Geometry of N=1 Super Yang-Mills Theory in Curved Superspace.

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Abstract
We give a new description of N=1 super Yang-Mills theory in curved superspace. It is based on the induced geometry approach to a curved superspace in which it is viewed as a surface embedded into C^{4|2}. The complex structure on C^{4|2} supplied with a standard volume element induces a special Cauchy-Riemann (SCR)-structure on the embedded surface. We give an explicit construction of SYM theory in terms of intrinsic geometry of the superspace defined by this SCR-structure and a CR-bundle over the superspace. We write a manifestly SCR-covariant Lagrangian for SYM coupled with matter. We also show that in a special gauge our formulation coincides with the standard one which uses Lorentz connections. Some useful auxiliary results about the integration over surfaces in superspace are obtained.

1 Introduction
The main purpose of this work is to apply the induced geometry approach to N=1 supergravity to the construction of N=1 SYM theory on a curved superspace. This approach was introduced in paper [1]. We refer the reader to that paper for all details about the geometric constructions we use in the present one. However we give all necessary definitions and the paper can be read independently. The main benefit of induced geometry approach is that it does not require any additional constraints like those one has to impose on a curvature and torsion tensors in the standard formulation of supergravity (see [2] for a good exposition). The induced SCR-structure on a superspace incorporates all
the necessary constraints and seems to be a more natural geometric construction then the Lorentz connections of the conventional approach. The induced geometry approach to N=1 supergravity was further developed in papers [3], [4], [5]. In paper [6] the application to the construction of SYM theory over curved superspace was proposed. In the present work we develop another point of view on this problem.

The paper is organized as follows. In section 2 we give some auxiliary results about the integration over integral surfaces defined by (0, 2)-dimensional distribution. The derivation of these results is postponed until appendix A. Then we give the description of a curved superspace and CR-structure on it within the framework of induced geometry approach. After that we explain how the results about the integration over integral surfaces can be applied to the construction of a chiral projecting operator. In section 3 we introduce some more geometric notions and define the gauge fields as sections of CR-bundles over the superspace. Then we formulate the main result and explain all the ingredients. The derivation of the Lagrangian is given in appendix B. We finish this section by considering the proper reality conditions imposed on the fields used in the construction. As a result of these restrictions the set of fields reduces to the standard one. This procedure is similar to the one introduced in [7]. In section 4 we compare our constructions to the conventional ones.

2 Integration over integral surfaces determined by (0,2)-dimensional distributions

Let $\mathcal{M}$ be a real $(m, n)$-dimensional supermanifold. Consider a pair of odd vector fields $E_{\alpha}$ (where $\alpha = 1, 2$) which are closed with respect to the anticommutator, i.e.

$$\{E_{\alpha}, E_{\beta}\} = c_{\alpha\beta}^{\gamma}E_{\gamma}$$

(1)

where $c_{\alpha\beta}^{\gamma}$ are some odd functions on $\mathcal{M}$. By a super version of the Frobenius theorem, condition (1) means that the vector fields $E_{\alpha}$ determine an integrable distribution, i.e. for every point in $\mathcal{M}$ there exists a (0, 2)-dimensional surface $\Sigma$ going through it such that its tangent plane at each point coincides with the one spanned by $E_{\alpha}$ (equivalently our distribution defines a foliation which has the surfaces $\Sigma$ as leaves). We are interested in the functions on $\mathcal{M}$ which are stable under the action of the vector fields $E_{\alpha}$, i.e. the functions $\Phi$ such that $E_{\alpha}\Phi = 0$. Here and below we use the same symbol for vector fields and for the first order differential operators corresponding to them.

On every surface $\Sigma$ we have a natural volume element defined by the requirement that the value of the volume form taken on the fields $E_1$ and $E_2$ is equal to one. Given an arbitrary function $\Phi$ one can get a function with the property we want by integrating $\Phi$ over the leaves $\Sigma$. We will denote this operation by $\square = E^\alpha\Phi$ when applied to $E$. The explicit formula for $\square = E^\alpha\Phi$ in terms of given $E_\alpha$ and
functions $c_{\alpha\beta\gamma}$ reads as

$$\Box E \Phi = E^\alpha E _\alpha \Phi + e^{\alpha\sigma} e_{\alpha\sigma} E_\sigma \Phi +$$

$$\Phi \left( \frac{2}{3} E_{\alpha\sigma} c_{\alpha\sigma} ^\sigma + \frac{1}{3} e^{\alpha\sigma} e_{\alpha\sigma} ^\beta c_{\alpha\beta} ^\sigma + \frac{1}{6} e_{\alpha\beta\sigma} c_{\alpha\beta\sigma} ^\sigma + \frac{1}{12} e_{\alpha\beta\sigma} c_{\alpha\beta\sigma} ^\sigma \right)$$

(2)

Here we are raising indices by means of the spinor metric tensor $e^{\alpha\beta}$ ($e^{\alpha\beta} = -e^{\beta\alpha}$, $e^{12} = 1$) and lowering by means of the inverse matrix $e_{\alpha\beta}$. We give a proof of this formula in appendix A. It is worth noting that the expression in parentheses in formula (2) is just the volume of the leaf evaluated with respect to the volume element specified above.

In order to get a function which is invariant with respect to a given integrable distribution it is not necessary to use the natural volume element related to the chosen pair of vector fields when integrating. One can perform the integration using any volume element as well. But since it differs from the natural volume element by multiplication by some function, all possible freedom is reflected by the following formula for a generic operation $\Box$:

$$\Box \Phi = \Box \rho \Phi \equiv \Box E (\rho \Phi)$$

(3)

where $\rho$ is some function. In short, this freedom is the same as modifying the initial function by multiplying it by some fixed function.

