Non-Abelian (2,0)-Super-Yang-Mills Coupled to Non-Linear σ-Models

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Considering a class of (2,0)-super-Yang-Mills multiplets that accommodate a pair of independent gauge potentials in connection with a single symmetry group, we present here their non-Abelian coupling to ordinary matter and to non-linear σ-models in (2,0)-superspace. The dynamics and the couplings of the gauge potentials are discussed and the interesting feature that comes out is a sort of "chirality" for one of the gauge potentials whenever light-cone coordinates are chosen.

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In a previous paper [1], we discussed the dynamics and the couplings of the Abelian vector potentials of a class of $(2,0)$-gauge super multiplets([2]-[8]) in connection with a single U(1)-symmetry group. Since a number of interesting features came out, it was a natural question to ask how these fields would behave if the non-Abelian version of the theory was to be considered.

We can see that indeed some subtle changes occur. As we wish to make a full comparison between the two aspects(Abelian and non-Abelian) of the same sort of theory, all the characteristics of the original formulation were kept, namely, the coordinates we choose to parametrise the $(2,0)$-superspace are given by:

$$z^A \equiv (x^{++}, x^{--}; \theta, \bar{\theta}), \tag{1}$$

where x^{++}, x^{--} denote the usual light-cone variables, whereas θ, $\bar{\theta}$ stand for complex right-handed Weyl spinors. The supersymmetry covariant derivatives are taken as:

$$D_+ \equiv \partial_\theta + i\bar{\theta}\partial_{++} \tag{2}$$

and

$$\bar{D}_+ \equiv \partial_{\bar{\theta}} + i\theta\partial_{++}, \tag{3}$$

where ∂_{++} (or ∂_{--}) represents the derivative with respect to the space-time coordinate x^{++} (or x^{--}). They fulfill the algebra:

$$D_+^2 = \bar{D}_+^2 = 0 \qquad \{D_+, \bar{D}_+\} = 2i\partial_{++}. \tag{4}$$

With these definitions for D and $\bar{D}$, one can check that:

$$e^{i\theta\bar{\theta}\partial_+} D_+ e^{-i\theta\bar{\theta}\partial_+} = \partial_\theta, \tag{5}$$

$$e^{-i\theta\bar{\theta}\partial_+} \bar{D}_+ e^{i\theta\bar{\theta}\partial_+} = \partial_{\bar{\theta}}. \tag{6}$$

The fundamental non-Abelian matter superfields are the "chiral" scalar and left-handed spinor superfields, whose respective component-field expressions are given by:

$$\begin{aligned}\Phi^i(x;\theta,\bar{\theta}) &= e^{i\theta\bar{\theta}\partial_{++}}(\phi^i + \theta\lambda^i),\\ \Psi^I(x;\theta,\bar{\theta}) &= e^{i\theta\bar{\theta}\partial_{++}}(\psi^I + \theta\sigma^I),\end{aligned} \tag{7}$$

the fields ϕ^i and σ^I are scalars, whereas λ^i and ψ^I stand respectively for right- and left-handed Weyl spinors. The indices i and I label the representations where the correspondenting matter fields are set.

We present below the gauge transformations of both Φ and Ψ, assuming that we are dealing with a compact and simple gauge group, $\mathcal{G}$, with generators G_a that fulfill the algebra $[G_a, G_b]=if_{abc}G_c$. The transfomations read:

$$\Phi'^i = R(\Lambda)^i_j \Phi^j, \qquad \Psi'^I = S(\Lambda)^I_J \Psi^J, \tag{8}$$

where R and S are matrices that respectively represent a gauge group element in the representations under which Φ and Ψ transform. Taking into account the chiral constraint on Φ and Ψ, and bearing in mind the exponential representation for R and S in terms of the group generators, we find that the gauge parameter superfields, Λ^a, must satisfy the same sort of constraint. They can therefore be expanded as follows:

$$\Lambda^a(x;\theta,\bar{\theta}) = e^{i\theta\bar{\theta}\partial_{++}}(\alpha^a + \theta\beta^a), \tag{9}$$

where α^a are scalars and β^a are right-handed spinors.

