On Lagrangian Formulations for Arbitrary Bosonic HS Fields on Minkowski Backgrounds

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Abstract—We review the details of unconstrained Lagrangian formulations for Bose particles propagated on an arbitrary dimensional flat space-time and described by the unitary irreducible integer higher-spin representations of the Poincare group subject to Young tableaux \( Y(s_1, ..., s_k) \) with \( k \) rows. The procedure is based on the construction of scalar auxiliary oscillator realizations for the symplectic \( sp(2k) \) algebra which encodes the second-class operator constraints subsystem in the HS symmetry algebra. Application of an universal BRST approach reproduces gauge-invariant Lagrangians with reducible gauge symmetries describing the free dynamics of both massless and massive bosonic fields of any spin with appropriate number of auxiliary fields.

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1. INTRODUCTION

Growth of the interest to higher-spin (HS) field theory is mainly stipulated by the hopes to reconsider the problems of an unique description of variety of elementary particles and all known interactions especially due to expected output of LHC on the planned capacity. Remind, that it suspects both the proof of supersymmetry in infinite tower of bosonic and fermionic HS fields it can be treated as an approach to study a structure of superstring theory, which operates with an infinite tower of bosonic and fermionic HS fields it can be treated as an approach to study a structure of superstring theory from field-theoretic viewpoint. Some of the aspects of current state of HS field theory are discussed in the reviews [2]. The paper considers the last results of constructing Lagrangian formulations (LFs) for free integer both massless and massive mixed-symmetry tensor HS fields on flat \( \mathbb{R}^{1,d-1} \)-space-time subject to arbitrary Young tableaux (YT) \( Y(s_1, ..., s_k) \) in Fronsdal metric-like formalism within BFV-BRST approach [3], and based on the results presented in [4] (see [4], for detailed bibliography).

It is known that for \( d > 4 \) space-time dimensions, there appear, besides totally symmetric irreducible representations of Poincare or (Anti)-de-Sitter ((A)dS) algebras the mixed-symmetry representations determined by more than one spin-like parameters [5, 6]. Whereas for the former ones the LFs both for massless and massive free higher-spin fields is well enough developed [7–11] as well as on base of BFV-BRST approach, e.g. in [12–15], for the latter the problem of their field-theoretic description is not completely solved. So, the main result within the problem of LF for arbitrary massless mixed-symmetry HS fields on a Minkowski space-time was obtained in [16] with use of unfolded form of equations of motion for the field in “frame-like” formulation. In the “metric-like” formulation corresponding Lagrangians were derived in closed manner for only reducible Poincare group \( ISO(1, d - 1) \) representations in [17].

2. INTEGER HS SYMMETRY ALGEBRA FOR BOSONIC FIELDS

A massless integer spin Poincare group irrep in \( \mathbb{R}^{1,d-1} \) is described by rank \( \sum_{i=1}^{k} s_i \) tensor field \( \Phi_{(\mu^1 \nu_1)(\mu^2 \nu_2) ... (\mu^k \nu_k)}(x) \) with generalized spin \( s = (s_1, s_2, ..., s_k) \), \( s_1 \geq s_2 \geq ... \geq s_k > 0, k \leq (d/2) \) subject to a YT with \( k \) rows of lengths \( s_1, s_2, ..., s_k \)

\[
\begin{array}{cccccc}
\mu^1 & \mu^2 & \cdots & \cdot & \cdots & \mu_i \\
\mu^2 & \mu^3 & \cdots & \cdot & \cdots & \mu_{i+1} \\
\vdots & \vdots & \ddots & \cdot & \cdots & \vdots \\
\mu_i & \mu_{i+1} & \cdots & \cdot & \cdots & \mu_k \\
\end{array}
\]

