The Effective Action
For
Brane Localized Gauge Fields

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Abstract

The low energy effective action including gauge field degrees of freedom on a non-BPS $p=2$ brane embedded in a $N=1$, $D=4$ target superspace is obtained through the method of nonlinear realizations of the associated super-Poincaré symmetries. The invariant interactions of the gauge fields and the brane excitation modes corresponding to the Nambu-Goldstone degrees of freedom resulting from the broken space translational symmetry and the target space supersymmetries are determined. Brane localized matter field interactions with the gauge fields are obtained through the construction of the combined gauge and super-Poincaré covariant derivatives for the matter fields.
When a target manifold contains an embedded defect its symmetry group is spontaneously broken to those of the world volume of the defect and its complement. The low energy quantum fluctuations of the defect in directions associated with broken symmetry generators correspond to Nambu-Goldstone degrees of freedom. It is well known that defects, such as an embedded domain wall, may cause the localization of additional massless as well as massive scalar and fermionic degrees of freedom on their world volume \[1, 2, 3\]. The idea that such localized degrees of freedom constitute the matter of our universe is explored in brane-world scenarios \[4, 5\]. While in a field theoretical context defects embedded in a flat (non-compact) target space do not yield localized gauge fields as zero modes \[6\], several alternative mechanisms that result in (quasi-) localized non-Abelian gauge fields on the defect have been proposed \[7, 8, 9\]. (Kaluza-Klein localization of the gauge field zero mode for compact or finite volume spaces is readily realized \[10\].) In any realistic brane-world scenario localized gauge fields are a physical necessity, and, regardless of the underlying physics, gauge fields are required to be present in the effective world volume field theory. The invariance of the world volume action under target space and gauge symmetries then dictates the form of the interaction between the gauge fields and the Nambu-Goldstone fields.

The purpose of this brief report is to extend the construction of the low energy effective action of a non-BPS p=2 brane embedded in a N=1, D=4 target superspace of reference \[11\] to include gauge fields in addition to the Nambu-Goldstone and matter degrees of freedom. This is done in a model independent way through the method of non-linear realizations of the spontaneously broken target space super-Poincaré symmetries. Actions for non-BPS D-branes have been constructed in reference \[12\]. There the Green-Schwarz action \[13\], excluding the Wess-Zumino term, was utilized to find the action for a Dp-brane embedded in D=10 superspace. The relation of the generalized Green-Schwarz method of \[12\] to the coset method employed here, and in reference \[11\], remains to be elucidated. However, in the BPS superparticle case, the relation between the Green-Schwarz action and that obtained using nonlinear realization techniques has been clarified in references \[14\]. Returning to the non-BPS
p=2 brane at hand, the Poincaré symmetries of the D=3 world volume space-time of the brane, which has static gauge coordinates $x^m, m = 0, 1, 2$, are realized linearly. The broken translation symmetry Nambu-Goldstone boson is denoted by $\phi(x)$ and the broken SUSY Goldstino D=3 (real) Majorana fields are denoted by $\theta_i(x)$ and $\lambda_i(x), i = 1, 2$. In addition there is an auxiliary Nambu-Goldstone D=3 vector field, $v^m(x)$, associated with the broken D=4 Lorentz transformations. It can be expressed in terms of the independent degrees of freedom $\phi, \theta$ and $\lambda$ through the “inverse Higgs mechanism” [15].

The non-linear realization of the N=1, D=4 super-Poincaré symmetry group $G$, with its broken space translation symmetry and the broken supersymmetries, induces a Nambu-Goldstone field dependent world volume general coordinate transformation

$$dx'^m = dx^n G^m_n$$

where [11]

$$G^m_n = \frac{\partial x'^m}{\partial x^n} = \delta^m_n - i(\partial_n \theta^0 \gamma^m \xi + \partial_n \lambda^0 \gamma^m \eta) - \partial_n \phi b^m + \epsilon_n^{\,ms} \alpha_s.$$  (2)