The kinetic action for Φ^i and Ψ^I can be made invariant under the local transformations (8) by minimally coupling gauge potential superfields, $\Gamma^a_{--}(x;\theta,\bar{\theta})$ and $V^a(x;\theta,\bar{\theta})$, according to the minimal coupling prescriptions:

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$$S_{inv}=\int d^2xd\theta d\bar{\theta}\{i[\bar{\Phi}e^{hV}(\nabla_{--}\Phi)-(\bar{\nabla}_{--}\bar{\Phi})e^{hV}\Phi]+\bar{\Psi}e^{hV}\Psi\}, \tag{10}$$

where the gauge-covariant derivatives are defined in the sequel.

The Yang-Mills supermultiplets are introduced by means of the gauge-covariant derivatives which, according to the discussion of ref. [7], are defined as below:

$$\nabla_+ \equiv D_+ + \Gamma_+, \qquad \bar{\nabla}_+ \equiv \bar{D}_+, \tag{11}$$

$$\nabla_{++} \equiv \partial_{++} + \Gamma_{++} \quad and \quad \nabla_{--} \equiv \partial_{--} - ig\Gamma_{--}, \tag{12}$$

with the gauge superconnections Γ_+, Γ_{++} and Γ_{--} being all Lie-algebra-valued. Note that Γ_{++} does not enter the Lagrangian density of eq.10). The gauge couplings, g and h, can in principle be taken different; nevertheless, this would not mean that we are gauging two independent symmetries. There is a single simple gauge group, $\mathcal{G}$, with just one gauge-superfield parameter, Λ. It is the particular form of the $(2,0)$-minimal coupling (realised by the exponentiation of V and the connection present in ∇_{--}) that opens up the freedom to associate, in principle, different coupling parameters to the gauge superfields V and Γ_{--}. Γ_+ and Γ_{++} can be both expressed in terms of the real scalar superfield, $V(x;\theta,\bar{\theta})$, according to [1]:

$$\Gamma_+ = e^{-gV}(D_+e^{gV}) \tag{13}$$

and

$$\Gamma_{++} = -\frac{i}{2}\bar{D}_+[e^{-gV}(D_+e^{gV})]. \tag{14}$$

To establish contact with a component-field formulation and to actually identify the presence of an additional gauge potential, we write down the θ-expansions for V^a and Γ^a_{--}:

$$V^a(x;\theta,\bar{\theta}) = C^a + \theta\xi^a - \bar{\theta}\bar{\xi}^a + \theta\bar{\theta}v^a_{++} \tag{15}$$

and

$$\begin{aligned}\Gamma^a_{--}(x;\theta,\bar{\theta}) &= \frac{1}{2}(A^a_{--}+iB^a_{--}) + i\theta(\rho^a + i\eta^a)\\ &+ i\bar{\theta}(\chi^a + i\omega^a) + \frac{1}{2}\theta\bar{\theta}(M^a + iN^a)\end{aligned} \tag{16}$$

A^a_{--}, B^a_{--} and v^a_{++} are the light-cone components of the gauge potential fields; ρ^a, η^a, χ^a and ω^a are left-handed Majorana spinors; M^a, N^a and C^a are real scalars and ξ^a is a complex right-handed spinor.