The field is symmetric with respect to the permutations of each type of Lorentz indices \( \mu^i \), for \( \eta_{\mu \nu} = \text{diag}(+, -, ..., -) \), \( \mu, \nu = 0, 1, ..., d - 1 \) and obeys to the Klein–Gordon, divergentless (2), traceless (3) and...
mixed-symmetry equations (4) [for $i,j = 1, \ldots, k$; $l$, $m_i = 1, \ldots, s_j$]:

$$\partial^{\mu} \partial_{\mu} \Phi_{(\mu_1 \nu_1, \mu_2 \nu_2, \ldots, \mu_k \nu_k)} = 0,$$  \hspace{1cm} (2)

$$\eta^{\mu_1 \nu_1}_{\mu_2 \nu_2}(\Phi_{(\mu_1 \nu_1, \mu_2 \nu_2, \ldots, \mu_k \nu_k)}) = \eta^{\mu_2 \nu_2}_{\mu_1 \nu_1}(\Phi_{(\mu_1 \nu_1, \mu_2 \nu_2, \ldots, \mu_k \nu_k)}) = 0,$$  \hspace{1cm} (3)

$$l < m_i, \quad i < j, \quad 1 \leq l_j \leq s_j,$$  \hspace{1cm} (4)

where the bracket below denote that the indices in it do not include in symmetrization.

Simultaneous description of all ISO$(1,d-1)$ group irreps maybe reformulated in a standard manner with an auxiliary Fock space $\mathcal{H}$, generated by $k$ pairs of bosonic creation $a_\mu^i(x)$ and annihilation $a_{\nu}^{ij}(x)$ operators, $i,j = 1, \ldots, k, \mu^i, \nu^i = 0, 1, \ldots, d - 1$: $[a_\mu^i, a_{\nu}^{ij}] = -\eta^{\mu_\nu}_{\nu_\nu} \delta^{ij}$ and a set of constraints for an arbitrary string-like (so called basic) vector $\Phi \in \mathcal{H}$,

$$|\Phi\rangle = \sum_{s_i, i=0}^{s_{k}} \cdots \sum_{s_k = 0}^{s_k} \bigg(\sum_{\Phi^{(1)}} \cdots \sum_{\Phi^{(k)}} \prod_{a=1}^{s} \frac{1}{l_i} \prod_{i=1}^{s} \frac{1}{l_i} \prod_{i=1}^{s} \frac{1}{l_i} \prod_{i=1}^{s} |0\rangle \bigg),$$  \hspace{1cm} (5)

$$(l_{\mu}, l', l^{ij}, l^{ij}) |\Phi\rangle = (\partial^{\mu} \partial_{\mu}, -i\eta^{\mu_\nu}_{\mu_\nu} \frac{1}{2} \partial^{\mu_{s'}} \partial_{\mu_{s'}}, a_\mu^i, a_{\nu}^{ij}) |\Phi\rangle = 0,$$  \hspace{1cm} (6)

The set of $(k(k+1)+1)$ primary constraints (6), $\{|a_\mu^i, a_{\nu}^{ij}\}$, are equivalent to Eqs. (2)–(4) for all spins. In turn, additional condition, $g_0^i |\Phi\rangle = \left(s_i + \frac{d}{2}\right) |\Phi\rangle$ for number particles operators, $g_0^i = -a_\mu^i a_{\nu}^{ij} + \frac{d}{2}$, makes (6) to be equivalent to (2)–(4) for given spin $s$.

The procedure of LF implies the Hermiticity of BFV-BRST operator $Q, Q = \overline{c^a} a^a + \ldots$, that means the extension of the set $\{a_\mu^i\}$ up to one of $\{|a_\mu^i, a_{\nu}^{ij}\}$, which is closed with respect to hermitian conjugation related to standard scalar product on $\mathcal{H}$ and commutator multiplication $. [, ]$. Operators $a_{\nu}^{ij}$ satisfy to the Lie-algebra commutation relations, $[a_{\mu}^i, a_{\nu}^{ij}] = f_{\mu}^{\kappa} a_{\nu}^{ij}$, for structure constants $f_{\mu}^{\kappa} = -f_{\mu}^{\kappa}$, to be determined from the multiplication Table 1.