Concurrent with this are the non-linear transformations of the Nambu-Goldstone fields

$$\Delta \phi = z + (\xi^0 \lambda - \theta^0 \eta) - b^m x_m$$

$$\Delta \theta_i = \xi_i - i b^m (\gamma^m \lambda)_i - i \rho \lambda_i - i b^m (\gamma^m \theta)_i$$

$$\Delta \lambda_i = \eta_i + i b^m (\gamma^m \theta)_i + i \rho \theta_i - i \alpha_m (\gamma^m \lambda)_i$$

$$\Delta v^m = + \frac{\sqrt{v^2}}{\tanh \sqrt{v^2}} \left( b^m - \frac{v^r b_r v^m}{v^2} \right) + \frac{v^r b_r v^m}{v^2} + \epsilon^{mnr} \alpha_n v_r.$$  (3)

where the set of transformation parameters \{z, a^m, (\xi_i, \eta_i), b^m, \alpha^m, \rho\} correspond to the \{broken D=4 translation, unbroken D=3 space-time translations, (broken SUSY, broken SUSY), broken D=4 Lorentz rotations, unbroken D=3 Lorentz rotations, unbroken $R$\} transformations of the N=1, D=4 super-Poincaré symmetries. (Since the unbroken symmetries are those of the D=3 Poincaré group, the D=4 transformation parameters and the Nambu-Goldstone fields are expressed in terms of their D=3 Lorentz group transformation properties.) The intrinsic variation of the fields,
\[ \delta \varphi \equiv \varphi'(x) - \varphi(x), \] is related to the above total variation, \( \Delta \varphi = \varphi'(x') - \varphi(x), \) by the Taylor expansion shift in the space-time coordinates: \( \delta \varphi = \Delta \varphi - \delta x^m \partial_m \varphi, \) with \( \delta x^m = x'^m - x^m = a^m - i(\bar{\xi} \gamma^m \theta + \bar{\eta} \gamma^m \lambda) - \phi b^m + \epsilon^{mnr} \alpha_n x_r. \)

According to the coset construction method [16, 17], the dreibein, the covariant derivatives of the Nambu-Goldstone fields and the spin connection can be obtained from the Maurer-Cartan one-forms. The covariant world volume coordinate differentials, \( \omega^a, \) are found in terms of the dreibein, \( e^a_m, \) and the world-volume coordinate differentials, \( dx^m, \)

\[ \omega^a = dx^m e^a_m, \quad (4) \]

where the dreibein has the factorized form \( e^a_m = \hat{e}^b_m N^a_b \ [11]. \) The Akulov-Volkov dreibein \( \hat{e}^a_m \)

\[ \hat{e}^a_m = A^a_m = \delta^a_m + i \partial_m \theta \gamma^0 \gamma^a \theta + i \partial_m \lambda \gamma^0 \gamma^a \lambda, \quad (5) \]

while the Nambu-Goto dreibein \( N^b_a \) has the form

\[ N^b_a = \delta^b_a + \left[ \cosh \sqrt{v^2} - 1 \right] \frac{v_a v^b}{v^2} + \left( \hat{D}_a \phi + \hat{D}_a \theta \gamma^0 \lambda - \theta \gamma^0 \hat{D}_a \lambda \right) v^b \frac{\sinh \sqrt{v^2}}{\sqrt{v^2}}, \quad (6) \]

where \( \hat{D}_a = \hat{e}^{-1}_a \partial_m = A^{-1}_a \partial_m \) is the Akulov-Volkov partial (SO(1,2)) covariant derivative. In addition the Maurer-Cartan one-form defines the covariant derivatives of the Nambu-Goldstone fields. In particular, the fully (SO(1,3)) covariant derivative, \( \nabla_a \phi, \) of the translational Nambu-Goldstone boson \( \phi \) can be expanded in terms of the Akulov-Volkov partially covariant coordinate differential basis one-forms, \( \hat{\omega}^a = dx^m \hat{e}^a_m, \)

\[ N^b_a \nabla_a \phi = \cosh \sqrt{v^2} \left[ \left( \hat{D}_a \phi + \hat{D}_a \theta \gamma^0 \lambda - \theta \gamma^0 \hat{D}_a \lambda \right) + v_a \frac{\tanh \sqrt{v^2}}{\sqrt{v^2}} \right]. \quad (7) \]