The infinitesimal gauge transfomations for V^a and Γ^a are given by

$$\delta V^a = \frac{i}{h}(\bar{\Lambda}-\Lambda)^a - \frac{1}{2}f^{abc}(\bar{\Lambda}+\Lambda)_bV_c \tag{17}$$

and

$$\delta\Gamma^a_{--} = -f^{abc}\Lambda_b\Gamma_{c--} + \frac{1}{g}\partial_{--}\Lambda_a. \tag{18}$$

No derivative acts on the $\Lambda)^a$'s in eq.(17), which suggests the possibility of choosing a Wess-Zumino gauge for V^a. It such a choice is adopted, it can be shown that the gauge transformations of the θ-component fields above read as follows:

$$\begin{aligned}
&= \frac{2}{h}\Im m\alpha,\\
\delta v^a_{++} &= \frac{2}{h}\partial_{++}\alpha^a - f_{abc}\alpha^b v^c_{++},\\
\delta A^a_{--} &= \frac{2}{g}\partial_{--}\alpha^a - f_{abc}\alpha^b A^c_{--},,\\
\delta B^a_{--} &= -f_{abc}\alpha^b B^c_{--},\\
\delta\eta^a &= -f_{abc}\alpha^b\eta^c,\\
\delta\rho^a &= -f_{abc}\alpha^b\rho^c,\\
\delta M^a &= -f_{abc}\alpha^b M^c + f_{abc}\partial_{++}\alpha^b B^c_{--},\\
\delta N^a &= \frac{2}{g}\partial_{++}\partial_{--}\alpha^a - f_{abc}\alpha^b N^c - f_{abc}\partial_{++}\alpha^b A^c_{--},\\
\delta\chi^a &= -f_{abc}\alpha^b\chi^c,\\
\delta\omega^a &= -f_{abc}\alpha^b\omega^c,
\end{aligned} \tag{19}$$

and they suggest that we should take $h=g$, so that the v^a_{++}-component could be identified as the light-cone partner of A^a_{--},

$$v^a_{++} \equiv A^a_{++}; \tag{20}$$

this procedure yields two component-field gauge potentials: $A^{\mu} \equiv (A^0, A^1) = (A^{++}; A^{--})$ and B_{--}.

It is interesting to point out here that the first difference between the Abelian and the non-Abelian version of the theory arises. In the Abelian version [1], it was shown that both fields χ and ω were gauge invariant and the fields M and N could be identified with a combination of A_{--} and B_{--}. This combination, which was naturally dictated by the form of the gauge transformations, ensured the symmetry of the Lagrangian. In the present situation, the gauge transformations do not undertake that we express M and N in terms of A_{--} and B_{--}, as it was done before but; on the other hand, the χ-and ω-fields are no longer auxiliary fields as they were in the Abelian version.

To discuss the field-strength superfields, we start analysing the algebra of the gauge covariant derivatives. So, the field strengths are defined such that:

$$
\begin{aligned}
\{\nabla_+, \nabla_+\} &\equiv \mathcal{F} = 2D_+\Gamma_+, \\
\{\nabla_+, \bar{\nabla}_+\} &\equiv 2i\nabla_{++}, \\
[\nabla_+, \nabla_{++}] &\equiv W_- = D_+\Gamma_{++} - \partial_{++}\Gamma_+, \\
[\nabla_+, \nabla_{--}] &\equiv W_+ = -igD_+\Gamma_{--} - \partial_{++}\Gamma_+ - ig[\Gamma_+, \Gamma_{--}], \\
[\bar{\nabla}_+, \nabla_{++}] &\equiv U_+, \\
[\bar{\nabla}_+, \nabla_{--}] &\equiv U_- = -ig\bar{D}_+\Gamma_{--}, \\
[\nabla_{++}, \nabla_{--}] &\equiv \mathcal{Z}_{+-} = -ig\partial_{++}\Gamma_{--} - \partial_{--}\Gamma_{++} - ig[\Gamma_+, \Gamma_{--}].
\end{aligned}
\tag{21}
$$

