ISSN 0029-3865

CBPF-NF-052/84

TOWARDS AN ALTERNATIVE UNIFICATION OF MASSLESS AND MASSIVE VECTOR BOSONS

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Abstract

A possible extension of the gauge principle is presented where two distinct gauge potentials are introduced in association with a single U(1) gauge group, each of them being taken to interact with a different kind of matter field. In such a picture, a massive vector boson naturally shows up in the physical spectrum. A massive photon without Higgs can be introduced. Renormalizability is seen to be a feature of the model. Possible supersymmetrizations are also contemplated.

Key-words: Two distinct gauge potentials in a single group.

1. Introduction

We could emphasize the relevance of gauge theories in connection with three aspects. First, the primary fact that the requirement of gauge invariance dictates the interactions governing the basic dynamics of a prescribed set of elementary particles [1]. Second, the emergence of the self-interactions among the mediating gauge bosons which yield the remarkable property of asymptotic freedom [2] and the appearance of magnetic monopoles as regular solutions of the classical field equations [3]. Moreover, the good agreement between the quantum-corrected mass formulas for the weak gauge bosons predicted by the Weinberg-Salam-Glashow model [4] and the recent results of the UA1 and UA2 collaborations [5] is another remarkable success of the gauge principle. Finally, the differential-geometric interpretation of gauge theories in terms of fiber bundles with connections suggests the possibility of a geometrical picture for the fundamental forces of Nature [6].

These facts, among many others we did not quote above, encourage us to follow the gauge method in our attempts of giving a representation for the interactions among some groups of elementary particles thought of as the building-blocks of the known matter. By then adopting the gauge principle, one could pursue an enlargement of the concept of gauge by associating *different* families of *gauge potentials* to different types of building-blocks fulfilling irreducible representations of a *single gauge group* [7]. Take, for instance, two gauge potentials, A_μ and B_μ, having the following transformation laws:

$$A_\mu \longrightarrow A'_\mu = UA_\mu U^{-1} + \frac{i}{g} (\partial_\mu U) U^{-1} \tag{1.1}$$

and

$$B_\mu \longrightarrow B'_\mu = UB_\mu U^{-1} + \frac{i}{g} (\partial_\mu U) U^{-1} \ , \tag{1.2}$$

which may originate a field strength tensor $G_{\mu\nu}$ given by

$$G_{\mu\nu} \equiv \partial_\mu A_\nu - \partial_\nu B_\mu + g[B_\mu, A_\nu] \ . \tag{1.3}$$

Considering this type of gauge theories, where more than a gauge potential is associated with the same gauge group, we intend to develop a proposal: it is to study the interaction of different families of matter fields by coupling them to different vector bosons. The simplest idea would be to couple the vector fields A_μ and B_μ either to fermionic and bosonic families respectively or to fermionic fields describing different families of spin-$\frac{1}{2}$ particles. We shall here adopt the first viewpoint and try to discuss the possibility of setting a sensible (renormalizable and unitary) field theory based on the association of two different gauge potentials to the same gauge group.

The central problem to be faced is the following: the presence of more than one gauge potential in just one gauge group implies that undesirable spin-zero modes, intrinsically contained in representation $(\frac{1}{2},\frac{1}{2})$ of Lorentz group, do not decouple from the physical degrees of freedom. They can plague the theory with negative-metric ghosts, spoiling the renormalizability and the unitarity of the gauge models discussed in [7]. This a thorny problem to be tackled in this class of models.

However, in the U(1) - case that we shall study here, it is easy to bypass this difficulty, as the theory automatically offers the freedom of coupling one of the vector fields to a *globally* conserved matter current. This leads to the decoupling of the spuriou scalar mode that the gauge freedom was not enough to kill off. Our Abelian model could be interpreted as a Q.E.D. for spin-$\frac{1}{2}$ fermions and scalar bosons with the conservation of a global charge interpreted as a fermionic minus bosonic number.

The outline of our paper is as follows: in Section 2, we discuss the issue of coupling the different gauge potentials to matter and establish a $U(1)_{local} \otimes U(1)_{global}$ Lagrangian with the property of being ghost-free. The power - counting renormalizability, the Ward identities and the gauge invariance of the renormalized theory are the content of Section 3. In Section 4, we contemplate possible supersymmetric extensions of the model, motivated by the fact that the original theory naturally starts from fermions and bosons as truly elementary fields transforming equally under the action of the gauge group. Section 5 contains a few conclusive remarks.