Setting \( \nabla_a \phi = 0 \) results in the “inverse Higgs mechanism” [15], allowing \( v^a \) to be expressed in terms of the independent Nambu-Goldstone degrees of freedom, \( \phi, \theta \) and \( \lambda, \)

\[ v_a \frac{\tanh \sqrt{v^2}}{\sqrt{v^2}} = - \left( \hat{D}_a \phi + \hat{D}_a \theta \gamma^0 \lambda - \theta \gamma^0 \hat{D}_a \lambda \right). \quad (8) \]

The covariant world volume coordinate differentials \( \omega^a \) transform under the N=1, D=4 super-Poincaré group \( G \) as structure group vectors

\[ \omega^b = \omega^a L^b_a, \quad (9) \]
where

\[ L^b_a = (e^{-i\beta^c \tilde{M}_c})^b_a, \] (10)

with the D=3 Lorentz vector representation matrix \((\tilde{M}_c)^b_a = i\epsilon_{ca} b.\) The field dependent transformation parameter \(\beta^c\) is determined from the non-linear realization of \(G\) on the coset element \(\Omega \in G/H\) through \(g\Omega = \Omega'h\), with the group element \(g \in G\) corresponding to the transformation, \(\Omega' \in G/H\) the transformed coset element, and \(h \in H = SO(1, 2) \otimes R\), the unbroken D=3 Lorentz and \(U_R(1)\) symmetry groups. For infinitesimal D=4 Lorentz transformations, with D=3 parameters \(\alpha^m\) for the unbroken transformations and \(b^m\) for the broken ones, the induced local Lorentz transformation parameter was determined in [11] to be

\[ \beta^a(g, v) = \alpha^a - \frac{1}{2} \tanh \frac{1}{2} \sqrt{v^2} \frac{1}{2} \sqrt{v^2} \epsilon^{abc} b_v. \] (11)

Correspondingly, under a \(G\)-transformation the dreibein, \(e^a_m\), transforms with one world index and one tangent space (structure group) index as

\[ e'^a_m = G^{-1n}_m e^n_b L^b_a, \] (12)

and likewise for the inverse dreibein

\[ e'^{-1m}_a = L^{-1b}_a e^b_n G^n_m. \] (13)

Utilizing the flat tangent space metric, \(\eta_{ab}\), the induced world volume metric tensor is given in terms of the dreibein as

\[ g_{mn} = e^a_m \eta_{ab} e^b_n. \] (14)

Consequently the invariant interval can be expressed as \(ds^2 = dx^m g_{mn} dx^n\). The leading term in the D=4 super-Poincaré invariant action is given by integrating the constant brane tension, \(\sigma\), over the area of the brane

\[ \Gamma = -\sigma \int d^3x \det e. \] (15)

The matter fields localized on the brane are characterized by their D=3 Lorentz group (with generators \(M^a\)) transformation properties. A scalar field, \(S(x)\), is in the
trivial representation of the Lorentz group: $M^a \to (\tilde{M}^a) = 0$. Fermion fields, $\psi_i(x)$, are in the spinor representation: $M^a \to (\tilde{M}^a)_{ij} = -1/2\gamma^a_{ij}$. Each matter field, $M(x)$, transforms under $G$ as

$$M'(x') \equiv \tilde{h}M(x),$$

(16)

where $\tilde{h}$ is given by

$$\tilde{h} = e^{i\beta_a(g,v)\tilde{M}^a},$$

(17)

while the $R$-transformation properties have been suppressed. The covariant derivative for the matter field is defined using the Maurer-Cartan spin connection one-form $\omega^a_M = (\cosh \sqrt{v^2} - 1)\epsilon^{abc}v_d\omega^c_{ab}$ as

$$\nabla M \equiv (d + i\omega^a_M \tilde{M}_a)M.$$  

(18)

