Covariant $\kappa$–Symmetry Gauge Fixing and the Classical Relation Between Physical Variables of the NSR String and the Type II GS Superstring

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The main goal of this paper is the manifestly covariant derivation of the classical relation between the gauge-fixed Grassmann variables of the NSR string and the Type II GS superstring. To this end we analyze the superembedding equation for the Type II superstring to derive the relation between the original variables of the NSR string and the Type II GS superstring and, further, by means of Lorentz harmonic variables we fix $\kappa$–symmetry of the GS superstring in the manifestly Lorentz covariant way.

1. INTRODUCTION

There exist two ways of unification of supersymmetry and string theory. One is able either to introduce local supersymmetry on the string worldsheet resulting in the Neveu-Schwarz-Ramond (NSR) model \cite{1}, which can be regarded as the coupling of 2d induced supergravity with matter, or global supersymmetry in a target-space which yields the Green-Schwarz (GS) superstring \cite{2}. The introduction of local supersymmetry in the target-space leads to a superstring model interacting with target-space supergravity. Models with both worldsheet (local) and target-space (global or local) supersymmetry, the so-called spinning superstrings \cite{3}, \cite{4}, in some sense can be regarded as a product of NSR and GS models. They possess wider spectrum of quantum states in comparison with the NSR and GS models.

The NSR string and the GS superstring are in fact different string theories at the classical level. It is remarkable, however, that upon first quantization they describe the same set of quantum states provided the NSR string is subject to GSO projection \cite{5}, \cite{6}. Therefore it seems interesting to establish a direct classical relation between these models at the level of variables and actions and to find a classical analogue of the GSO projection. This problem was put forward in late 80-s by D.V. Volkov and his collaborators and the solution was found for the case of superparticle and spinning particle models \cite{7}, \cite{8}, \cite{9}. Since then interesting results on the relation of the NSR and the GS superstring have been obtained in \cite{10}, \cite{11} and \cite{12} which we shall generalize in this contribution.

In the present paper following the guidelines of pioneer works \cite{7}, \cite{8}, \cite{9} and using the methods that were developed afterwards (the superembedding approach (see review \cite{13} and references therein) and Lorentz harmonic variables \cite{14} adapted to the description of strings \cite{15}, \cite{16}) we deduce the relation between the physical variables (i.e. those surviving local symmetry gauge fixing) of the NSR string and the Type II GS superstring. This requires analysis of the superembedding equation for the $n = (1|1)$ worldsheet superspace embedded into the flat $D = 10$ Type II target-superspace to obtain the relation between the original Grassmann variables of the models which include not only physical but also pure gauge ones (Section 2). $SO(1,9)$ Lorentz covariant $\kappa$–symmetry gauge fixing of the twistor-like Lorentz harmonic formulation of the Type II GS superstring \cite{12} is the subject of Section 3. Finally in Section 4 we compile all these data to find the covariant relation between the physical variables of the NSR string and the Type II GS

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superstring.

2. ANALYSIS OF THE SUPEREMBEDDING EQUATION FOR \( n = (1|1) \) WORLDSHET SUPERSPACE EMBEDDED INTO \( D = 10 \) TYPE II TARGET-SUPERSPACE

Let us consider the embedding of \( n = (1|1) \) superworldsheet into flat \( D = 10 \) Type II target-superspace. \( n = (1|1) \) worldsheets super spac is locally parametrized by the coordinates \( z^M = (\xi^m, \eta^\alpha) \), where Grassmann coordinate \( \eta^\alpha \) is the 2d Majorana spinor that consists of two one-component Majorana-Weyl spinors of opposite chiralities (that explains the notation \( n = (1|1) \)). The choice of \( n = (1|1) \) worldsheets super spac is dictated by the fact that in such worldsheets superspace the NSR string is formulated. \( D = 10 \) Type II target-superspace is parametrized by coordinates \( Z^M = (X^\mu, \theta^\alpha, \theta^{\dot{\alpha}}) \) in the Type IIA case and by \( Z^M = (X^\mu, \theta^\alpha, \theta^{\dot{\alpha}}) \) in the Type IIB case. From the superembedding approach it follows that in order to describe conventional string (brane) models one should impose certain restrictions on the embedding, namely, the superembedding equation (6) which reads that the pull-back of the target-space supervielbein bosonic components along the superworldsheet Grassmann directions vanishes