2. The Classical Theory

We start from two massless matter fields: a Dirac spinor, ψ, and a complex scalar, ϕ, which undergo the following phase transformations:

$$\psi \longrightarrow \psi' = e^{ig_1\alpha_1}\psi \qquad (2.1)$$

and

$$\phi \longrightarrow \phi' = e^{ig_2\alpha_2}\phi \ , \qquad (2.2)$$

generated by charges whose interpretation we shall discuss later.

Now, we wish to render these transformations local and, via the principle of gauge invariance, to introduce two distinct potentials and analyse the dynamics which emerges from this scheme. By imposing that ψ and ϕ have the local U(1) transformations,

$$\psi \longrightarrow \psi' = e^{ig_1\alpha(x)}\psi \qquad (2.3)$$

and

$$\phi \longrightarrow \phi' = e^{ig_2\alpha(x)}\phi \ , \qquad (2.4)$$

the first step is, as it is usual to do, the definition of the covariant derivatives. In this case, we define two different kinds:

$$D[A]\psi \equiv (\partial + ig_1 A)\psi \qquad (2.5)$$

and

$$D[B]\phi \equiv (\partial + ig_2 B)\phi \ , \qquad (2.6)$$

with

$$A_\mu \longrightarrow A'_\mu = A_\mu - \partial_\mu\alpha \qquad (2.7)$$

and

$$B_\mu \longrightarrow B'_\mu = B_\mu - \partial_\mu \alpha \ . \qquad (2.8)$$

By then defining the field strengths associated to the vector fields A_μ and B_μ,

$$A_{\mu\nu} \equiv \partial_\mu A_\nu - \partial_\nu A_\mu \qquad (2.9)$$

and

$$B_{\mu\nu} \equiv \partial_\mu B_\nu - \partial_\nu B_\mu \ , \qquad (2.10)$$

one can propose the following Lagrangian density which exhibits invariance under the local U(1) transformations given above. It is:

$$\mathcal{L} = -\frac{1}{2} A^{\mu\nu} A_{\mu\nu} - \frac{1}{2} B^{\mu\nu} B_{\mu\nu} + \frac{1}{2} m^2 (A-B)^2 +$$

$$+ \bar{\psi}\, i\gamma . D[A]\psi + (D[B]\phi)^* (D[B]\phi) \ . \qquad (2.11)$$

Some remarks are worthwhile:

(i) The factors $\frac{1}{2}$ in the vector-field kinetic terms have been chosen in order to reproduce the usual term $-\frac{1}{4} F^{\mu\nu} F_{\mu\nu}$ when the Lagrangian will be rewritten in terms of the physical vector fields, as we shall see later.

(ii) A term like $A^{\mu\nu} B_{\mu\nu}$ is also allowed by gauge invariance. However, when re-expressed in terms of the physical fields, one can see that it has no physically significant consequence. So, we do not introduce it in our classical Lagrangian. At the quantum level, we shall come back to this point.

(iii) An interaction term of the form $(A-B)^4$ would be forbidden by the requirement of renormalizability. The reason being that the component (A^0-B^0) does not propagate. Consequently, it has its scale dimension fixed by the mass term, and not by the kinetic term, as it is the case for the components (A^i-B^i). The presence of an interaction like $(A^0-B^0)^4$ would therefore spoil the renormalizability of the theory we are attempting to set.

(iv) The canonical momenta conjugated to the gauge potentials are

$$\begin{cases} \Pi_o[A] = 0 \\ \Pi_i[A] = -A_{oi} \equiv -E_i^A \end{cases} \tag{2.12}$$

and

$$\begin{cases} \Pi_o[B] = 0 \\ \Pi_i[B] = -B_{oi} = -E_i^B \ . \end{cases} \tag{2.13}$$

This first analysis then leads to G degrees of freedom. Now, by using the gauge freedom we have at our disposal, we can reduce this number by 1, so that we can finally state that the gauge potentials carry altogether 5 physical degrees of freedom. Moreover, the Hamiltonian for the free vector fields is non-negative definite,

$$H = \frac{1}{2}(\vec{E}_A^2 + \vec{E}_B^2) + \frac{1}{2}(\vec{B}_A^2 + \vec{B}_B^2) + \frac{1}{2} m^2 [(A^0 - B^0)^2 + (\vec{A} - \vec{B})^2] \geq 0 , \quad (2.14)$$