The results obtained for the field-strengths are consistent with the Bianchi identities. The identity for U_+,

$$[\bar{\nabla}_+, \{\nabla_+, \bar{\nabla}_+\}] + [\nabla_+, \{\bar{\nabla}_+, \bar{\nabla}_+\}] + [\bar{\nabla}_+, \{\bar{\nabla}_+, \nabla_+\}] = 0 \tag{22}$$

gives immediately that $U_+ = 0$. The Bianchi identity for Z_{+-},

$$[\nabla_{--}, \{\nabla_+, \bar{\nabla}_+\}] + \{\nabla_+, [\bar{\nabla}_+, \nabla_{--}]\} - \{\bar{\nabla}_+, [\nabla_{--}, \bar{\nabla}_+]\} = 0, \tag{23}$$

allows us to express Z_{+-} as

$$Z_{+-} = -\frac{i}{2}\nabla_+U_- - \frac{i}{2}\bar{\nabla}_+W_-; \tag{24}$$

and, finally, the Bianchi identity

$$[\bar{\nabla}_+, \{\nabla_+, \nabla_+\}] + [\nabla_+, \{\nabla_+, \bar{\nabla}_+\}] + [\nabla_+, \{\bar{\nabla}_+, \nabla_+\}] = 0 \tag{25}$$

leads to

$$W_+ = \frac{i}{4}\bar{D}_+\mathcal{F}. \qquad (26)$$

These are the relevant results yielded by pursuing an investigation of the Bianchi identities.

The gauge field, A_μ, has its field strength, $F_{\mu\nu}$, located at the θ-component of the combination $\Omega \equiv W_- + \bar{U}_-$. This suggests the following kinetic action for the Yang-Mills sector:

$$\begin{aligned} S_{YM} &= \frac{1}{8g^2}\int d^2x d\theta d\bar{\theta} Tr \Omega\bar{\Omega} \\ &= \int d^2x Tr[\frac{-1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{i}{4}\Sigma \overleftrightarrow{\partial}_{++} \bar{\Sigma} + \frac{1}{4}M^2], \end{aligned} \qquad (27)$$

where $\Sigma = \rho + i\eta + \bar{\chi} - i\bar{\omega}$ and $A \overleftrightarrow{\partial} B \equiv (\partial A)B - A(\partial B)$.

Choosing now a supersymmetry-covariant gauge-fixing, ins tead of the Wess-Zumino, we propose the following gauge-fixing term in superspace:

$$\begin{aligned} S_{gf} &= -\frac{1}{2\alpha}\int d^2x d\theta d\bar{\theta} Tr[\Pi\bar{\Pi}] \\ &= -\frac{1}{2\alpha}\int d^2x\{[(\partial_\mu A^\mu)^2 + (\partial_\mu A^\mu)N + \frac{1}{4}N^2] \\ &+ \frac{1}{4}[M^2 - 2M\partial_{++}B_{--} + (\partial_{++}B_{--})^2] \\ &- i(\rho + i\eta) \overleftrightarrow{\partial}_{++} (\bar{\rho} - i\bar{\eta})\}, \end{aligned} \qquad (28)$$

where $\Pi = -iD_+\Gamma_{--} + \frac{1}{2}D_+\partial_{--}V$.

So, the total action reads as follows:

$$\begin{aligned} S &= \int d^2x Tr\{-\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - \frac{1}{2\alpha}(\partial_\mu A^\mu)^2 - \frac{1}{2\alpha}(\partial_\mu A^\mu)N - \frac{1}{8\alpha}N^2 \\ &+ \frac{1}{4}(1 - \frac{1}{2\alpha})M^2 + \frac{1}{4\alpha}M(\partial_{++}B_{--}) - \frac{1}{8\alpha}(\partial_{++}B_{--})^2 \\ &+ \frac{i}{2\alpha}(\rho + i\eta) \overleftrightarrow{\partial}_{++} (\bar{\rho} - i\bar{\eta}) + \frac{i}{4}\Sigma \overleftrightarrow{\partial}_{++} \bar{\Sigma}\}. \end{aligned} \qquad (29)$$

