeness, we briefly discuss some particular direct interactions between a scalar field and a Kalb-Ramond field involving non-linear terms (which are quadratic in the Kalb-Ramond field).

First of all we note that \( n \)-linear (\( n > 1 \)) terms in \( \theta_{\mu\nu} \) appearing in the covariant derivative have to be gauge-invariant in order to assure the invariance of the Lagrangian (neglecting terms which are total divergences). Indeed, when considering the \( X_\mu \) field where, for example, non gauge-invariant bilinear terms in \( \theta_{\mu\nu} \) appears, we recognize that, under a gauge transformation (7), two derivatives of two gauge (4-vectors) functions \( \lambda_\mu \) come out. In general, the corresponding gauge terms in \((D^\mu \phi)^\dagger (D_\mu \phi)\) cannot be absorbed by a phase transformation (containing only one scalar function) of the scalar field \( \phi \). Therefore, simple choices for \( X_\mu \) involve only the gauge-invariant field \( H_{\mu\nu\rho} \) and its dual \( \theta_\mu \) and, as a consequence, the gauge group is the identity.

Some example of quadratic interactions are as follows:

\[ X_\mu = \theta_\mu \partial^\rho \theta_\mu, \]
\[ X_\mu = \epsilon^{\alpha\beta\gamma\delta} \partial^\rho H_{\mu\alpha\beta} H_{\gamma\delta\eta}, \]
\[ X_\mu = \epsilon_{\mu\rho\sigma\gamma} \partial^\rho H^{\alpha\beta\gamma} \theta_\gamma. \]

Note that in the case considered the Eq. (25) (and for similar terms), \( X_\mu \) vanishes only when \( H_{\gamma\delta\eta} \) satisfies its equation of motion in absence of interaction Eq. (6). Finally, we point out that allowing quadratic terms for \( X_\mu \) as those in Eqs. (24)-(26) (and similar ones), the coupling constant \( k \) in the covariant derivative must have the dimensions of the inverse of the 4th power of a mass, so that the non-renormalizability of the theory is greatly worsened with respect to the linear coupling case.

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\(^2\)More in general, we can allow a phase transformation \( \phi \rightarrow \phi' = \phi e^{i\beta} \), with a constant \( \beta \).
IV. INTERACTING FIELD DYNAMICS.

The dynamics of a charged scalar field interacting directly with a Kalb-Ramond field is described by the following Lagrangian:

\[
L = (D_\mu \phi)^\dagger (D^\mu \phi) - m^2 \phi^\dagger \phi - \frac{1}{12} H_{\mu\nu\rho} H^{\mu\nu\rho},
\]  

(27)

where the covariant derivative is written as in Eq. (15) and, for simplicity, we have neglected a scalar potential term. By expliciting the \(X_\mu\) term, we can rewrite Eq. (27) as:

\[
L = (\partial_\mu \phi)^\dagger (\partial^\mu \phi) - (m^2 - k^2 X^2) \phi^\dagger \phi - \frac{1}{12} H_{\mu\nu\rho} H^{\mu\nu\rho},
\]

(28)

where the second term in the sum accounts for the effective mass acquired by the scalar particle interacting with the Kalb-Ramond field, while the following term describes this interaction.

The Euler-Lagrange equation for \(\phi\) immediately follows:

\[
\partial_\mu (\partial^\mu \phi) + (m^2 - k^2 X^2) \phi = 2i k \phi \partial_\mu \phi + \sqrt{2} k |\phi|^2 X_{\mu
u},
\]

(29)

Note that, in the linear coupling case, the last term vanishes, since \(\partial_\mu X_{\mu\nu} = 0\).

Instead, the equations of motion for the interacting Kalb-Ramond field are:

\[
\partial_\sigma H_{\sigma\mu\nu} = J_{\mu\nu},
\]

(30)

with

\[
J_{\mu\nu} = \partial^\sigma H_{\sigma\mu\nu} = 2 \left[ -ik (\phi \partial^\mu \phi - \phi^\dagger \partial_\mu \phi) + 2k^2 |\phi|^2 X_{\sigma} \frac{\partial X_{\sigma}}{\partial (\partial_\mu \theta_{\mu\nu})} \right] - \phi \partial_\sigma (\phi^\dagger \partial_\rho \phi + 2k^2 |\phi|^2 X_{\sigma} \frac{\partial X_{\sigma}}{\partial (\partial_\rho \theta_{\mu\nu})}).
\]

(31)

Given the symmetry properties of the \(\theta_{\mu\nu}\) field, it follows that the current \(J_{\mu\nu}\) is antisymmetric while exchanging the indices \(\mu\) and \(\nu\):

\[
J_{\mu\nu} = -J_{\nu\mu}.
\]

(32)

Moreover, from the field equations (30) we deduce that, despite of the explicit form of \(X_\mu\), \(J_{\mu\nu}\) is a conserved current:

\[
\partial^\mu J_{\mu\nu} = 0.
\]

(33)

Note that such a property follows, again, from the antisymmetry of \(H_{\sigma\mu\nu}\).