It has the same transformation properties as the matter field itself,

$$\nabla M' (x') = \tilde{h}\nabla M(x).$$

(19)

The spin connection one-form, $\omega^a_M$, can be expanded in terms of the covariant coordinate differential basis, $\omega^a, \omega^a_M = \omega^a \Gamma^b_a (= dx^m \Gamma^b_m)$ with components

$$\Gamma^b_a = \left(\cosh \sqrt{v^2} - 1\right)\epsilon^{bcd}v_e\omega^e_{bc},$$

(20)

where $\mathcal{D}_a = e^{-1m}\partial_m$. As well, expanding the covariant derivative one-form in terms of $\omega^a$, the component form of the covariant derivative is obtained

$$\nabla_a M = \left(\mathcal{D}_a + i\Gamma^b_{a} \tilde{M}_b\right)M.$$  

(21)

The component form of the covariant derivative has the $G$ transformation law

$$\nabla_a M' (x') = \tilde{h}L^{-1b}_{a} \nabla_b M(x).$$

(22)

The definition of the covariant derivative must be extended when the matter fields also belong to representations of a local internal symmetry group $G$

$$M'^\alpha(x) = (U(\epsilon))^\alpha_\beta M^\beta(x),$$

(23)
where the representation matrix

$$(U(\epsilon))^{\alpha}_{\beta} = (e^{i g A(x) T^A})^{\alpha}_{\beta}, \quad (24)$$

is given in terms of the world volume local transformation parameters $e^A(x)$, the
gauge coupling constant $g$ and the representation generator matrices $(T^A)^{\alpha}_{\beta}$, $A = 1, 2, \ldots, \text{dim}[G]$. The generators obey the associated Lie algebra $[T^A, T^B] = i f^{ABC} T^C$, and are normalized so that $\text{Tr}[T^A T^B] = 1/2 \delta^{AB}$. In order to extend the invariance of the action to include gauge transformations, the world volume Yang-Mills gauge potential one-form, $A(x)$, must be introduced, where

$$A(x) = dx^m A_m(x) = dx^m (i T^A A^A_m(x)). \quad (25)$$

Under $G$-transformations the Yang-Mills one-form is invariant: $A'(x') = A(x)$, that
is the world index gauge field transforms as a coordinate differential

$$A'_m(x') = G^{-1}_{m n} A_n(x). \quad (26)$$

Under $G$-transformations the Yang-Mills field transforms as a gauge connection

$$A = U(\epsilon) A U^{-1}(\epsilon) + \frac{1}{g} (d U(\epsilon)) U^{-1}(\epsilon). \quad (27)$$

Hence, the gauge and super-Poincaré covariant derivative of the matter field is ob-
tained

$$\nabla M = [d + i \omega^a_M \tilde{M}_M a - g A] M. \quad (28)$$

Under super-Poincaré transformations the covariant derivative transforms as does $M$,

$$(\nabla M)'(x') = \tilde{h}(\nabla M)(x). \quad (29)$$

Likewise, under gauge transformations the covariant derivative remains in the same
matter field representation of $G$