Now let us consider the case of a complex supermanifold $\mathcal{M}$ with a pair of complex vector fields $E_\alpha$ on it satisfying the integrability condition (1). In this case we have to modify the consideration above slightly. Formula (2) again defines an operation yielding a $E_\alpha$-invariant function if one assumes that $E_\alpha$ is a holomorphic vector field and $\Phi$ is a holomorphic function. As before by the Frobenius theorem we have an integral complex surface $\Sigma_C$ going through every point in $\mathcal{M}$. Our vector fields define a natural holomorphic volume element on $\Sigma_C$. For a generic real submanifold $\Sigma_R$, a real basis in a tangent space to $\Sigma_R$ can be considered as a complex basis in a tangent space to $\Sigma_C$. This means that the holomorphic volume form determines a nondegenerate volume element in a tangent space to $\Sigma_R$. Given such a generic submanifold $\Sigma_R \subset \Sigma_C$, one can perform an integration of a holomorphic function over it. One can show that when we have a purely odd supermanifold $\Sigma_C$, the result of integration does not depend on the particular choice of $\Sigma_R$. Therefore we see that in the case of a complex manifold, the operation $\Box E_\alpha$ (formula (2)) can be also interpreted in terms of integration over the leaves.

We are interested in applications of formula (2) in the framework of the induced geometry approach to the description of curved superspace. In this approach a curved superspace is described as a generic real $(4,4)$-dimensional surface $\Omega$ embedded into $\mathbb{C}^{4|2}$ (see [1] for details). The complex structure on $\mathbb{C}^{4|2}$ induces a CR-structure on $\Omega$. This means that a complex plane is singled out in every tangent space to $\Omega$. Namely if $T_z(\Omega)$ is the real $(4,4)$-dimensional tangent
space at a point \( z \) and \( J \) denotes the linear operator given by the multiplication by \( i \), then \( T_z(\Omega) \cap JT_z(\Omega) \) is the maximal complex subspace contained in \( T_z(\Omega) \). It is not difficult to figure out that this complex subspace is of dimension \((0, 2)\). One can choose vector fields

\[
E_\alpha, \bar{E}_\dot{\alpha}, E_c \quad (\alpha, \dot{\alpha} = 1, 2; c = 1, \ldots, 4)
\]

(4)
tangent to \( \Omega \) such that the fields \( E_\alpha \) form a (complex) basis of the complex subspace at each point, the fields \( \bar{E}_\dot{\alpha} \) define a basis in the complex conjugate plane and the fields \( E_c \) complete \( E_\alpha, \bar{E}_\dot{\alpha} \) to a (real) basis of the whole tangent space. The (anti)commutator of two vector fields tangent to \( \Omega \) is also a vector field tangent to \( \Omega \). Thus we have

\[
\{E_A, E_B\} = c_{AB}^D E_D
\]

(5)
where \( c_{AB}^D \) are some functions and the indices take on the values of the indices \( \alpha, \dot{\alpha}, c \). The fact that our CR-structure defined on \( \Omega \) is induced by a \( GL(4, 2|\mathbb{C}) \)-structure in the ambient space implies

\[
\{E_\alpha, E_\beta\} = c_{\alpha\beta}^\gamma E_\gamma
\]

(6)
and the corresponding complex conjugate equations. This means that we are dealing with an integrable CR-structure on \( \Omega \). We call a function \( \Phi \) defined on \( \Omega \) chiral if \( \bar{E}_\dot{\alpha} \Phi = 0 \). A function \( \Phi^+ \) is called antichiral if \( E_\alpha \Phi^+ = 0 \). Note that the restriction to \( \Omega \) of any holomorphic function defined in some neighborhood of \( \Omega \) in \( \mathbb{C}^{4|2} \) is a chiral function (the converse is also true in some sense, see \[\text{Footnote}\]). Formula (4) looks exactly like formula (1). The only difference is that in the case at hand the complex fields \( E_\alpha \) are defined on a real manifold \( \Omega \). Formula (2) can be used to construct an (anti)chiral function from an arbitrary given one. Moreover one can give an interpretation of formula (2) similar to those we gave for the purely real and complex cases using the complexification of \( \Omega \).

3 Formulation of N=1 SUYM in curved superspace in terms of induced geometry and CR-bundles

To construct the Lagrangian of N=1 super-Yang-Mills theory in the induced geometry approach, first we need to say more about induced CR-structures and introduce some useful geometric notions.

Note that the basis (3) of tangent vectors defining a CR-structure on \( \Omega \) is fixed up to linear transformations of the form

\[
E'_a = g_a^b E_b + g_a^\delta E_\beta + \bar{g}_a^\dot{\beta} \bar{E}_\dot{\beta}
\]

(7)

\[
E'_\alpha = g_\alpha^\beta E_\beta, \quad \bar{E}'_\dot{\alpha} = \bar{g}_\dot{\alpha}^\dot{\beta} \bar{E}_\dot{\beta}
\]

(8)
where \((g^b_a)\) is a real matrix and \((\bar{g}^\dot{\beta}_{\dot{\alpha}})\) is the complex conjugate matrix of \((g^\beta_\alpha E^\beta_\alpha)\). If \(\mathbb{C}^{4|2}\) is equipped with a volume element one can choose \(E_\alpha, E^\alpha\) to be a unimodular complex basis in the tangent space to \(\mathbb{C}^{4|2}\). This allows one to restrict the transformations (7) by the requirement

\[
\det(g^b_a) = \det(g^\dot{\beta}_{\dot{\alpha}})
\]

In this case we say that there is an induced SCR-structure on \(\Omega\). From now on we will assume that the basis (\(E_\alpha, E^\alpha\)) defines a SCR-structure.

For the functions \(c_{AB}^D\) defined by formula (3) in the case of induced SCR-structure we have the following identities

\[
c_{\alpha\beta}^\dot{\gamma} = c_{\alpha\beta}^d = 0, \quad c_{\alpha\beta}^\dot{\beta} = c_{\alpha\beta}^d
\]

and the corresponding complex conjugate ones.

We define the Levi matrix of the surface \(\Omega\) by the expression

\[
\Gamma^a_b = i\sigma^\alpha_{b\dot{\beta}} c^a_{\alpha\beta}
\]

where \(\sigma^b\) are the Pauli matrices for \(b = 1, 2, 3\) and the identity matrix for \(b = 0\). The matrix \(\Gamma^a_b\) coincides with the matrix of the Levi form defined in the standard way (see [1]).