Let us now consider the explicit interesting case in which \(X_\mu\) is given by Eq. (22). After some algebra we obtain the following expression for the current:

\[
J_{\mu\nu} = \partial^\sigma T_{\sigma\mu\nu}.
\]

(34)

with

\[
T_{\sigma\mu\nu} = -2 (ik \epsilon_{\sigma\rho\nu} \phi^\dagger \partial^\rho \phi + k^2 |\phi|^2 H_{\sigma\mu\nu}).
\]

(35)

It is remarkable that, in the case considered, the current \(J_{\mu\nu}\) is the gradient of an antisymmetric rank-3 tensor \(T_{\sigma\mu\nu}\). Indeed, as a consequence of this, we have that Eq. (33) holds independently of the equations of motion (by taking the divergence of Eq. (34) we obtain a vanishing R.H.S due to the symmetry properties of \(T_{\sigma\mu\nu}\)), so that \(J_{\mu\nu}\) is a conserved topological current. Moreover on substituting Eq. (34) into Eq. (30) we find that

\[
\partial^\mu \tilde{H}_{\mu\rho\sigma} = 0.
\]

(36)

with

\[
\tilde{H}_{\mu\rho\sigma} = H_{\mu\rho\sigma} - T_{\mu\rho\sigma}.
\]

(37)

i.e. the novel field \(\tilde{H}_{\mu\rho\sigma}\) can be interpreted as a “dressed” Kalb-Ramond field satisfying the free field equation (36).
A. The symmetry group of the topological current

Let us now turn back to the topological current $J_{\mu\nu}$ in Eqs (34) and (35) and write it as sum of a term $J_{\mu\nu}^0$, which does not depend explicitly on the Kalb-Ramond field, and a term $J_{\mu\nu}^1$ which vanishes for vanishing $H_{\alpha\mu\nu}$:

$$J_{\mu\nu} = -2k(J_{\mu\nu}^0 + kJ_{\mu\nu}^1)$$  \hspace{1cm} (38)

$$J_{\mu\nu}^0 = \epsilon_{\alpha\beta\mu
u}(\partial^\alpha \phi^\dagger)^{\beta}$$  \hspace{1cm} (39)

$$J_{\mu\nu}^1 = \partial^\alpha (|\phi|^2 H_{\alpha\mu\nu})$$  \hspace{1cm} (40)

The "free" current $J_{\mu\nu}^0$ remains invariant if we perform a rotation with an imaginary angle $i\alpha$ of the fields $\phi, \phi^*$.:

$$\begin{pmatrix} \phi \\ \phi^* \end{pmatrix} \rightarrow \begin{pmatrix} \phi' \\ \phi'^* \end{pmatrix} = \begin{pmatrix} \cos i\alpha & \sin i\alpha \\ -\sin i\alpha & \cos i\alpha \end{pmatrix} \begin{pmatrix} \phi \\ \phi^* \end{pmatrix}$$  \hspace{1cm} (42)

However the term $J_{\mu\nu}^1$, and hence the whole current $J_{\mu\nu}$, is not invariant under the group defined above, so that the corresponding topological symmetry can be only viewed as an approximate invariance of the conserved current in the limit of a small coupling constant (the $J_{\mu\nu}^1$-term is quadratic in $k$). The general transformation for the scalar field $\phi$ leaving invariant the whole current $J_{\mu\nu}$ is the following:

$$\phi \rightarrow \phi' = \phi e^{i(\alpha|\phi|^2 + \beta)}$$  \hspace{1cm} (43)

with $\alpha, \beta$ two real parameters, and, in the limit $\alpha = 0$, we recover the simple global phase transformation. However, since Eq. (43) is highly non-linear for $\alpha \neq 0$, it is not evident the physical meaning corresponding to the transformation in Eq. (43).

V. CONCLUSIONS

In this paper we have introduced a direct interaction between a scalar particle and a Kalb-Ramond field by defining a covariant derivative $D_\mu$, where an appropriate auxiliary vector field $X_\mu$ (depending on the Kalb-Ramond field) appears. Several possible choices for $X_\mu$ have been studied, leading to linear or quadratic coupling with the scalar field. In the simple, viable models considered here, the gauge group underlying the theory is the identity, so that no conserved Noether charge exists. However, due to the antisymmetry properties of the Kalb-Ramond field, a conserved (antisymmetric) topological current arises in the simplest model, which appears in the equations of motion for $H_{\mu\nu\rho}$. Since this current is a divergence of a rank-3 antisymmetric tensor, it is possible to define a "dressed" Kalb-Ramond field strength, obeying the free field equations, which describes the dynamics of the interacting field. Some possible applications of our results are concerned with the theory of gravity, where a zero mass Kalb-Ramond field is the source of torsion in Einstein-Cartan theory [13]. Moreover, in recent years, it has been pointed out that the presence of Kalb-Ramond fields in the background space-time leads to several interesting astrophysical and cosmological phenomena like cosmic optical activity and neutrino helicity flip [11]. This motivates the study of some important problems related to the standard Friedman-Robertson-Walker cosmological model in light of Kalb-Ramond cosmology [12], [2] in an inflationary framework, where the coupling to a scalar field is crucial.

3 More in general, the current $J_{\mu\nu}^0$ is left unchanged if the fields $\phi, \phi^*$ undergo the following linear transformation:

$$\begin{pmatrix} \phi \\ \phi^* \end{pmatrix} \rightarrow \begin{pmatrix} \phi' \\ \phi'^* \end{pmatrix} = \begin{pmatrix} a & b \\ b^* & a^* \end{pmatrix} \begin{pmatrix} \phi \\ \phi^* \end{pmatrix} + \begin{pmatrix} c \\ c^* \end{pmatrix}$$  \hspace{1cm} (41)

with $a, b, c$ three arbitrary complex numbers satisfying the constraint $|a|^2 - |b|^2 = 1$. Note that, since the determinant of the transformation matrix is equal to 1, the group of transformations defined in (41) is that of equiaffinities in the complex plane, which is a subgroup of $SL(2, C)$. 

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ACKNOWLEDGMENTS

We are indebted with Dr. G. Fiore for his kind cooperation and many interesting suggestions. Useful discussions with Dr.s R. Marotta, O. Pisanti and Prof. G. Miele have been appreciated as well.

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