$$(\nabla M)' = U(\epsilon)(\nabla M). \quad (30)$$
The matter field derivative can be expanded in terms of the world volume coordinate differentials $dx^m$ as
\[
\nabla_m M = \left( \partial_m + i e_m^a \Gamma_a^b \tilde{M}^b_m - g A_m \right) M. \quad (31)
\]
The fully covariant derivatives for the scalar, $S(x)$, and fermion, $\psi_i(x)$, matter fields have the explicit form
\[
(\nabla_m S)^\alpha = \partial_m S^\alpha - ig A_m^A (T^A)^\alpha_\beta S^\beta,
\]
\[
(\nabla_m \psi_i)^\alpha = \partial_m \psi_i^\alpha - \frac{i}{2} \Gamma_m^a (\gamma_a)_{ij} \psi_j^\alpha - ig A_m^A (T^A)^\alpha_\beta \psi_i^\beta. \quad (32)
\]
Employing the world volume induced metric, $g^{mn}$, and dreibein, $e_a^{-1m}$, the matter field fully invariant action is given by
\[
\Gamma_{\text{matter}} = \int d^3 x \det e \mathcal{L}_{\text{matter}}, \quad (33)
\]
where the fully invariant matter field Lagrangian $\mathcal{L}_{\text{matter}}$, which takes the form
\[
\mathcal{L}_{\text{matter}} = \text{Tr} \left[ (\nabla_m S)^\dagger g^{mn} (\nabla_n S) \right] - V(S) + i \bar{\psi} \gamma^a e_a^{-1m} \nabla_m \psi - \bar{\psi} m \psi + Y(S, \bar{\psi} \psi), \quad (34)
\]
includes a globally $G$-invariant scalar field potential $V(S)$, and (if possible to form) globally $G$-invariant fermion mass terms $\bar{\psi} m \psi$, and generalized Yukawa couplings $Y(S, \bar{\psi} \psi)$.

The Yang-Mills field strength two-form, $F$, is defined as
\[
F \equiv dA + gA \wedge A. \quad (35)
\]
As a two-form, $F$ is invariant under $G$-transformations while under $G$-transformations it is in the adjoint representation
\[
F' = U(\epsilon) F U^{-1}(\epsilon). \quad (36)
\]
Expanding $F$ in terms of the coordinate differential basis $dx^m$, $F = \frac{1}{2} dx^m \wedge dx^n (iT^A F^A_{mn})$, the world index field strength tensor is obtained
\[
F^A_{mn} = \partial_m A^A_n - \partial_n A^A_m + gf^{ABC} A^B_m A^C_n. \quad (37)
\]
The fully invariant Yang-Mills action (constructed in \[18\] when only supersymmetry is realized nonlinearly) is secured as

\[ \Gamma_{Y-M} = \int d^3x \det e \mathcal{L}_{Y-M}, \]  

with the invariant Lagrangian \( \mathcal{L}_{Y-M} \) given by

\[ \mathcal{L}_{Y-M} = -\frac{1}{2} \text{Tr}[F_{mn} g^{mr} g^{ns} F_{rs}], \]

There are two other useful one-form bases in which to express the derivatives and gauge fields. The basis consisting of the fully covariant world volume coordinate differentials \( \omega^a = dx^m e_m^a \) and the basis consisting of the partially (SO(1,2)) covariant Akulov-Volkov one-forms \( \hat{\omega}^a \) with \( \omega^a = \hat{\omega}^b N_b^a \). The exterior derivative can be expanded in these bases as 

\[ d = dx^m \partial_m = \omega^a D_a = \hat{\omega}^a \hat{D}_a \]

where the derivatives are related by the Nambu-Goto dreibein, \( D_a = N_a^{-1b} \hat{D}_b \). Likewise, the gauge field one-form has the expansion

\[ A = dx^m A_m = \omega^a A_a = \hat{\omega}^a \hat{A}_a. \]

As previously noted, the fully covariant basis \( \omega^a \) transforms according to the vector representation of the D=3 (local) Lorentz structure group, \( \omega'{}^a = \omega^b L_b^a \) with the Nambu-Goldstone field dependent matrix \( L_b^a \) given in equations (10) and (11). So the fully covariant gauge field transforms analogously, \( A'_a(x') = L_a^{-1b} A_b(x) \). The partially covariant basis differentials \( \hat{\omega}^a \) transform SO(1,2) covariantly but not SO(1,3) covariantly,

\[ \hat{\omega}'{}^a = \hat{\omega}^b \hat{L}_b^a \]

\[ = \hat{\omega}^b \left( \delta_b^a + \alpha^c \epsilon_{cb}^{} \right) \frac{\tanh \sqrt{v^2}}{\sqrt{v^2}}, \]

with \( v^a \) related to the Nambu-Goldstone fields via the “inverse Higgs” mechanism, equation (8), and likewise for the partially covariant gauge fields

\[ \hat{A}'_a(x') = \hat{L}_a^{-1b} \hat{A}_b(x). \]