To construct a Yang-Mills theory on \(\Omega \subset \mathbb{C}^{4|2}\), we start with two complex vector bundles \(\mathcal{F}\) and \(\mathcal{F}^+\) with structure group \(G\). Denote the Lie algebra corresponding to \(G\) by \(\mathfrak{g}\). Trivializing our bundles we can represent their sections \(\Phi, \Phi^+\) locally as vector functions, called fields. They describe matter and charge conjugated matter respectively. Gauge transformations correspond to the change of trivialization. They have the form

\[
\Phi' = e^{i\Lambda} \Phi \quad (\Phi^+) = e^{-i\bar{\Lambda}} \Phi^+
\]

where \(\Lambda\) and \(\bar{\Lambda}\) are some functions (sections of corresponding homomorphism bundles) with values in the representation of the Lie algebra \(\mathfrak{g}\) corresponding to the field \(\Phi\). We want to stress the fact that for now we consider \(\mathcal{F}\) and \(\mathcal{F}^+\) separately, not requiring them to be complex conjugate bundles (the functions \(\Lambda, \bar{\Lambda}\) in (11) are also independent). By SUYM fields we understand two pairs of semiconnections

\[
\nabla_\alpha \Phi^+ \equiv (E_\alpha + A_\alpha) \Phi^+ \quad \nabla_\alpha \Phi \equiv (\bar{E}_\alpha + \bar{A}_\alpha) \Phi
\]

restricted by the conditions

\[
\{\nabla_\alpha, \nabla_\beta\} = c_{\alpha\beta}^\gamma \nabla_\gamma \quad \{\bar{\nabla}_{\dot{\alpha}}, \bar{\nabla}_{\dot{\beta}}\} = c_{\dot{\alpha}\dot{\beta}}^\dot{\gamma} \bar{\nabla}_{\dot{\gamma}}
\]

These conditions mean that the corresponding semiconnections have vanishing curvature. It can be shown (see for example [5]) that semiconnections \(\nabla_\alpha\) satisfying (12) determine a CR-bundle structure on \(\mathcal{F}^+\), i.e. \(\mathcal{F}^+\) can be pasted
together from trivial bundles by chiral gluing functions. One can define the chiral sections as those annihilated by $\bar{\nabla}_\alpha$. Then condition (12) guarantees that there are sufficiently many of them. Analogously, given $\nabla_\alpha$ satisfying (12), one obtains a CR-bundle. By CR-bundle we mean a bundle whose gluing functions are antichiral. Moreover, one can take this property as a definition of CR and CR-bundles (see [5] for details). Thus the basic geometrical objects we start with are the surface $\Omega \subset C^{4|2}$, the CR-bundle $F^+$ and the CR-bundle $F$, both defined over $\Omega$. The solutions to the zero curvature equations (12) can be written locally as

$$ A_\alpha = e^{U} E_\alpha e^{-U} \quad A_\alpha = e^{-\bar{U}} \bar{E}_\alpha e^{\bar{U}} $$

where $U$ and $\bar{U}$ are some $G$-valued fields. Note that the fields $e^{-U}$ and $e^{\bar{U}}$ are determined by (13) only up to the left multiplication by arbitrary antichiral and chiral fields respectively. If we want to write a gauge invariant Lagrange function for chiral fields we will immediately encounter the difficulty in writing the kinetic term, which for the case of free chiral fields is simply $\Phi^+ \Phi$. This difficulty is due to the fact that in the case at hand these fields are sections of different bundles. Therefore we are forced to identify CR and CR bundles choosing a section of the bundle $F^+ \otimes F^*$. Here $F^*$ is the dual to the bundle $F$. This section we denote by $e^V$. Under the gauge transformations the field $e^V$ transforms in the following way

$$ e^{V'} = e^{-i\Lambda} e^{V} e^{i\bar{\Lambda}} $$

Now we can take the gauge invariant combination $\Phi^+ e^V \Phi$ as a kinetic density term.

Our next goal is to describe a gauge invariant theory in terms of the fields $\Phi_i, \Phi^+_i, e^{-U}, e^{\bar{U}}, e^V$ defined on our curved superspace $\Omega$. This will be done in a manifestly SCR-covariant way, i.e. independently of the choice of basis vector fields (3) up to local SCR-transformations (7). Instead of the customary Lorentz connections in our construction of the Lagrangian, we use only objects defined by the internal geometry of the superspace $\Omega$, namely the Levi matrix $\Gamma$ and the functions $c_{AB}^D$. We postpone the details of this construction until appendix B. Now we want to formulate the main result. The Lagrangian has the following form

$$ S = \int dV \left[ \frac{1}{k} (\det \Gamma)^{-1} (\Box_D e^G E^\alpha e^{-G}) (\Box_D e^G E_\alpha e^{-G}) \right] + $$

$$ + \int dV \left[ \frac{1}{4} \det \Gamma \right] \Phi^+ e^V \Phi + $$

$$ + \int dV \left[ a_i \Phi_i + \frac{1}{2} m_{ij} \Phi_i \Phi_j + \frac{1}{3} g_{ijk} \Phi_i \Phi_j \Phi_k \right] + h.c. $$

(15)
Here as in flat space the N=1 SUYM Lagrangian contains a Lagrangian of gauge fields, a kinetic term of chiral fields and a term describing the interaction between chiral fields. In (15) we are using the following notations: $k$ is a coupling constant, $e^G = e^{U} - e^{V}$, $a_i, m_{ij}, g_{ijk}$ are coupling constants which must be chosen in a way that ensures the gauge invariance of matter-matter interaction, $\Box_E$ is a chiral projector whose general form was described in section 2, and $\Box_F$ is a “covariant” chiral projector which is constructed as $\Box_E$ with derivatives $\overline{E}_\alpha$ replaced by “covariant” derivatives $\overline{D}_\alpha$. More precisely, $\overline{D}_\alpha$ acts on an arbitrary tensor $V$ carrying the spinor index $\alpha$ in the following way

$$\overline{D}_\alpha V_\beta = \overline{E}_\alpha V_\beta - \tilde{c}_{\dot{\alpha}, \sigma} V_\sigma,$$

where $\tilde{c}_{\dot{\alpha}, \sigma}$ are some functions which can be expressed in terms of $c_{ABD}$ (see formulae (36), (37), (38) in appendix B). Note that the integrands in (15) are chiral functions (for the first term this fact follows from its construction, which is explained in details in appendix B). The integration in (15) should be understood as a chiral integration. If a chiral function can be extended to a holomorphic function in some domain in $\mathbb{C}^{4\mid2}$ where it can be integrated with the holomorphic volume element over a real $(4,2)$-submanifold contained in $\Omega$.