Using eq.(29), we are ready to write down the propagators for A^a, B^a_{--}, N^a, M^a, ρ^a, η^a, χ^a and ω^a:

$$\begin{aligned} \langle AA \rangle &= -\frac{i}{2\Box(1-\Box)}\omega^{\mu\nu}, \\ \langle BB \rangle &= -\frac{i(2\alpha-1)}{4\alpha(1-\alpha)}\frac{\partial^2_{--}}{\Box^2}, \\ \langle AN \rangle &= \frac{i\alpha}{\Box(1-\Box)}(1-\Box+\alpha)\partial^\mu, \\ \langle NA \rangle &= \frac{i}{(1-\Box)\Box}\partial^\nu \\ \langle NN \rangle &= -\frac{2i\alpha^2}{(1-\Box)} \\ \langle MM \rangle &= -\frac{i}{16\alpha}\frac{1}{(1-\alpha)} \\ \langle MB \rangle &= -\langle BM \rangle = \frac{i}{8\alpha(1-\alpha)}\frac{\partial_{--}}{\Box} \\ \langle(\rho+i\eta)(\bar{\rho}-i\bar{\eta})\rangle &= -\frac{2\alpha}{(\alpha-1)}\frac{\overleftrightarrow{\partial}_{--}}{4\Box} \\ \langle(\rho+i\eta)(\chi+i\omega)\rangle &= -\frac{\alpha}{4}\frac{\overleftrightarrow{\partial}_{--}}{\Box} \\ \langle(\bar{\chi}-i\bar{\omega})(\bar{\rho}-i\bar{\eta})\rangle &= +\frac{\alpha}{4(\alpha-1)}\frac{\overleftrightarrow{\partial}_{--}}{\Box} \\ \langle(\bar{\chi}-i\bar{\omega})(\chi+i\omega)\rangle &= +\frac{(\alpha+2)}{4}\frac{\overleftrightarrow{\partial}_{--}}{\Box}. \end{aligned} \qquad (30)$$

Expressing the action of equation (10) in terms of component fields, and coming back to the $(2,0)$-version of the Wess-Zumino gauge, the matter-gauge sector Lagrangian reads:

$$\begin{aligned}\mathcal{L}_{matter-gauge} &= 2\phi^{*i}\Box\phi_i - ig[\phi^{*i}A^a_{--}(G_a)^j_i\partial_{++}\phi_j - c.c] + \bar{\sigma}^i\sigma_i + \\ &- ig[\phi^{*i}A^a_{++}(G_a)^j_i\partial_{--}\phi_j - c.c] - g\phi^{*i}M^a(G_a)^j_i\phi_j + \\ &- \frac{i}{2}g^2\phi^{*i}A^a_{++}A^b_{--}\phi_i d_{abc}G_c - g\bar{\lambda}^i A^a_{--}(G_a)^j_i\lambda_j + \\ &- \frac{1}{2}\phi^{*i}A^a_{++}B^b_{--}\phi_i f_{abc}G_c + 2i\bar{\lambda}^i\partial_{--}\lambda_i + \\ &- ig\phi^{*i}[(\chi^a + \bar{\rho}^a + i\omega^a - i\bar{\eta}^a)(G_a)^j_i\lambda_j - c.c] + \\ &- 2i\bar{\psi}^i\partial_{++}\psi_i - g\bar{\psi}^i A^a_{++}(G_a)^j_i\psi_j,\end{aligned} \tag{31}$$

where d_{abc} are the(representation-dependent) symmetric coefficients associated to $\{G_a, G_b\}$.