where

$$\vec{B}_A \equiv \vec{\nabla} \times \vec{A}$$

and

$$\vec{B}_B \equiv \vec{\nabla} \times \vec{B} . \quad (2.15)$$

At this point, our treatment will become more transparent if we re-express the Lagrangian proposed in (2.11) in terms of the physical vector fields, that is, those which diagonalize the mass matrix. They are simply

$$C_\mu \equiv A_\mu - B_\mu \quad (2.16)$$

and

$$D_\mu \equiv A_\mu + B_\mu , \quad (2.17)$$

which are massive (mass m) and massless, respectively. D_μ is the genuine gauge field of the theory,

$$D_\mu \longrightarrow D'_\mu = D_\mu - 2\partial_\mu \alpha , \quad (2.18)$$

and it has 2 physical degrees of freedom, whereas C_μ is gauge invariant and carries 3 physical degrees of freedom (a Proca field). In terms of these physical fields, we clearly see that , despite the presence of the gauge potentials A_μ and B_μ, there is only one true gauge boson in the model.

The Lagrangian now reads:

$$\mathcal{L} = -\frac{1}{4}D^{\mu\nu}D_{\mu\nu} - \frac{1}{4}C^{\mu\nu}C_{\mu\nu} + \frac{1}{2}m^2C^2 +$$

$$+ \bar{\psi}i\gamma.\partial\psi - \frac{1}{2}g_1\bar{\psi}\gamma\psi.D - \frac{1}{2}g_1\bar{\psi}\gamma\psi.C +$$

$$+ \partial\phi^*.\partial\phi - \frac{i}{2}g_2\phi^*\overleftrightarrow{\partial}\phi.D + \frac{i}{2}g_2\phi^*\overleftrightarrow{\partial}\phi.C +$$

$$+ \frac{1}{4}g_2^2\phi^*\phi D^2 + \frac{1}{4}g_2^2\phi^*\phi C^2 - \frac{1}{2}g_2^2\phi^*\phi C.D \quad . \qquad (2.19)$$

The presence of the gauge-invariant massive field C_μ in our model may in principle lead to the suspicion that it lacks of the renormalizability property. Let us then consider this issue more carefully.

The Euler-Lagrange equations of motion for the physical vector fields are

$$\partial^\mu D_{\mu\nu} = \frac{1}{2}g_1\bar{\psi}\gamma_\mu\psi + \frac{i}{g}g_2\phi^*\overleftrightarrow{\partial}_\nu\phi - \frac{1}{2}g_2^2\phi^*\phi D_\nu +$$

$$+ \frac{1}{2}g_2^2\phi^*\phi C_\nu \equiv J_\nu \qquad (2.20)$$

and

$$\partial^\mu C_{\mu\nu} = -m^2C_\nu + \frac{1}{2}g_1\bar{\psi}\gamma_\nu\psi - \frac{i}{g}g_2\phi^*\overleftrightarrow{\partial}_\nu\phi +$$

$$+ \frac{1}{2}g_2^2\phi^*\phi D_\nu - \frac{1}{2}g_2^2\phi^*\phi C_\nu \equiv -m^2C_\nu + \mathcal{J}_\nu \, . \qquad (2.21)$$

The current $\mathcal{J}_\mu$ to which the Proca field couples is not necessari ly conserved and we could conclude by the non-renormalizability of the model.

However, let us recall that we started from two global U(1) symmetries, of which only one was truly gauged and this is supported by the existence of a single genuine gauge field, D_μ. We have still to exploit the remaining global U(1) invariance we have in our model.

Notice that if the fields ψ and ϕ have opposite values of the global U(1) charge, the globally conserved current turns out to be

$$j = \bar{\psi}\gamma\psi - i\phi^*\overleftrightarrow{\partial}\phi + g_2\phi^*\phi D - g_2\phi^*\phi C \; . \tag{2.22}$$

So, if we choose the coupling constants g_1 and g_2 equal,

$$g_1 = g_2 \equiv g \; , \tag{2.23}$$

the current j_μ is nothing but the current J_μ up to an overall factor and (2.21) becomes

$$\partial^\mu C_{\mu\nu} = m^2 C_\nu + j_\nu \; , \tag{2.24}$$

with

$$\partial . j = 0 \; , \tag{2.25}$$

that is, C_μ couples to a conserved current. Consequently, our model enjoys the property of being renormalizable [8], provided we take

$g_1=g_2$ and matter fields with opposite values of the global U(1) charge, as this leads to a decoupling of the spin-zero mode carried by C_μ from the physical degrees of freedom. In the path integral formalism, this means that one can suitably redefine the matter fields and then integrate over the longitudinal component of C_μ, ending up with an effective interacting theory for the new matter fields and the transverse component of C_μ.