We have constructed a Lagrangian depending on the field $V$, the specific combination of the fields $e^U, e^\tilde{U}, e^V$ which we denoted by $e^G$, and matter fields $\Phi_i, \Phi_i^\dagger$. All these fields are complex. In order to perform a functional integration over the fields $U, \tilde{U}, V, \Phi_i$ one has to restrict these fields to a real surface in functional space. Once this surface is chosen this will restrict our large gauge group to a smaller one. The real surface we choose is given by the equations

$$\overline{\Phi}_i = \Phi_i^\dagger, \ V = \overline{V}^t, \ -\tilde{U} = \overline{U}$$

where the upper bar denotes complex conjugation. The gauge transformations preserving these reality conditions are of the form (11), (14) where $\Lambda$ and $\tilde{\Lambda}$ are complex conjugates of each other and $T = \tilde{S}^+ = S$. Still we have a rather large gauge group. Let us do a partial gauge fixing by requiring that $e^U = 1$. By reality conditions (16) this implies $e^{-\tilde{U}} = 1$ and therefore $e^{-G} = e^V$, which means that our Lagrangian (15) contains only the field $V$ in this partial gauge fixing. The remaining gauge group contains the transformations with $i\Lambda = \tilde{S}^+, \ -i\tilde{\Lambda} = S$, i.e. the antichiral transformations of the fields $\Phi_i$ and the corresponding complex conjugate chiral transformations of the fields $\Phi_i^\dagger$.

4 Comparison with the conventional approach

In this section we want to compare our constructions with the conventional (Wess-Zumino) approach to supergravity and super Yang-Mills theory in curved superspace which is presented in [2] in great detail. We start with a comparison of formula (3) with the standard formula for (anti)chiral projection operators
(see [2], chapter 19). But first let us recall briefly the main constituents of the conventional approach. In this approach we have a connection defined on a tangent bundle over (4,4)-dimensional real superspace with the Lorentz group as structure group. Another dynamical variable in this approach is a vielbein \( E_A^M \) which defines the covariant derivatives \( D_M \) on the tangent bundle and also identifies the Lorentz bundle with the tangent one allowing to transform world indices into Lorentz indices and vice-versa. Here we are working with Lorentz indices. Under a certain set of constraints (see [2]) the vector fields \( E_\alpha = E_A^M \partial_M \) are closed with respect to the anticommutator. Moreover the functions \( c_{\alpha\beta} \) defining the anticommutation relations satisfy the following condition

\[
c_{\alpha\beta\gamma} = \omega_{\alpha\beta\gamma} + \omega_{\beta\alpha\gamma}
\]

(17)

where \( \omega_{\alpha\beta\gamma} \) are the connection coefficients for the covariant differentiation of spinor fields, having only Lorentz indices.

The antichiral projector acts on arbitrary function \( \Phi \) as follows

\[
(D^\alpha D_\alpha - 8 R^+) \Phi
\]

(18)

and gives an antichiral function as a result. Here \( R^+ = \frac{1}{24} R_{\alpha\beta}^{\alpha\beta} \) denotes the invariant obtained from the curvature tensor

\[
R_{\alpha\beta\delta}^{\gamma} = E_\alpha \omega_{\beta\delta}^{\gamma} + E_\beta \omega_{\alpha\delta}^{\gamma} + \omega_{\alpha\sigma}^{\gamma} \omega_{\beta\delta}^{\sigma} + \omega_{\beta\sigma}^{\gamma} \omega_{\alpha\delta}^{\sigma} - c_{\alpha\beta} \omega_{\sigma\delta}^{\gamma}
\]

After the contraction of indices this gives the following expression for \(-8 R^+\)

\[
-8 R^+ = -(E_\alpha \omega_{\beta}^{\alpha\beta} + E_\beta \omega_{\alpha}^{\alpha\beta} + \omega_{\beta\sigma}^{\gamma} \omega_{\alpha\delta}^{\sigma} + \omega_{\delta\sigma}^{\gamma} \omega_{\alpha\beta}^{\sigma}) + \omega_{\delta\beta}^{\gamma} \omega_{\alpha\sigma}^{\alpha}
\]

(19)

The induced geometry approach to supergravity has been shown to be equivalent to the Wess-Zumino one ([1]). Thus one can use formula (2) to obtain (18). We identify the pair of vector fields \( E_\alpha \) appearing in the Wess-Zumino approach with those in definition of the CR-structure induced on \( \Omega \) (i.e. complex conjugate to the corresponding CR-structure). The formulae (18) and (2) must be equivalent at least up to the freedom described by formula (3). Indeed as one can easily check, substituting relation (17) in (2), we will get exactly formula (18) with the term \(-8 R^+\) expressed as in (19).