One immediately checks that the extra gauge field, B_{--}, does not decouple from the matter sector. Our point of view of leaving the superconnection Γ_{--} as a complex superfield naturally introduced this extra gauge potential in addition to the usual gauge field, A_μ: B_{--} behaves as a second gauge field. The fact that it yelds a massless pole of order two in the spectrum may harm unitarity. However, the mixing with the M-component of Γ_{--}, which is a compensating field, indicates that we should couple them to external currents and analyse the imaginary part of the current-current amplitude at the pole. In so doing, this imaginary part turns out to be positive-definite, and so no ghosts are present. It is very interesting to point out that, in the Abelian case, B_{--} showed the same behaviour [1]. It coupled to C instead of M, but these two fields show the same kind of behaviour: C (in the Abelian case) and M (in the non-Abelian case) are both compensating fields. This ensures us to state that B_{--} behaves as a physical gauge field: it has dynamics and couples both to matter and the gauge field A^μ. Its only peculiarity regards the presence of a single component in the light-cone coordinates. The B-field plays rather the rôle of a "chiral gauge potential". Despite the presence of the pair of gauge fields, a gauge-invariant mass term cannot be introduced, since B does not carry the B_{++}-component, contrary to what happens with A^μ.

Let us now turn to the coupling of the two gauge potentials, A_μ and B_{--}, to a non-linear σ-model always keeping a sypersymmetric scenario. It is our main purpose henceforth to carry out the coupling of a $(2,0)$ σ-model to the relaxed gauge superfields of the ref. [7], and show that the extra vector degrees of freedom do not decouple from the matter fields (that is, the target space coordinates)[9][10][11][12]. The extra gauge potential, B_{--}, obtained upon relaxing constraints can therefore acquire a dynamical significance by means of the coupling between the σ-model and the Yang-Mills fields of ref.[7]. To perform the coupling of the σ-model to the Yang-Mills fields we reason along the same considerations as i ref.[1] and find out that:

$$\begin{aligned}\mathcal{L}_\xi &= \partial_i[K(\Phi,\tilde{\Phi}) - \xi(\Phi) - \tilde{\xi}(\tilde{\Phi})]\nabla_{--}\Phi^i + \\ &- \tilde{\partial}_i[K(\Phi,\tilde{\Phi}) - \xi(\Phi) - \tilde{\xi}(\tilde{\Phi})]\nabla_{--}\tilde{\Phi}^i,\end{aligned} \tag{32}$$

where $\xi(\Phi)$ and $\bar{\xi}(\bar{\Phi})$ are a pair of *chiral* and *antichiral* superfields, $\tilde{\Phi}_i \equiv exp(i\mathbf{L}_{V.\bar{k}})\bar{\Phi}_i$ and $\nabla_{--}\Phi^i$ and $\nabla_{--}\tilde{\Phi}^i$ are defined in perfect analogy to what is done in the case of the bosonic σ-model:

$$\nabla_{--}\Phi_i \equiv \partial_{--}\Phi_i - g\Gamma^\alpha_{--}k^i_\alpha(\Phi) \tag{33}$$

and

$$\nabla_{--}\tilde{\Phi}_i \equiv \partial_{--}\tilde{\Phi}_i - g\Gamma^\alpha_{--}\bar{k}_{\alpha i}(\tilde{\Phi}). \tag{34}$$

The interesting point we would like to stress is that the extra gauge degrees of freedom accommodated in the component-field $B_{--}(x)$ of the superconnection Γ_{--} behave as a genuine gauge field that shares with A^μ the feature of coupling to matter and to σ-model [7]. This result can be explicitly read off from the component-field Lagrangian projected out from the superfield Lagrangian $\mathcal{L}_\xi$. We therefore conclude that our less constrained $(2,0)$-gauge theory yields a pair of gauge potentials that naturally transform under the action of a single compact and simple gauge group and may be consistently coupled to matter fields as well as to the $(2,0)$ non-linear σ-models by means of the gauging of their isotropy and isometry groups. Relaxing constraints in the $N = 1$- and $N = 2 - D = 3$ supersymmetric algebra of covariant derivatives may lead to a number of peculiar features of the gauged $O(3)$-σ-model[13] in the presence of Born-Infeld terms for the pair of gauge potentials; of special interest are the self-dual equations[14].

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