As a conclusion, the classical Lagrangian density we arrive at is

$$\begin{aligned}\mathcal{L} = & -\frac{1}{4}D^{\mu\nu}D_{\mu\nu} - \frac{1}{4}C^{\mu\nu}C_{\mu\nu} + \frac{1}{2}m^2C^2 + \\ & + \overline{\psi}\gamma.i\partial\psi - \frac{1}{2}g\,\overline{\psi}\gamma\psi.D - \frac{1}{2}g\,\overline{\psi}\gamma\psi.C + \\ & + \partial\phi^*.\partial\phi - \frac{i}{g}g\,\phi^*\overleftrightarrow{\partial}\phi.D + \frac{i}{g}g\,\phi^*\overleftrightarrow{\partial}\phi.C + \\ & + \frac{1}{4}g^2\,\phi^*\phi D^2 + \frac{1}{4}g^2\,\phi^*\phi C^2 - \frac{1}{2}g^2\,\phi^*\phi D.C\ ,\end{aligned} \tag{2.26}$$

which is $U(1)_{local} \otimes U(1)_{global}$ invariant.

The fields D_μ and C_μ couple respectively to the currents J^μ and j^μ. D_μ could be interpreted as the photon field (universality of the coupling g) and J^μ as the electromagnetic current. The conserved global U(1) charge can be taken as

$$Q = F - B\ , \tag{2.27}$$

where F and B are respectively a fermionic and a bosonic number.

3. Quantum Properties

To study the behaviour of the model previously discussed upon quantum corrections, we have first of all to adjoin the gauge-fixing term for the "physical" gauge field of the theory. We choose to work in the covariant class of gauges,

$$\mathcal{L}_{g.f.} = -\frac{1}{2\alpha}(\partial.D)^2 \ , \qquad (3.1)$$

and the Faddeev-Popov ghosts, as in the usual U(1) case, completely decouple from the other fields.

Recalling that the massive field C_μ couples to a conserved current, the power-counting formula can be readily derived. The superficial degree of divergence, δ, of the primitively divergent graphs is dictated by the expression

$$\delta = 4 - E_D - E_C - \frac{3}{2} E_\psi - E_\phi \ , \qquad (3.2)$$

which evidenciates the power-counting renormalizability of the model. In (3.2), E stands for the number of external legs corresponding to the field appearing as a subscript.

In the task of computing quantum loop corrections, we can choose to work with either the Lagrangian (2.11) with $g_1=g_2=g$ or the one given by the expression (2.26). In the latter picture, the matter-matter interactions can be easily observed as both ψ and ϕ couple to both D_μ and C_μ. In the former representation, the fermionic matter couples only to A_μ whereas the bosonic one couples only to B_μ. The ψ-ϕ interaction can however be visualized through a mediating mixed A_μ-B_ν-propagator, whose explicit form is

$$\langle T(A_\mu B_\nu)\rangle = \frac{1}{2}\frac{m^2}{k^2(k^2+m^2)}\eta_{\mu\nu} + \frac{1}{2m^2}\frac{k^4+(1-\alpha)m^2(k^2+m^2)}{k^4(k^2+m^2)}k_\mu k_\nu , \qquad (3.3)$$

already in the Wick-rotated momentum space. Notice the positions of the poles at $k^2=0$ and $k^2=-m^2$, indicating the presence of a massless gauge boson and a massive vector particle. Though in this picture the vertices read very easily, the expressions for the propagators are not suitable for computations beyond the tree--level.

The power-counting formula (3.2) indicates that divergent quantum corrections of the type $C^{\mu\nu}D_{\mu\nu}$, $(C^\mu C_\mu)^2$ and $(\phi^*\phi)^2$ can be induced already at the one-loop level, whose superficial degrees of divergence are respectively quadratic and logarithmic. At this point, we should consult the Ward identities of the theory.