If one starts with the induced geometry approach then in order to get the Lorentz gauge group one has to require the following gauge condition

\[
c_{\alpha\beta}^{\gamma} = 2i \sigma_{\alpha\beta}^{\gamma}
\]

where \( \sigma_{\alpha\beta}^{\gamma} \) are the Pauli matrices. This condition fixes the SCR-basis ([4]) up to transformations of the form ([5]) where \( \det(g^\alpha_\beta) = \det(g_\alpha^\beta) = \det(g^\beta_\alpha) = 1 \), i.e. up to the Lorentz transformations. In this gauge 4(\det \Gamma)^{-1} = 1, which as it is shown in appendix B (see formulae (39) and (40)) implies

\[
\tilde{c}_{\gamma,\sigma\gamma}^{\sigma\delta} = \tilde{c}_{\gamma,\sigma\gamma}^{\delta\sigma} = 0
\]
Moreover, the quantities $-\tilde{c}_{\alpha,\beta\sigma}^{\dot{\sigma}}$, $-\tilde{c}_{\alpha,\sigma\beta}^{\dot{\sigma}}$ transform now as coefficients of the Lorentz connection (as can be seen from (43) for Lorentz transformations). Thus it seems reasonable to identify $-\tilde{c}_{\alpha,\beta\sigma}^{\dot{\sigma}}$, $-\tilde{c}_{\alpha,\sigma\beta}^{\dot{\sigma}}$ with the connection coefficients $\omega_{\alpha,\beta}^{\dot{\sigma}}$, $\omega_{\alpha,\sigma\beta}^{\dot{\sigma}}$, respectively, from the conventional approach. The last assumption implies

$$\bigtriangledown_\alpha V_\alpha = (\bar{D}_\gamma \bar{D}_{\dot{\gamma}} - 8R) V_\alpha$$

and the corresponding complex conjugate identity. This shows that in the gauge specified above, the Yang-Mills Lagrangian term from (15) reduces to the standard one, which is written in terms of field strengths $\bar{W}_\alpha$. Indeed, the identities $-\tilde{c}_{\alpha,\beta\sigma}^{\dot{\sigma}}, -\tilde{c}_{\alpha,\sigma\beta}^{\dot{\sigma}}$ are true. One can derive them from the standard set of torsion constraints

$$T_{\alpha\beta}^a = 0, T^a_{\alpha\beta} = T^a_{\dot{\alpha}\dot{\beta}} = 0, T^a_{\alpha\dot{\beta}} = 2i\sigma^a_{\alpha\dot{\beta}}, T^a_{\dot{\alpha}\beta} = 0, T^a_{\dot{\alpha}\dot{\beta}} = 0$$

(20)

where $\alpha$ denotes either $\alpha$ or $\dot{\alpha}$. Finally, the volume element $dV$ we used in (15) is nothing but the chiral volume element of the conventional approach, usually denoted as $\mathcal{E}d^2\Theta$. This completes the derivation of the conventional picture from the induced geometry one.

### A Explicit formula for the chiral projector

To prove formula (2) let us first consider the case of anticommuting vector fields, i.e. the functions $c_{\alpha,\beta\sigma}^{\dot{\sigma}}$ are identically zero.

Let $X^M$ be coordinates on $M$ ($M = 1, \cdot \cdot \cdot , m + n$), so we have $E_\alpha = E_\alpha^M \partial_M$ where $\partial_M$ is the derivative with respect to $X^M$ and $E_\alpha^M$ are the coordinates of the field $E_\alpha$ et the surface $\Sigma$ be defined in a parametric form $X^M(\xi)$ where $\xi = (\xi^\alpha), (\alpha = 1, 2)$ are odd coordinates on $\Sigma$. Then we have the following equation for $\Sigma$

$$\frac{\partial X^M(\xi)}{\partial \xi^\alpha} = E_\alpha^M(X(\xi))$$

(21)

Since we assume the fields $E_\alpha$ anticommutate we can integrate this equation expanding $X^M(X_0, \xi)$ in $\xi$ (here $X_0$ is an initial data for the system (21) ). The solution to (21) reads as follows

$$X^M(X_0, \xi) = X_0^M + \xi^\alpha E_\alpha^M + \frac{1}{4} \xi \xi E^\sigma(E_\sigma^M)$$

(22)

where $E_\alpha^M$ and $E^\sigma(E_\sigma^M)$ are evaluated at the point $\xi = 0$ and $\xi \xi = \xi^\alpha \xi_\alpha$. The natural volume element on $\Sigma$ in $\xi$-coordinates is simply $d\xi d\xi$. Expanding the
given function $\Phi$ in $\xi$ and performing the integration we get

$$\langle \square E \Phi \rangle(X_0) = \int \Phi(X(X_0, \xi))d\xi d\xi = E^\alpha E_\alpha(\Phi)$$

(23)

which is obviously annihilated by $\tilde{E}_\alpha$. As one can easily see formula (23) coincides with (2) when $c_{\alpha \beta}^\gamma = 0$. Thus for the case of anticommuting vector fields formula (2) is proved.

Now let us consider two different pairs of vector fields $E_\alpha$ and $\tilde{E}_\alpha$ defining the same integrable distribution. Let $c_{\alpha \beta}^\gamma$ and $\tilde{c}_{\alpha \beta}^\gamma$ be the corresponding functions defining the anticommutation relations for each pair. The fields $E_\alpha$ and $\tilde{E}_\alpha$ are connected by means of some nondegenerate even matrix $A$:

$$\tilde{E}_\alpha = A^\alpha_\beta E_\beta \quad (24)$$

The transformation law for the functions $c_{\alpha \beta}^\gamma$ reads as follows

$$\tilde{c}_{\alpha \beta}^\gamma = A^\delta_\alpha A^\beta_\beta (A^{-1})^\gamma_\mu c_{\delta \sigma}^\mu + (A^{-1})^\gamma_\beta (A^\delta_\alpha E_\alpha A^\beta_\beta + A^\delta_\beta E_\sigma A^\sigma_\alpha) \quad (25)$$

Since the ratio of the natural supervolume element corresponding to the vector fields $\tilde{E}_\alpha$ to the one of the fields $E_\alpha$ equals $\det A$, the operations $\square E$ and $\square \tilde{E}$ are related by the following formula

$$\square \tilde{E} \Phi = \square E (\det A \Phi) \quad (26)$$

Note that since our distribution is integrable, given the vector fields $E_\alpha$ we can perform a linear transformation (24) such that the new vector fields $\tilde{E}_\alpha$ will anticommute. Thus if we prove that the right hand side of (2) under the transformations (24) satisfies (2) then we will have proved formula (2). We are going to show this for the case of infinitesimal transformations $A^\beta_\beta = \delta^\beta_\beta + a^\beta_\alpha$, where $a^\beta_\alpha$ is an infinitesimally small matrix. But this also proves (4) because every finite transformation (24) can be obtained as a succession of infinitesimally small ones.