\[
D_a Z^M E_M^a (Z(z)) = 0, \tag{1}
\]

where \( E_M^a \) and \( E_M^\dot{a} \) are the ordinary GS variables, \( \epsilon_M^a \) are the target-space supervielbein and the worldsheet superzweinbein respectively. Their inverse are \( E^M_a = (E_M^a, E_M^{\dot{a}}, E_2^M) (E^2_{M}) \), \( E^M_B E_M^a = \delta_B^a \) and \( e^M_B = (e^M_1, e^M_2) \), \( e^M_B e_B^M = \delta_B^A \). So that the worldsheet Grassmann covariant derivative entering (5) looks like \( D_a = e^a_M \partial_M \). Using the explicit expression for the bosonic components of the flat target-space supervielbein

\[
(IIA) \quad \Pi^a = dZ^M E_M^a = \delta_M^a (dX^m - 2 \rho^m A \Theta^2_{\dot{\alpha}}), \tag{2}
\]

\[
(IIB) \quad \Pi^2 = dZ^M E_M^a = \delta_M^a (dX^m - 2 \rho^m A \Theta^2_{\dot{\alpha}}) - i \theta^{\alpha_1} e^{\alpha_2 \dot{a}} \Theta^1_{\alpha_2} - i \theta^{\dot{\alpha_1}} e^{\dot{\alpha_2} a} \Theta^2_{\dot{\alpha_2}} \tag{3}
\]

and splitting the tangent space indices of the superworldsheet in light-like basis \( e_M^a = (e_M^\pm, e_M^\mp) \), \( e_M^a = (e_\pm M, e_{\mp}^M) \) one obtains from (5)

\[
D^{\pm} X^m - i D_{\pm} \Theta^{1 \alpha} e^{\alpha \dot{a}} \Theta^{1 \dot{a}} - i D_{\pm} \Theta^{2 \alpha} e^{\alpha a} \Theta^{2 a} = 0 \tag{4}
\]

for the Type IIA case and

\[
D^{\pm} X^m - i D_{\pm} \Theta^{1 \alpha} e^{\alpha \dot{a}} \Theta^{1 \dot{a}} - i D_{\pm} \Theta^{2 \alpha} e^{\alpha a} \Theta^{2 a} = 0 \tag{5}
\]

for the Type IIB case. \( n = (1|1) \) supergravity in \( 2d \) is well studied and the basic fact is that it is superconformally flat. Thus, the component expansion of the worldsheetsuperfields \( X^m (z^M) \), \( \Theta^{1 \alpha} (z^M) \) acquire the most simple form

\[
X^m (\xi^m, \eta^\alpha) = x^m (\xi^m) + \frac{i}{\sqrt{8}} \eta^\alpha \psi^m (\xi^m), \tag{6}
\]

\[
\Theta^{1 \alpha} (\xi^m, \eta^\alpha) = \theta^{1 \alpha} (\xi^m) + \eta^\alpha \lambda^{1 \alpha} (\xi^m) + i \eta^\alpha \eta^\beta \rho^{1 \alpha \beta} (\xi^m), \tag{7}
\]

where \( x^m \) and \( \theta^{1 \alpha} \) are the ordinary GS variables, \( \psi^m \) are the NSR Grassmann variables, \( \lambda^{1 \alpha} \) are the stringy twistor-like variables, and \( F^m \) and \( \rho^{1 \alpha \beta} \) are auxiliary fields. For further analysis we impose the chirality conditions on \( \Theta^{1 \alpha} \) superfields

\[
D_- \Theta^{1 \alpha} = D_+ \Theta^{2 \alpha} = 0, \tag{8}
\]