In similar fashion, the fully covariant derivative transforms as \( D'_a = L_a^{-1b} D_b \) while the SO(1,2) partially covariant Akulov-Volkov derivative transforms as \( \hat{D}'_a = \hat{L}_a^{-1b} \hat{D}_b \).
The invariant interval can be expressed in each basis by use of a metric specific to it
\[ ds^2 = dx^m g_{mn} dx^n = \omega^a n_{ab} \omega^b, \] (43)
with \( \eta_{ab} \) the flat tangent space Minkowski metric and \( n_{ab} \) the Nambu-Goto metric made from the Nambu-Goto dreibein \( N_a^b \), \( n_{ab} = N_a^c \eta_{cd} N_b^d = \eta_{ab} - \frac{n_{ab}}{v^2} \tanh^2 \sqrt{v^2} \).

The matter field covariant derivatives can be expressed in each basis as
\[
(\nabla_a S)^\alpha &= D_a S^\alpha - ig A_a^A (T^A)^\alpha_\beta S^\beta \\
(\hat{\nabla}_a S)^\alpha &= \hat{D}_a S^\alpha - ig \hat{A}_a^A (T^A)^\alpha_\beta S^\beta \\
(\nabla_a \psi)_{i}^\alpha &= D_a \psi_{i}^\alpha - \frac{i}{2} \Gamma_a^{bc} \gamma_{ij} \psi_{j}^\alpha - ig A_a^A (T^A)^\alpha_\beta \psi_{j}^\beta \\
(\hat{\nabla}_a \psi)_{i}^\alpha &= \hat{D}_a \psi_{i}^\alpha - \frac{i}{2} \hat{\Gamma}_a^{bc} \gamma_{ij} \psi_{j}^\alpha - ig \hat{A}_a^A (T^A)^\alpha_\beta \psi_{j}^\beta, \]
(44)
with \( \hat{\Gamma}_a^{bc} = N_a^c \Gamma_c^b \). With these replacements, the fully invariant matter field Lagrangian, equation (34), can be written in the two bases as
\[
\mathcal{L}_{\text{matter}} &= \text{Tr} \left[ (\nabla_a S) \eta^{ab} (\nabla_b S) \right] - V(S) \\
&\quad + i \bar{\psi} \gamma^a \nabla_a \psi - \bar{\psi} m \psi + Y(S, \bar{\psi} \psi) \\
\mathcal{L}_{\text{matter}} &= \text{Tr} \left[ (\hat{\nabla}_a S) \eta^{ab} (\hat{\nabla}_b S) \right] - V(S) \\
&\quad + i \bar{\psi} \hat{\gamma}^a \nabla_a \psi - \bar{\psi} m \psi + Y(S, \bar{\psi} \psi), \]
(45)
with the Dirac matrices in the partially covariant basis defined by means of the Nambu-Goto dreibein \( \gamma^a = \gamma^b N_b^{-1a} \).

The Yang-Mills fields expressed in terms of the different bases have correspondingly modified gauge variations
\[
A_a' &= U(\epsilon) A_a U^{-1}(\epsilon) + \frac{1}{g} (D_a U(\epsilon)) U^{-1}(\epsilon) \\
\hat{A}_a' &= U(\epsilon) \hat{A}_a U^{-1}(\epsilon) + \frac{1}{g} (\hat{D}_a U(\epsilon)) U^{-1}(\epsilon). \]
(46)
Also, the field strength tensor takes the various forms
\[
F_{ab}^A &= D_a A_b^A - D_b A_a^A + gf^{ABC} A_b^B A_a^C \\
\hat{F}_{ab}^A &= \hat{D}_a \hat{A}_b^A - \hat{D}_b \hat{A}_a^A + g f^{ABC} \hat{A}_b^B \hat{A}_a^C. \]
(47)
As with the matter field Lagrangian, the fully invariant Yang-Mills Lagrangian, equation (39), in these bases becomes

\[
\mathcal{L}_{Y-M} = -\frac{1}{2} \text{Tr}[F_{ab} \eta^{ac} \eta^{bd} F_{cd}]
\]

\[
\mathcal{L}_{Y-M} = -\frac{1}{2} \text{Tr}[\hat{F}_{ab} \hat{n}^{ac} \eta^{bd} \hat{F}_{cd}].
\] (48)

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References

[1] A. M. Polyakov, JETP Lett. 20, 194 (1974) [Pisma Zh. Eksp. Teor. Fiz. 20, 430 (1974)].