Keeping only the terms of the first order in $a$, we have the following transformations

$$\tilde{E}_\alpha = E_\alpha + a^\beta_\alpha E_\beta, \quad \tilde{E}^\alpha = E^\alpha - \bar{a}^\alpha_\beta E^\beta$$

$$\tilde{c}_{\alpha \beta}^\gamma = c_{\alpha \beta}^\gamma + a^\sigma_\alpha c_{\sigma \beta}^\gamma + a^\gamma_\beta c_{\alpha \sigma}^\gamma - a^\sigma_\beta c_{\alpha \gamma}^\sigma + E_\alpha a^\gamma_\beta + E_\beta a^\gamma_\alpha \quad (27)$$

where $\bar{a}^\alpha_\beta = \epsilon_{\alpha \nu} \epsilon^{\beta \mu} a^\nu_\mu$. Note that $a^\beta_\alpha - \bar{a}^\beta_\alpha = \delta^\beta_\alpha tr(a)$ where $\delta^\beta_\alpha$ is the Kronecker symbol. Another useful identity is

$$E^\delta E_\sigma a^\gamma_\delta = \frac{1}{2} E^\delta E_\beta (tr(a)) + \frac{1}{2} \epsilon^\delta_\alpha \gamma E_\gamma a^\alpha_\delta$$

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For the bilinear combinations in $c_{\gamma}^{\alpha \beta}$ which enter (2) and for the operator $E^\alpha E_\alpha$, up to the first order in $a$ we have

\[
\tilde{c}_{\sigma}^{\alpha \beta} \tilde{c}_{\alpha \sigma}^{\beta} = (1 + \text{tr}(a))c_{\sigma}^{\alpha \beta} c_{\alpha \beta}^{\sigma} + 2c_{\sigma}^{\alpha \beta} E_\alpha a_\beta^\sigma + 2c_{\alpha \beta}^{\sigma} E_\beta a_\alpha^\sigma \\
\tilde{c}_{\alpha \beta}^{\sigma} \tilde{c}_{\alpha \beta}^{\alpha \sigma} = (1 + \text{tr}(a))c_{\sigma}^{\alpha \beta} c_{\sigma}^{\alpha \beta} + 2c_{\sigma}^{\alpha \beta} E_\alpha a_\beta^\sigma - 2c_{\sigma}^{\alpha \beta} E_\beta a_\alpha^\sigma \\
\tilde{c}_{\alpha \beta}^{\sigma} E_\alpha = (1 + \text{tr}(a))E_\alpha E_\alpha + (E^\alpha E_\alpha E_\alpha) (28)
\]

Note that the appearance of the combination $1 + \text{tr}(a)$ in (28) is natural because for the volume preserving transformations (24) the contraction of upper and lower indices becomes an invariant operation with respect to the derivative independent part of transformation (25). Substituting the expressions (27) and (28) in the right hand side of (2) using the identities mentioned above and keeping terms only up to the first order in $a$, we obtain the following expression

\[
(1 + \text{tr}(a))E^\alpha E_\alpha \Phi + ((1 + \text{tr}(a))c^\alpha_{\sigma} c_\sigma^\alpha + 2(E^\alpha \text{tr}(a)))E_\alpha \Phi + \\
\Phi(E^\alpha E_\alpha \text{tr}(a) + c^\alpha_{\sigma} E_\sigma \text{tr}(a) + (1 + \text{tr}(a))\mathcal{V}(\Phi)) (29)
\]

where for compactness of notation we denoted the term multiplying $\Phi$ in (2) by $\mathcal{V}(\Phi)$. Now after some trivial transformations one can easily single out the factor $(1 + \text{tr}(a))$ in (29) and get

\[
E^\alpha E_\alpha (\Phi(1 + \text{tr}(a)) + c^\alpha_{\sigma} E_\sigma \Phi(1 + \text{tr}(a))) + \Phi(1 + \text{tr}(a))\mathcal{V}(\Phi) = \Box E(\Phi(1 + \text{tr}(a))) (30)
\]

The last equation means that the right hand side of (2) satisfies the infinitesimal version of (26), so that by the above considerations it really represents the operation $\Box_E$.

**B Construction of the Lagrangian**

We start with the term describing the interaction between matter chiral fields $\Phi_i$. It has the form

\[
\int dV \mathcal{U}(\Phi) + h.c. = \int dV \left[ a_i \Phi_i + \frac{1}{2} m_{ij} \Phi_i \Phi_j + \frac{1}{3} g_{ijk} \Phi_i \Phi_j \Phi_k \right] + h.c. (31)
\]

where the function $\mathcal{U}(\Phi)$ is assumed to be gauge invariant. The integration here is as described in section 3. As we do not have any metric, we have to restrict ourselves to SCR-transformations in order to preserve the volume element $dV$.