which on the component level are equivalent to

\[
\lambda^{1 \alpha} = \rho^{1 \alpha} = 0, \quad \lambda^{2 \alpha} = \rho^{2 \alpha} = 0; \tag{9}
\]

\[
\partial_{\alpha_1} \theta^{1 \alpha} = \partial_{\alpha_1} \lambda^{1 \alpha} = \partial_{\alpha_1} \lambda^{2 \alpha} = 0, \tag{10}
\]

\[
\partial_{\alpha_2} \theta^{2 \alpha} = \partial_{\alpha_2} \lambda^{1 \alpha} = \partial_{\alpha_2} \lambda^{2 \alpha} = 0. \tag{11}
\]
The conditions \[ \{ \] contain equations of motion and thus put the variables on the mass shell. When eq. \[ \{ \] is imposed the superembedding equation yields
\[
\psi_+^m = \sqrt{8} \lambda_+^m \sigma^m \partial^2 \bar{\beta}^2, \quad \psi_-^m = \sqrt{8} \lambda_-^m \sigma^m \partial^1 \bar{\beta}; \tag{12}
\]
\[
\Pi^m_{\pm 2} = \lambda_+ \sigma^m \lambda_+; \tag{13}
\]
\[
F^m = 0. \tag{14}
\]
For the Type IIA case we have \( \psi_+^m = \sqrt{8} \lambda_+^m \sigma^m \partial^2 \bar{\beta}^2, \quad \Pi^m_{\pm 2} = \lambda_- \sigma^m \lambda_- \). Applying then equations \[ \{ \] and \[ \{ \] one gets the NSR string fermionic equations of motion
\[
\partial_{+2} \psi_+^m = 0, \quad \partial_{-2} \psi_-^m = 0. \tag{15}
\]

Let us outline some properties of the obtained formulae. Representation \[ \{ \] connects in a natural way the Grassmann variables of the NSR string and the Type II GS superstrings. Its main drawback is the mismatch between the number of degrees of freedom on the l.h.s and the r.h.s. The balance is, however, restored on the constraint shell. Indeed, the NSR string Grassmann vectors \( \psi_+^m \) contain 9+9 components as a result of two supercurvilinear constraints. However, as will be seen below, a proper solution of these constraints ensures dropping away the two extra components of \( \psi_+^m \) from the NSR string action. So, only 8+8 physical components of \( \psi_+^m \) contribute to the action. On the other hand, among \( 16 + 16 \) components of the two 10D MW spinors \( \theta^{1,2\alpha} \) there remain 8+8 components after explicit gauge fixing \( \kappa \)-symmetry (see Sec.3). Thus, on the constraint shell there is the same number of the Grassmann degrees of freedom in both formulations of string theory, as it should be. In Section 4 we will find manifest expressions for the physical variables in the NSR and GS model and relate them to each other. Representation \[ \{ \] solves the Type II GS superstring Virasoro constraints since the vectors \( \lambda_+ \sigma^m \lambda_+ \) and \( \lambda_- \sigma^m \lambda_- \) are light-like due to the famous 10D permutation relation
\[
\sigma^m \sigma^m \lambda_+ \lambda_+ + \sigma^m \sigma^m \lambda_- \lambda_- = 0. \tag{16}
\]
The NSR string and GS superstring equations of motion are satisfied by virtue of \[ \{ \] and \[ \{ \].