If we denote by $\Gamma[C_\mu,D_\nu;\overline{\psi},\psi;\phi^*,\phi]$ the effective action of the theory (generating functional of the 1P.I. diagrams), the Ward identities it obeys is given by

$$\frac{1}{\alpha}\Box\partial.D(x) - \partial.\frac{\delta\Gamma}{\delta D(x)} + ig\overline{\psi}(x)\frac{\delta\Gamma}{\delta\psi(x)} +$$

$$-ig\psi(x)\frac{\delta\Gamma}{\delta\overline{\psi}(x)} + ig\phi^*(x)\frac{\delta\Gamma}{\delta\phi^*(x)} - ig\phi(x)\frac{\delta\Gamma}{\delta\phi(x)} = 0 \qquad (3.4)$$

Besides the usual non-renormalization of the longitudinal component of D_μ, they also imply that a term of the form $C^{\mu\nu}D_{\mu\nu}$ must necessarily have the form

$$C^\mu(-k)(k^2\eta_{\mu\nu}-k_\mu k_\nu)\,\Pi(k)\,D^\nu(k) \quad , \qquad (3.5)$$

with $\Pi(k)$ being only logarithmically divergent. Since $C^{\mu\nu}D_{\mu\nu}$ is a gauge-invariant quantity, one could include at the tree-level Lagrangian a term

$$\lambda C^{\mu\nu}D_{\mu\nu} ,$$

and then reabsorb this infinity into a logarithmic renormalization of the dimensionless parameter λ.

As for the corrections of the form $(C^{\mu}C_{\mu})^{4}$ and $(\phi^{*}\phi)^{2}$, gauge invariance does not reduce their superficial degree of divergence; a correction-term like $(C^{\mu\nu}D_{\mu\nu})^{2}$ can be induced, but it contributes a finite amount to the effective action.

As a final result following from the analysis of the Ward identities, by taking functional derivatives with respect to suitable classical fields, one can obtain useful relationships among the various wave-function and vertex renormalization factors. These relations ensure that all gauge couplings renormalize effectively in the same way, that is, the final Lagrangian accounting for all the renormalizations exhibit only one renormalized coupling constant, g_k, as it should be. This consequently indicates that the renormalized version of the theory is indeed gauge invariant. The outcome of this analysis is that our $U(1)_{local} \otimes U(1)_{global}$ - invariant theory exhibits the desirable features of renormalizability and unitarity.

4. Supersymmetric Proposals

As the theory we are trying to establish starts from fermions and bosons and they transform equally under the action of the U(1) gauge group, let us go a little further with the similarities between the fermionic and bosonic sectors and let us look at a possible supersymmetric extension of the Lagrangian (2.26).

Let ϕ_1 and ϕ_2 be the chiral scalar superfields which accomodate ψ and ϕ respectively. They are parametrized by the following θ-expansions:

$$\Phi_1 = e^{\theta\sigma\overline{\theta}.i\partial}(z + \theta\psi + \theta^2 h) \qquad (4.1)$$

and

$$\Phi_2 = e^{\theta\sigma\overline{\theta}.i\partial}(\phi + \theta\chi + \theta^2 g) \quad , \qquad (4.2)$$

where z and χ are the physical supersymmetric partners of ψ and ϕ respectively, whereas h and g are complex auxiliary fields.

The gauge field D_μ and the massive vector boson C_μ are located in the vector superfields V and U respectively, whose complete θ-expansions are given below:

$$V = R + \theta\chi + \overline{\theta}\overline{\chi} + \theta^2 M + \overline{\theta}^2 M^* + \theta\sigma\overline{\theta}.D + \theta^2\overline{\theta}(\overline{\lambda} + \frac{i}{2}\partial\chi.\sigma) +$$
$$+ \overline{\theta}^2\theta(\lambda + \frac{i}{2}\sigma.\partial\overline{\chi}) + \theta^2\overline{\theta}^2(E + \frac{1}{4}\Box A) \qquad (4.3)$$

and

$$U = S + \theta\xi + \overline{\theta}\overline{\xi} + \theta^2 N + \overline{\theta}^2 N^* + \theta\sigma\overline{\theta}.C + \theta^2\overline{\theta}(\overline{\xi} + \frac{i}{2}\partial\xi.\sigma)$$

$$+ \overline{\theta}^2\theta(\xi + \frac{i}{2}\sigma.\partial\overline{\xi}) + \theta^2\overline{\theta}^2(F + \frac{1}{4}\Box B) \quad . \qquad (4.4)$$

The physical components of V are the gauge boson D_μ and the gauge fermion λ; R, χ and M are pure gauge modes: they can be eliminated by a suitable choice of the gauge parameters, in this case the components of an arbitrary chiral scalar superfield. As for the vector superfield U, it is a gauge-invariant quantity so that it cannot obviously accomodate any pure gauge mode. Its 8 fermionic and 8 bosonic degrees of freedom are distributed among physical an auxiliary components: the pseudoscalar m S, the Majorana spinors $(m\xi)$ and ξ, and the transverse component of C_μ describe physical particles of mass m (the longitudinal component of C_μ is a ghost which decouples from the physical sector of the Fock space), whereas N, N^* and the pseudoscalar F play the rôle of auxiliary fields.