The products $B_{ij}^{l_{ij}}, A_{ij}^{l_{ij}}, F_{ij}^{l_{ij}}, L_{ij}^{l_{ij}, h}$ in the Table 1 are given by the relations,

$$B_{ij}^{l_{ij}} = (g_0^i - g_0^j) \delta^{ij},$$  \hspace{1cm} (7)

$$A_{ij}^{l_{ij}} = l_{ij}^{ij} \delta^{ij} - l_{ij}^{ij} \delta^{ij},$$  \hspace{1cm} (8)

$$L_{ij}^{l_{ij}, h} = \frac{1}{4} \delta^{ij} \delta^{ij} [2g_0^i \delta^{ij} + g_0^i + g_0^j] - (\delta^{ij} l_{ij}^{ij} \theta^{ij} + l_{ij}^{ij} \delta^{ij} + l_{ij}^{ij} \theta^{ij} + (j_2 \leftrightarrow j_1)),$$  \hspace{1cm} (9)

...
with Heaviside \( \theta \)-symbol \( \theta^\delta \). From the Hamiltonian analysis of the dynamical systems the operators \( \{ o_d \} \) contain \( 2k^2 \) second-class \( \{ l^i_j, I^{i^\delta}_i \} \), \( 2k + 1 \) first-class \( \{ l_0, I^{i^\delta}_i \} \) constraints subsystems and \( k \) elements \( g^i_0 \) forming the non-vanishing in \( H \) matrix

\[
\Delta_{ab} (g^i_0) \text{ in } \{ o_d, g^i_0 \} - \Delta_{ab}.
\]

We called in [4] the algebra of the operators \( O_i \) as integer higher-spin symmetry algebra in Minkowski space with a Young tableaux having \( k \) rows and denoted it as \( \mathcal{A}(Y(k), \mathbb{R}^{1,d-1}) \).

The subsystem of the second-class constraints \( \{ o_d \} \) together with \( \{ g^i_0 \} \) forms the subalgebra in \( \mathcal{A}(Y(k), \mathbb{R}^{1,d-1}) \) to be isomorphic to symplectic \( sp(2k) \) algebra (see details in [4]).

Having constructed the HS symmetry algebra, we can not still construct BRST operator \( Q \) with respect to the elements \( o_d \) from \( \mathcal{A}(Y(k), \mathbb{R}^{1,d-1}) \) due to second-class constraints \( \{ o_d \} \) presence in it. One should to convert symplectic algebra \( sp(2k) \) of \( \{ o_d, g^i_0 \} \) into enlarged set of operators \( O_i \) with only first-class constraints.

3. NEW OSCILLATOR REALIZATION FOR \( sp(2k) \)

Considering an additive conversion procedure developed within BRST approach, (see e.g. [13]), which implies the enlarging of \( o_d \) to \( O_i = o_d + o'_i \), with additional parts \( o'_i \) to be given on a new Fock space \( \mathcal{H}' \).

Now, the elements \( O_i \) are given on \( \mathcal{H} \otimes \mathcal{H}' \) so that a requirement for \( O_i \) \( \{ O_i, O_j \} = O_k \), leads to the same algebraic relations for \( O_i \) and \( o'_i \) as those for \( o_d \).

Leaving aside the details of Verma module (special representation space [12]) construction for the symplectic algebra \( sp(2k) \) of new operators \( o'_i \) considered in [4], we present here theirs explicit oscillator form in terms of new \( 2k^2 \) creation and annihilation operators

\[
(B_i^F; B_i^F) = (b_{i}^\dagger, d_{i}; b_{i}, d_{i}^\dagger), \ i, j, r, s = 1, \ldots, k; \ i \neq j; r < s \text{ as follows for } (k_0 \equiv l)
\]

\[
g^i_0 = \sum_{l \leq m} d_{l}^\dagger d^{l^\dagger} + \sum_{r < s} d_{r}^{s^\dagger} d^{r^\dagger} + h^i, \quad (10)
\]

\[
l^{i^\dagger}_0 = b^{i^\dagger}_0, \quad l^{i^\dagger}_0 = d^{i^\dagger}_m - \sum_{n=1}^{l-1} d_{n} d^{n^\dagger} + \sum_{n=1}^{k} (1 + \delta_{n, 0}) b^{n^\dagger}_m b_{n^\dagger}, \quad (11)
\]

\[
l^{i^\dagger}_0 = \sum_{m=1}^{l-1} d_{m}^\dagger d_{m^\dagger} + \sum_{p=0}^{m-l} \cdots \sum_{k_{s} = l - p}^{m-l} C_{k_{s}}^{p} (d^{s^\dagger}_p, d^{p})^{j}_{j=1} \Delta_{k_{l}, k_{s}} = \sum_{n=1}^{k} (1 + \delta_{n, 0}) b^{n^\dagger}_m b_{n^\dagger}, \quad (12)
\]