### 3. Twistor-like Lorentz Harmonic Formulation of Type II GS Superstrings and Covariant \( \kappa \)-Symmetry Fixing

The twistor-like Lorentz harmonic formulation of the Type IIB GS superstring, which is classically equivalent to the original formulation was constructed in \[ \{ \]:
\[
S = \int e \left( - (\alpha')^{-1/2} \epsilon_m \nu_m \Pi^m_{\pm 2} + c \right) \]
\[
- \frac{1}{c\alpha'} \int e^{m n} \Pi^m_{\pm 2} \left( \partial_n \theta^1 \sigma_m \theta^1 - \partial_n \theta^2 \sigma_m \theta^2 \right) \]
\[
+ \frac{1}{c\alpha'} \int e^{m n} \partial_n \theta^1 \sigma_m \theta^1 \partial_n \theta^2 \sigma_m \theta^2. \tag{17}
\]
In the Type IIA case one should replace \( \partial_n \theta^2 \sigma_m \theta^2 \) with \( \partial_n \theta^2 \sigma_m \theta^2 \). In addition to the variables present in the standard GS superstring formulation \[ \{ \] it contains the worldsheet zweinbein \( e_a^m \) and the light-like Lorentz frame vectors \( u^m_{\pm} \), tangent to the string worldsheet. These light-like Lorentz frame vectors together with \( u^m_{\pm} (i = 1, \ldots, 8) \), orthogonal to the worldsheet, constitute a complete orthonormal basis
\[
u_{\pm 2} \cdot u_{\pm 2} = 0, \quad u_{\pm 2} \cdot u_{\pm 2} = 2,
\]
\[
u_{\pm 2} \cdot u_i = 0, \quad u_i \cdot u_i = -\delta^{ij} \tag{18}
\]
which one can use to expand any \( D = 10 \) Minkowski vector. The Lorentz frame vectors can be presented as bilinear combinations of the spinor harmonics \( \nu_{\pm} \) or their inverse \( (\nu^{-1})_{\pm} \) or \( (\nu^{-1})_{\pm} \) or \( (\nu^{-1})_{\pm} \) or \( (\nu^{-1})_{\pm} \) or \( (\nu^{-1})_{\pm} \) or \( (\nu^{-1})_{\pm} \). In addition to the variables present in the standard GS superstring formulation \[ \{ \] it contains the worldsheet zweinbein \( e_a^m \) and the light-like Lorentz frame vectors \( u^m_{\pm} \), tangent to the string worldsheet. These light-like Lorentz frame vectors together with \( u^m_{\pm} (i = 1, \ldots, 8) \), orthogonal to the worldsheet, constitute a complete orthonormal basis
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u_{\pm 2} \cdot u_i = 0, \quad u_i \cdot u_i = -\delta^{ij} \tag{18}
\]
From the embedding theory point of view (19), (20) $\Omega^I_m$ can be identified with the torsion (third fundamental form) components, $\Omega^{+2}_m$ and $\Omega^{-2}_m$ with the second fundamental form components and $\Omega^{(0)}_m$ with the 2d spin connection. Integrability conditions of eqs. (17) are the Gauss, Peterson-Kodacci and Ricci equations (19), (20).

The action (17) is invariant under $\kappa$–symmetry transformations with local parameters $\kappa^+_A$ and $\kappa^-_A$ (see (18)). For the Type IIA case $\kappa$–symmetry transformations read:

$$
\delta \text{I}_A = v^+_{2A} \kappa^+_A, \ \delta \text{II}_A = v^-_{2A} \kappa^-_A,
$$

$$
\delta w^{\alpha} = i \left( \kappa^+_A \bar{w} - \frac{\alpha}{2} \sigma^{\alpha} \theta^+_A \right), \ \delta (ee^{m+2}) = \frac{4i}{c(\alpha')^{1/2}} \kappa^+_{\alpha m} \partial_\alpha \theta^+_A + \Omega^m_{\alpha A},
$$

$$
\delta (ee^{m-2}) = -\frac{4i}{c(\alpha')^{1/2}} \kappa^+_{\alpha m} \partial_\alpha \theta^-_A + \Omega^m_{\alpha A},
$$

$$
\delta u^{\alpha^+}_m = -\frac{2i}{c(\alpha')^{1/2}} \kappa^+_{\alpha^+ m} u^\alpha, \ \delta u^{-}_m = \frac{2i}{c(\alpha')^{1/2}} \kappa^+_{\alpha^+ m} u^\alpha.
$$