To construct the kinetic term one needs to apply first a chiral projector $\Box$ to the gauge invariant quantity $\Phi^+ e^{\mathcal{V}} \Phi$ and then perform the integration
described above. The explicit construction of a chiral projector in terms of internal geometry of the surface \( \Omega \) was given in section 2. It was shown there that the action of a generic chiral projector can be written as follows

\[
\Box_E \rho \Phi^+ e^V \Phi
\]

where \( \Box_E \) is a chiral projector corresponding to a natural volume element defined by the basis fields \( \bar{E}_\alpha \), the explicit construction of which is given by formula (2), and \( \rho \) is an arbitrary function. Note that under a change of basis, the vector fields \( \bar{E}_\alpha \) corresponding to (7) the operation \( \Box_E \) behave as follows

\[
\Box_{E'} = \Box_E \det(\tilde{g})
\]

(see also formulae (24) and (26) in appendix A). Therefore to have SCR covariance one has to choose the transformation law for \( \rho \) to be \( \rho' = \rho \det \Gamma \). It can be easily checked that \( \rho = |\det \Gamma|^{-\frac{1}{3}} \) has this property. Here \( |\det \Gamma| \) stands for the absolute value of the determinant of the Levi form corresponding to the surface \( \Omega \). Therefore the SCR covariant expression of the kinetic term reads as

\[
\int dV \left[ \Box_E \left| \frac{1}{4} \det \Gamma \right|^{-\frac{1}{3}} \Phi^+ e^V \Phi \right] + h.c. \tag{32}
\]

where \( \frac{1}{4} \) is inserted for normalization purposes.

Finally let us turn to the construction of the Lagrangian of gauge fields. This turns out to be the most technically complicated part of the whole construction. First note that due to our large gauge invariance group, (11) when \( \Lambda \) and \( \bar{\Lambda} \) are arbitrary functions with values in the gauge Lie algebra it is impossible to construct the gauge strength fields \( W_\alpha, \bar{W}_\alpha \) corresponding to each of fields \( e^V, e^{-U}, e^\bar{U} \). Consider the combination \( e^G = e^\bar{U} e^{-V} e^{-U} \), which is invariant under gauge transformations. Due to the arbitrariness in the choice of solutions \( U \) and \( \bar{U} \) to the zero curvature equations (12) mentioned in section 3, there is an ambiguity in the definition of \( e^G \) described by the transformation

\[
e^{G'} = e^T e^G e^{(S^*)^t} \tag{33}
\]

where \( T \) and \( S^+ \) are arbitrary respectively chiral and antichiral \( \mathcal{G} \)-valued functions.

The generic form of chiral field strengths for \( e^G \) reads as

\[
W_\alpha = \Box_E B^\beta_\alpha e^G E_\beta e^{-G} \tag{34}
\]

where \( B^\beta_\alpha \) is a matrix to be defined by the invariance properties of \( W_\alpha \). The transformation law for \( B^\beta_\alpha \) under (7) can be derived in terms of the invariance of construction (34) under (7), and has the form
\((B')^\beta_\alpha = \frac{1}{\det(\hat{g})} B^\beta_\gamma (g^{-1})^\gamma_\alpha \)  

(35)

Further restrictions on the matrix \( B \) come from the covariance requirement under the substitution (33). For \( W_\alpha \) to transform covariantly under (33), the operator \( \square_E B^\alpha_\beta E_\beta \) ought to annihilate chiral functions (the annihilation of antichiral functions is obvious). It is convenient to derive the restriction on the matrix \( B \) first in a special basis, namely where \( \{ \hat{E}_\alpha, \hat{E}_\beta \} = 0 \) (see appendix A for details). But first let us introduce some useful notations. One can choose the vectors \( E_{\alpha\beta} = \{ E_\alpha, E_\beta \} \) as a basis of the real tangent space to \( \Omega \). Then we have the following commutation relations

\[
[E_\gamma, E_{\alpha\beta}] = \hat{c}_{\gamma,\alpha\beta} \sigma^\sigma E_{\sigma\sigma} + \hat{c}_{\gamma,\alpha\beta} \sigma^\sigma + \hat{c}_{\gamma,\alpha\beta} \hat{E}_\gamma
\]

(36)

where \( \hat{c}_{\gamma,\alpha\beta} A \) are some functions. The connection between the two bases introduced in real space is given by the following formulas

\[
E_{\alpha\beta} = c_{\alpha\beta} \sigma^a E_a + c_{\alpha\beta} \gamma^a E_\gamma + c_{\alpha\beta} \hat{\gamma} \hat{E}_\gamma
\]

(37)

\[
E_a = \hat{c}_{a\beta} E_{\alpha\beta} - c_{\alpha\beta} \gamma^a E_\gamma - c_{\alpha\beta} \hat{\gamma} \hat{E}_\gamma
\]

where \( \hat{c}_{a} A = i(\Gamma^{-1})_a^b \sigma^\alpha \sigma^\beta c_{\alpha\beta} A \). We use the notation \( \sigma^\sigma \) for the Pauli matrices and \( \Gamma^{-1} \) is the inverse of the Levi form matrix \( \Gamma^b_a = i\sigma^\alpha \sigma^\beta c_{\alpha\beta} b \). Substituting the first expression (37) into the commutator on the left hand side of (36), one obtains the following expression for the quantity \( \hat{c}_{\gamma,\sigma\beta} \sigma^\sigma \) in terms of the functions \( c_{A} B \)

\[
\hat{c}_{\gamma,\sigma\beta} \sigma^\sigma = \hat{c}_{\sigma} \sigma^a E_{\gamma\sigma\beta} a + \hat{c}_{\sigma} \sigma^b c_{\sigma\beta} b - c_{\gamma} \hat{\gamma} \hat{\gamma}
\]

(38)

Contracting the upper index \( \hat{\sigma} \) with the lower \( \hat{\beta} \) we obtain the following important identity which we will need later

\[
\hat{c}_{\gamma,\sigma\beta} \hat{\sigma} \hat{\beta} = (\det \Gamma)^{-1} E_\gamma \det \Gamma
\]

(39)

where in the last step we used the identity (10). The complex conjugate to equation (39) reads as

\[
\hat{c}_{\gamma,\sigma\beta} \hat{\sigma} \hat{\beta} = (\det \Gamma)^{-1} E_\gamma \det \Gamma
\]

(40)