A straightforward supersymmetric version of the Lagrangian (2.26) could well be given by

$$\mathcal{L} = \int d^4\theta \, [-\frac{1}{16}(V D^\alpha \overline{D}^2 D_\alpha V + U D^\alpha \overline{D}^2 D_\alpha^U) +$$

$$- \frac{1}{2} m^2 U^2 - \frac{1}{16\alpha}(D^2 V)(\overline{D}^2 V) +$$

$$+ \overline{\Phi}_1 e^{g(V-U)} \Phi_1 + \overline{\Phi}_2 e^{-g(V-U)} \Phi_2] \quad . \qquad (4.5)$$

However, the non-polynomial interactions of the pseudoscalar physical component of U with the matter fields (coming from e^{gU} and e^{-gU}) leads to the non-renormalizability of the Lagrangian (4.5), as U does not obey even some broken Ward identities (remember that

U is a gauge-invariant object).

Then, to overcome the problems stemming from non-polynomial interactions, we propose the following supersymmetric Lagrangian:

$$\mathcal{L} = \int d^4\theta \, [-\frac{1}{16} (V D^{\alpha}\bar{D}^2 D_{\alpha} V + U D^{\alpha}\bar{D}^2 D_{\alpha} U) - \frac{1}{2} m^2 U^2 -$$

$$- \frac{1}{16\alpha} (D^2 V)(\bar{D}^2 V) + (\bar{\Phi}_1 e^{gV} \Phi_1 + \bar{\Phi}_2 e^{-gV} \Phi_2)(1 - gU + \frac{1}{2} g^2 U^2) \qquad (4.6)$$

which, upon a component-wise analysis, can be seen to reproduce the Lagrangian of expression (2.26) and exhibits vertices of the renormalizable kind.

5. Conclusions

We have in this article considered the possibility of establishing a renormalizable and unitary gauge theory with the unusual feature of introducing two different gauge potentials in association with a single gauge group, each of them interacts with a different family of matter fields. Such an attitude naturally leads to the emergence of a massive vector boson accompanied by a genuine massless) gauge field: this was the actual motivation behind the whole approach.

This picture may have some interest to the physics of electroweak interactions, where massive and massless vector bosons are present in the spectrum of physical particles and mediate the decay and scattering processes. However, for the time being, ours is not yet a very realistic claim, as we have been able to give a meaning to our attempt only in the Abelian case. Thus a physical aspect for (2.16) and (2.17) is to interpret them as massive and massless photons respectively. Although experiments as the geomagnetic data sets limits on the photon mass [9], there is no intuitive arguments against it. The model in (2.11) can also be extended for fermions fields, i.e., one can use a fermion $\chi(x)$ instead of $\phi(x)$. Observe that there are other models looking for a massive photon [10], where gauge symmetry can be broken explicitly or not. The contribution of this paper was that of generating healthy massive photon without gauge symmetry breaking.

The mismatch between the gauge freedom and the number of gauge potentials implies an abundance of spin-zero modes which may plague the theory with negative-metric ghost states. However, in the U(1) case, we could verify that the theory naturally offers the freedom for decoupling these modes from the S-matrix elements through the coupling of the massive vector boson of the theory to a globally conserved matter current. One can then end up with a sensible theory: renormalizable and unitary.

The most interesting case of non-Abelian gauge groups is now under investigation. One can readily realize that in such a case, we cannot trivially follow the same patterns we established in this paper: due to the presence of the three- and four-vector couplings, the undesirable ghost carried by the massive vector field do not decouple from the physical degrees of freedom, through they can decouple from the matter sector. Any progress in this direction shall be reported in a further publication.

Acknowledgements

It is for us a great pleasure to express our gratitude to Prof. M. Tonin, Prof. M. Flato, Prof. R. Iengo and Prof. Dai Yuan-ben for the friendly and very instructive discussions on various aspects of quantum gauge theories.

Both of us (R.M.D.) and (A.W.S.) would like to thank Prof. Jayme Tiomno for the hospitality at CBPF and the financial support of Conselho Nacional de Desenvolvimento Científico e Tecnológico (CNPq).

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