(20)

where

$$
w^\alpha_m = -\left( \kappa^+_{B^+ A^+} \partial M \theta^+_A \right) - \kappa^+_{B^+ A^+} \partial M \theta^+_A.
$$

To gauge fix $\kappa$–symmetry it is useful to expand the Grassmann variables in the basis of the spinor harmonics:

$$
\theta^+_{1A} = v^+_{2A} \theta^+_A + v^-_{2A} \theta^-_A,
$$

$$
\theta^-_{1A} = v^+_{2A} \theta^+_A - v^-_{2A} \theta^-_A,
$$

(21)

for the Type IIA case and

$$
\theta^{(2)}_{1A} = v^+_{2A} \theta^{(2)}_A + v^-_{2A} \theta^{(2)}_A
$$

(22)

for the Type IIB case.

Using (21) it is straightforward to show that among the variables introduced in (21) and (22) $\theta^+_A, \theta^-_A$ (IIA) and $\theta^+_A, \theta^-_A$ (IIB) are pure gauge, thus it is admissible to impose the following covariant $\kappa$–symmetry gauge fixing conditions:

$$
\theta^+_A = \theta^-_{A} = 0, \ \theta^{(2)}_A = 0.
$$

(23)

In this gauge the remaining variables $\theta^-_{A} \equiv \theta^-_{A}$, $\theta^+_{A} \equiv \theta^+_{A}$ (IIA) and $\theta^-_{A} \equiv \theta^-_{A}, \theta^+_{A} \equiv \theta^+_{A}$ (IIB) are $\kappa$–invariant. Note, that they are the worldsheet MW spinors, whereas original variables $\theta^{(1,2)}_{A}$ were the worldsheet scalars.

In conclusion let us present $\kappa$–symmetry-fixed version of (17). To this end analogously to the Grassmann variables we expand the bosonic coordinate $x^{\alpha}$ in the basis of the vector harmonics

$$
x^m = \frac{1}{2} x^m + x^m - x^m - x^m.
$$

(25)

Then the derivatives $\partial M x^{\alpha}$ entering the action (17) transform into the covariant derivatives containing the $SO(1,9)$ Cartan forms constructed from the harmonics

$$
D_m x^{\alpha} = \partial_m x^{\alpha} + \Omega_m^{\alpha(1,2)} x^{(1,2)}.
$$

(26)

The covariant derivatives of the fermionic variables are

$$
D_m \theta^+_A = \partial_m \theta^+_A - \frac{1}{2} \Omega_m^{(0)} \theta^+_A - \Omega_m^{ij} \theta^+_A \theta^+_B, \ \ D_m \theta^-_A = \partial_m \theta^-_A + \frac{1}{2} \Omega_m^{(0)} \theta^-_A - \Omega_m^{ij} \theta^-_A \theta^-_B.
$$

(27)

(28)

Then the $\kappa$–symmetry gauge fixed action for the Type II GS superstring takes the form

$$
S_{fixed} = \int d^2 \zeta \left[ \frac{1}{2} (D_m x^m - 2 i D_m \theta^- \theta^+) - (D_m x^m - 2 i D_m \theta^- \theta^+) + c \right]
$$

- \left( D_m x^m - 2 i D_m \theta^- \theta^+ \right) D_n \theta^- \theta^-

(29)

$$
\theta^+_A = \theta^-_A = 0, \ \theta^{(2)}_A = 0.
$$

(IIA)
The form of the action (29) resembles that of Ref. [14] for the superparticles. The action (29) contains the following light-cone-like terms quadratic in $\theta$

$$S_{l.c.} = \frac{2i}{c(\alpha')^{1/2}} \int \sum_{\pm} \left(1 + \frac{1}{c(\alpha')^{1/2}} D_{\pm}^2 x^{\mp 2}\right) D\tau \theta^\mp \theta^\mp$$ (30)