Now we are ready to explore the restrictions on the matrix \( B \) in (34). For the basis specified above, applying the operator \( \square_E B^\alpha_\beta E_\beta \) to some chiral function \( \Phi \) we get

\[
\square_E B^\alpha_\beta E_\beta \Phi = \hat{E}_\gamma \hat{E}_\gamma B^\alpha_\beta E_\beta \Phi = \{ \hat{E}_\gamma, \{ \hat{E}_\gamma, B^\alpha_\beta E_\beta \} \} \Phi = 2B^\alpha_\delta \epsilon^{\delta \gamma} ((B^{-1})_j^\gamma \hat{E}_\delta B^\alpha_\gamma + \hat{c}_{\delta,\hat{\sigma} \hat{\gamma}} \hat{\sigma}) E_{\sigma\gamma} \Phi + \hat{f}_a \hat{E}_\beta \Phi
\]

(41)
where by \( f^\alpha_\beta \) we denote some functions which can be expressed in terms of the matrix \( B \) and its derivatives, but are unessential for further analysis. The last remark is due to the fact that if the coefficient of \( E_{\sigma\bar{\sigma}} \) vanishes, then the whole operator vanishes on chiral functions. This is because the operator on the left hand side of (41) acts from chiral functions to chiral ones (for the fixed index \( \alpha \)) and it can be easily shown that an operator with this property and of the form \( f^\alpha_\beta E_\beta \) is identically zero. Therefore we have the following equation for the matrix \( B \)

\[
(B^{-1})^\gamma_\delta \bar{E}_\delta B^\gamma_\sigma = -\bar{c}_{\delta,\sigma}^\delta \dot{\sigma} \sigma
\]

(42)

It can be derived directly from the definition (36) that the transformation law for the coefficients \(-\bar{c}_{\delta,\sigma}^\delta \dot{\sigma} \sigma\) under (7) is as follows

\[
- (\dot{c}'_{\delta,\sigma}^\delta \dot{\sigma} \sigma) = -\bar{c}_{\delta,\sigma}^\delta \dot{\sigma} \sigma - \bar{c}_{\delta,\sigma}^\delta \dot{\sigma} \sigma
\]

where \( \delta \) is the Kronecker symbol. The first two terms in this expression are of the same form as the transformation law for semiconnection coefficients. One can easily check using (35) that the transformation rule for the expression on the left hand site of (42) is of the same form, which is of course what one should expect since the operator under consideration is defined invariantly. This invariance implies that it is sufficient to prove the solvability of (42) in the special basis which we have already used. In this basis the integrability condition for (42) takes the form

\[
\bar{E}_{\alpha} \bar{c}_{\bar{\beta},\bar{\sigma}} \dot{\sigma} \sigma + \bar{E}_\beta \bar{c}_{\alpha,\bar{\sigma}} \dot{\sigma} \sigma - \bar{c}_{\alpha,\bar{\sigma}} \dot{\alpha} \sigma \bar{c}_{\bar{\beta},\bar{\sigma}} \dot{\sigma} \sigma - \bar{c}_{\bar{\beta},\bar{\sigma}} \dot{\sigma} \sigma \bar{c}_{\alpha,\bar{\sigma}} \dot{\sigma} \sigma = 0 \quad (44)
\]

This formula follows from two identities

\[
\{ \bar{E}_{\alpha}, [\bar{E}_{\beta}, E_{\gamma\bar{\gamma}}] \} = 0
\]

[\bar{E}_{\beta}, E_{\gamma\bar{\alpha}}] + [\bar{E}_{\alpha}, E_{\gamma\bar{\beta}}] = 0
\]

(45)

which are true when \( \{ \bar{E}_{\alpha}, \bar{E}_{\beta} \} = 0 \). The second identity in (45) can be equivalently written as

\[
\bar{c}_{\alpha,\bar{\beta}}^A A + \bar{c}_{\alpha,\bar{\beta}}^A A = 0
\]

Writing the first identity in (45) in terms of the basis \( E_{\alpha}, \bar{E}_{\alpha}, E_{\alpha\alpha} \), for the coefficient of \( E_{\sigma\bar{\sigma}} \) we have

\[
-\bar{c}_{\bar{\beta},\bar{\gamma}}^\mu \bar{c}_{\alpha,\bar{\mu}} \dot{\sigma} \sigma + E_{\alpha} \bar{c}_{\bar{\beta},\bar{\gamma}} \dot{\sigma} \sigma + \delta^\beta_\alpha \bar{c}_{\bar{\beta},\bar{\gamma}} \dot{\sigma} \sigma = 0
\]

Contracting the indices \( \bar{\sigma} \) and \( \bar{\gamma} \) in the last expression, symmetrizing it over the indices \( \bar{\alpha}, \bar{\beta} \) and using the second identity in (45) we get (44). To summarize, we showed that equation (42) is solvable. Note however that there is a large
arbitrariness in the choice of solutions to (42). Namely, one can multiply any solution by a matrix with chiral entries and get another one. We will show how to narrow the choice of solution. But first let us write the Lagrangian $W^\alpha W_\alpha$ in terms of the construction (34) taking into account (42). After some trivial transformations we will get

$$W^\alpha W_\alpha = \det B(\Box_D e^G E^\alpha e^{-G})(\Box_{\bar{D}} e^G E_\alpha e^{-G})$$

where $\Box_D$ is obtained from the usual chiral projector $\Box_{\bar{E}}$ by substitution of the derivatives $\bar{E}_\alpha$ by covariant derivatives $\bar{D}_\alpha$ which act on a tensor $V$ carrying a spinor index $\alpha$ in the following way

$$\bar{D}_\alpha V_\beta = \bar{E}_\alpha V_\beta - \bar{\epsilon}_{\alpha,\beta}^{\dot{\sigma}\sigma} V_{\sigma}$$

The only quantity in (46) which depends on the choice of the solution to (42) is $\det B$. From (12) we have

$$\det B \bar{E}_\gamma (\det B)^{-1} = \bar{\epsilon}_{\gamma,\sigma}^{\dot{\sigma}\sigma}$$

As one can see from (40), the determinant of the Levi form $\Gamma$ satisfies exactly the same equation. Therefore one can take $\det \Gamma$ multiplied by a constant factor as a solution to the last equation. We choose $\det B = 4(\det \Gamma)^{-1}$. With this choice of constant factor we will get the standard form for our Lagrangian in a Wess-Zumino gauge for the basis vector fields.

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