[2] M. B. Voloshin, Sov. J. Nucl. Phys. 21, 687 (1975) [Yad. Fiz. 21, 1331 (1975)].

[3] R. Jackiw and C. Rebbi, Phys. Rev. D 13, 3398 (1976).

[4] V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B 125 (1983) 136.

[5] K. Akama, Lect. Notes Phys. 176, 267 (1982) [arXiv:hep-th/0001113].

[6] S. L. Dubovsky and V. A. Rubakov, Int. J. Mod. Phys. A 16, 4331 (2001) [arXiv:hep-th/0105243];

[7] G. R. Dvali and M. A. Shifman, Phys. Lett. B 396, 64 (1997) [Erratum-ibid. B 407, 452 (1997)] [arXiv:hep-th/9612128].

[8] G. R. Dvali, G. Gabadadze and M. A. Shifman, Phys. Lett. B 497, 271 (2001) [arXiv:hep-th/0101071].

[9] E. K. Akhmedov, Phys. Lett. B 521, 79 (2001) [arXiv:hep-th/0107223].

[10] I. Oda, Phys. Lett. B 496, 113 (2000) [arXiv:hep-th/0006203]; S. L. Dubovsky, V. A. Rubakov and P. G. Tinyakov, JHEP 0008, 041 (2000) [arXiv:hep-ph/0007179]; T. Kugo and K. Yoshioka, Nucl. Phys. B 594, 301 (2001) [arXiv:hep-ph/9912496].

[11] T. E. Clark, M. Nitta and T. ter Veldhuis, arXiv:hep-th/0208184.

[12] A. Sen, JHEP 9910, 008 (1999) [arXiv:hep-th/9909062].

[13] M. B. Green and J. H. Schwarz, Phys. Lett. B 136, 367 (1984); M. B. Green and J. H. Schwarz, Nucl. Phys. B 243, 285 (1984); M. Aganagic, C. Popescu and J. H. Schwarz, Nucl. Phys. B 495, 99 (1997) [arXiv:hep-th/9612080].
[14] E. Ivanov and S. Krivonos, Phys. Lett. B 453, 237 (1999) [arXiv:hep-th/9901003]; F. Delduc, E. Ivanov and S. Krivonos, Nucl. Phys. B 576, 196 (2000) [arXiv:hep-th/9912222]; F. Delduc, E. Ivanov and S. Krivonos, Talk given at 14th Max Born Symposium: New Symmetries and Integrable Systems, Karpacz, Poland, 21-24 Sep 1999. Published in “Karpacz 1999, New symmetries and integrable models,” p. 131 [arXiv:hep-th/9912292].

[15] E. A. Ivanov and V. I. Ogievetsky, Teor. Mat. Fiz. 25, 164 (1975).

[16] S. R. Coleman, J. Wess and B. Zumino, Phys. Rev. 177, 2239 (1969); C. G. Callan, S. R. Coleman, J. Wess and B. Zumino, Phys. Rev. 177, 2247 (1969).

[17] D. V. Volkov, Sov. J. Particles and Nuclei 4, 3 (1973); V. I. Ogievetsky, Proceedings of the X-th Winter School of Theoretical Physics in Karpacz, vol. 1, p. 227 (Wroclaw, 1974).

[18] T. E. Clark, T. Lee, S. T. Love and G. H. Wu, Phys. Rev. D 57, 5912 (1998) [arXiv:hep-ph/9712353].