\[
Q' = O^i_0 \epsilon^i_0 + \frac{1}{2} \epsilon^i_0 \epsilon^j_0 \int_{\mathcal{H}_k} \delta_{ij} \quad (13)
\]

with the constants \( f_{ij} \) from the Table 1, constraints

\[
\{ O_i, O_j \} = O_k \in (\mathbb{R}^{1,d-1}) \text{ for massless HS fields to } \mathcal{Y}(k), \mathbb{R}^{1,d-1}
\]

\[
\eta^i = \eta^i, \quad \theta^\delta = \theta^\delta, \quad \{ \theta^\delta, \lambda_{\alpha}^\delta \} = \delta_{\alpha}, \delta_{\alpha}, \quad \{ \lambda^\alpha, \lambda_{\alpha}^\delta \} = \delta_{\alpha}, \delta_{\alpha}, \quad (14)
\]

and non-vanishing anticommutators \( \{ \eta^\delta, \eta'_{\eta} \} = 1 \), \( \{ \eta^\delta_{\eta}, \eta'_{\eta} \} = i \delta^\delta_{\eta} \) for zero-mode ghosts.

To construct LF for bosonic HS fields in a \( \mathbb{R}^{1,d-1} \) Minkowski space we partially follow the algorithm of [13, 14], which is a particular case of our construction, corresponding to \( s_3 = 0 \). First, we extract the depen-

\[2\] The case of the massive bosonic HS fields whose system of second-class constraints contains additionally to elements of \( sp(2k) \) algebra the constraints of isometry subalgebra of Minkowski space \( l^i_0, l^i_0, l^i_0 \) may be treated via procedure of dimensional reduction of the algebra \( \mathcal{A}(Y(k), \mathbb{R}^{1,d-1}) \) for massive HS fields to one \( \mathcal{A}(Y(k), \mathbb{R}^{1,d-1}) \) for massive HS fields, (see [4]). Now, the wave equation in (2) is changed on Klein-Gordon equation corresponding to the constraint \( l^i_0 (l^0_0 = \epsilon^\dagger_0 c_{\mu}^\dagger + m^2) \) acting on the same basic vector \( |\Phi \rangle \) (5).

\[3\] The ghosts possess the standard ghost number distribution, \( gh(\epsilon^i_0) = -gh(\epsilon^j_0) = 1 \Rightarrow gh(Q') = 1 \).
dence of $Q'$ (13) on the ghosts $\eta_i^G, \mathcal{P}_i^G$, to obtain the BRST operator $Q$ only for the system of converted first-class constraints $\{O_j\} \setminus \{G_0\}$ and generalized spin operator $\sigma'$:

$$Q' = Q + \eta_i^G (\sigma' + h') + \mathcal{A} \mathcal{P}_i^G,$$  \hspace{1cm} (15)

with some $\mathcal{A}'$, where

$$Q = \left( \frac{1}{2} \eta_0 L_0 + \sum_{i=1}^{\infty} (\eta_i L_i + \sum_{m \leq n} \mathcal{P}_i^m + \mathcal{P}_i^m + h_c) \right) + \frac{1}{2} \mathcal{E}_i \mathcal{E}_j f_{ij} \mathcal{P}_k,$$

$$\sigma_i = G_i^0 - h_i - \sum_{m} \mathcal{P}_i^m - \eta_i \mathcal{P}_i + \sum_{m} (1 + \delta_{im}) \eta_i \mathcal{P}_i^m + \sum_{i \leq j} \left[ \delta_{ij} \lambda_i \phi_j - \mathcal{P}_i^m \lambda_j \phi_i - \mathcal{P}_i^m \phi_i \lambda_j \right]$$  \hspace{1cm} (16)

where \( \{ \mathcal{E}_i \} = \{ \mathcal{E}_j \} \{ \mathcal{E}_k \}, \{ \mathcal{P}_i^m \} \equiv \{ \mathcal{P}_j^m \} \{ \mathcal{P}_k^m \} \). Then, we choose a representation of $\mathcal{H}_{\text{tot}}$: \( \eta_i, \mathcal{P}_i^m, \lambda_i, \phi_i, \lambda_i, \phi_i \) \( \{ \{ \mathcal{P}_i^m \} \} \{ \{ \mathcal{P}_i^m \} \} \{ \{ \mathcal{P}_i^m \} \} \). Thus, the physical state having the ghost number \( \eta_i \) do not depend on ghosts $\eta_i$:

$$|\chi\rangle = \sum_{n, n_i} \prod_{i \leq j} \left( \eta_i^+ \right)^{n_i} \left( \mathcal{P}_i^m \right)^{n_i} \left( \mathcal{P}_i^m \right)^{n_i} \Phi = \Phi, \Phi = \Phi \times \Phi,$$

$$\times \prod_{i \leq j \leq m} \left( \lambda_i^+ \right)^{n_i} \left( \phi_i^+ \right)^{n_i} \left( \phi_i^+ \right)^{n_i} \Phi,$$

$$\times \prod_{r < s \leq m} \left( \lambda_i^+ \right)^{n_i} \left( \phi_i^+ \right)^{n_i} \left( \phi_i^+ \right)^{n_i} \Phi, \Phi = \Phi \times \Phi.$$

We denote by $|\chi_0\rangle$ the state (18) satisfying to $gh(|\chi_0\rangle) = -|\chi_0\rangle$. Thus, the physical state having the ghost number zero is $|\chi_0\rangle$, the gauge parameters $|\Lambda\rangle$ having the ghost number $-1$ is $|\chi_0\rangle$ and so on. The vector $|\chi_0\rangle$ must contain physical string-like vector $|\Phi\rangle = |\Phi\rangle$:

$$|\chi_0\rangle = |\Phi\rangle + |\Phi\rangle, \quad \text{where} \quad |\Phi\rangle = \Phi \times \Phi, \Phi = \Phi \times \Phi.$$

Independence of the vectors (18) on $\eta_i^G$ transforms the equation for the physical state $Q|\chi_0\rangle = 0$ and the BRST complex of the reducible gauge transformations, $\delta |\chi\rangle = Q|\chi_0\rangle = 0 |\chi_0\rangle$, $Q|\chi_0\rangle = 0 |\chi_0\rangle$, ..., $Q|\chi_0\rangle = 0 |\chi_0\rangle$, to the relations:

$$Q|\chi_0\rangle = 0, \quad \delta |\chi_0\rangle = Q|\chi_0\rangle = 0 |\chi_0\rangle, \quad \delta |\chi_0\rangle = Q|\chi_0\rangle = 0,$$

$$Q|\chi_0\rangle = 0, \quad (\sigma' + h')|\chi_0\rangle = 0, \quad (\sigma' + h')|\chi_0\rangle = 0,$$

$$Q|\chi_0\rangle = 0, \quad (\sigma' + h')|\chi_0\rangle = 0, \quad (\sigma' + h')|\chi_0\rangle = 0.$$  \hspace{1cm} (20)

let us fix some values of $n_i = s_i$. Then one should substitute $h_i'$ corresponding to the chosen $s_i$ (22) into $Q$ (16) and relations (20). Thus, e.g., the equation of motion (20) corresponding to the field with given spin $(s_1, ..., s_k)$ has the form $Q|\chi_0\rangle = 0$, with nilpotent $Q|\chi_0\rangle$ and the same for the rest relations in (20).

Following to bosonic one—[12] and two-row cases, [13, 14] one can show that last equation may be derived from the Lagrangian action for fixed spin $(n)_k = (s)_k$,

$$S_{(n)} = \int d\eta^{(n)} \left\langle \chi_0 \right| K_{(n)} Q_{(n)} \left| \chi_0 \right\rangle,$$

$$\frac{\delta \mathcal{S}_{(n)}}{\delta \eta^{(n)} | \chi_0 \rangle} = 0 \Rightarrow Q_{(n)} \left| \chi_0 \right\rangle = 0,$$

where the standard scalar product for the creation and annihilation operators in $\mathcal{H}_{\text{tot}} = \mathcal{H} \otimes \mathcal{H} \otimes \mathcal{H}$ is assumed and non-degenerate operator $K_{(n)}$ provides reality of the action.

Concluding, one can prove the action (23) indeed reproduces the basic conditions (20) (4) for massless (massive) HS fields. General action (23) gives, in principle, a straight recept to obtain the Lagrangian for any component field from general vector $|\chi_0\rangle$.

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