4. RELATION BETWEEN PHYSICAL VARIABLES OF THE NSR STRING AND TYPE II GS SUPERSTRING

First let us establish the connection between the Lorentz harmonic variables and the commuting spinors $\lambda^\pm$ of Section 2. The corresponding relation in the case of $n = (8|8)$ worldsheet superspace was established in [21]. Note that the variation of the action (17) with respect to the zweinbeins and harmonics produces component embedding equation

$$\Pi^m = \frac{c(\alpha')^{1/2}}{2} \left(\epsilon_m - \epsilon_m + \epsilon_m - \epsilon_m\right).$$ (31)

Upon choosing the conformal gauge for the zweinbein $e^a_m = e^{-\phi} \delta^a_m$ equation (31) transforms into

$$\Pi^m = \frac{c(\alpha')^{1/2}}{2} e^{-\phi} \epsilon_m^{\pm 2}.$$ (32)

Equation (32) coincides with (13) only if

$$\lambda^a = \frac{c(\alpha')^{1/2}}{2} e^{-\phi} \epsilon_m^{\pm 2},$$ (33)

which establishes the link with the discussion of the Section 3. To analyze the consequences of (33) let us expand $\lambda^a_\pm$ in the basis of the spinor harmonics

$$\lambda^a_\pm = k \left(v^a_\pm - \lambda_\pm + v^a_\pm \lambda_\pm \right),$$ (34)

where $k = \frac{c(\alpha')^{1/2}}{2} e^{-\phi}/2$. Then one finds that $\lambda_{\pm 2\pm} = \lambda_{\pm 2\mp} = 0$ and $\lambda_\pm \lambda_\mp = 1$, $\lambda_\pm \lambda_\mp = 1$ for the Type IIB case. Corresponding relations for the Type IIA case are $\lambda_{\pm 2\pm} = 0$ and $\lambda_\pm \lambda_\pm = \lambda_\mp \lambda_\mp = 1$.

The connection between $\kappa$–symmetry gauge fixed GS variables and NSR variables can be established upon the substitution of

$$\theta_{\pm} = v^a_\pm \theta^a_\pm, \quad \theta^a = v^a \theta^a,$$

$$\lambda^{a}_\pm = kv^a_\pm \lambda_\pm, \quad \lambda^{a}_\pm = kv^a_\pm \lambda_\pm$$ (35)

into (12) and the expansion of $\psi^m_\pm$ in the basis of the vector harmonics

$$\psi^m_\pm = \frac{1}{2} u^m \varphi^\mp + \frac{1}{2} u^m \varphi^\mp - u^m \varphi^\mp.$$ (36)

As a result we obtain

$$\varphi^\pm = \varphi^\mp = 0,$$

$$\varphi^\pm = -\sqrt{8} k \lambda_\pm \gamma^{\pm} \theta^\pm,$$ (37)

For the Type IIA case we have:

$$\theta_{\pm} = v^a_\pm \theta^a_\pm, \quad \theta^a_\mp = v^a_\pm \theta^a,$$

$$\lambda^{a}_\pm = kv^a_\pm \lambda_\pm, \quad \lambda^{a}_\pm = kv^a_\pm \lambda_\pm,$$ (38)

so

$$\varphi^\pm = \varphi^\mp = 0,$$

$$\varphi^\pm = -\sqrt{8} k \lambda_\pm \gamma^{\pm} \theta^\pm,$$ (39)

5. CONCLUSION

We have considered the superembedding equation for $n = (1|1)$ worldsheet superspace embedded into flat $D = 10$ Type II target-superspace. It was shown to contain the relation (12) between the original variables of the NSR string and the Type II GS superstring, as well as the solution (13) to the Type II GS superstring Virasoro constraints. Upon Lorentz covariant gauge fixing $\kappa$–symmetry with the use of the twistor-like Lorentz harmonic variables, which amounts to co-variantizing the light-cone gauge, eq.(12) reduces to the relation between $\kappa$–symmetry gauge fixed variables $\theta^\pm_\pm$ (Type IIA) and $\theta^\mp_\pm, \theta^\mp_\pm$ (Type IIB) of the GS superstring, and the transverse physical variables $\varphi^\pm$ of the NSR string.

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