

Unfolded equations for massive higher spin supermultiplets in AdS_3

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ABSTRACT: In this paper we give an explicit construction of unfolded equations for massive higher spin supermultiplets of the minimal $(1, 0)$ supersymmetry in AdS_3 space. For that purpose we use an unfolded formulation for massive bosonic and fermionic higher spins and find supertransformations leaving appropriate set of unfolded equations invariant. We provide two general supermultiplets $(s, s + 1/2)$ and $(s, s - 1/2)$ with arbitrary integer s , as well as a number of lower spin examples.

KEYWORDS: Field Theories in Lower Dimensions, Higher Spin Symmetry, Supersymmetric gauge theory

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1 Introduction

Among all the higher spin symmetries a supersymmetry still plays a distinguished role. It is enough to recall that all massive states in the superstring theory are nicely organized into massive higher spin supermultiplets. The classification of massless and massive supermultiplets (depending on the space-time dimensions, number of supersymmetries and type of fermions) is purely algebraic task and is very well understood now. But as far as the

concrete realization in terms of Lagrangians and/or equations of motion is concerned there is a striking difference between massless and massive cases. For the massless supermultiplets the realization can be straightforwardly constructed (both in components as well as in superfields) simply because the supertransformations always have the same simple pattern:

$$\delta B \sim F\eta, \quad \delta F \sim \partial B\eta$$

where B and F are bosonic and fermionic fields and η — parameter of the supertransformations. But switching from the massless to the massive case one has to introduce a lot of complicated higher derivatives corrections to the supertransformations without any evident pattern and the higher the spin of fields entering the supermultiplet the higher the number of derivatives one has to consider.

In the four-dimensional Minkowski space the solution in components was proposed in [1] based on the gauge invariant description of massive higher spin bosonic [2] and fermionic [3] fields.¹ The main idea was that the massive supermultiplet can be constructed out of the appropriate set of massless ones in the same way as the gauge invariant description of massive higher spin particles can be constructed using an appropriate set of massless ones. In spite of the large number of fields involved such construction appears to be pretty straightforward. At the same time the meaning of these complicated corrections to the supertransformations that one has to introduce working with non-gauge invariant description of massive particles becomes clear. Namely, they are just the restoring gauge transformations that appear when one tries to exclude all Stueckelberg fields fixing all gauge symmetries.

Recently, using the same approach, we have constructed an explicit Lagrangian formulation for massive higher spin supermultiplets in three-dimensional Minkowski space [19] based on our previous works on the gauge invariant formulation for massive bosonic [20] and fermionic [21] higher spins in three dimensions.² The aim of the current work is to investigate massive higher spin supermultiplets in the three-dimensional anti de Sitter space. Note that the approach under consideration leads to on-shell supersymmetric models where the auxiliary fields allowing to close the supersymmetry algebra are absent. As is well known [23], AdS_3 space is special because all AdS_3 superalgebras (as well as AdS_3 algebra itself) factorize into “left” and “right” parts. For the case of simplest (1,0) superalgebra (the one we are working here with) it has the form:

$$OSp(1, 2) \otimes Sp(2)$$

so that we have supersymmetry in the “left” sector only. It means that the minimal massive supermultiplet must contain just one bosonic and one fermionic degrees of freedom. To realize such supermultiplets we will use a so called unfolded formalism [24, 25] which

¹Construction of the Lagrangian superfield description for massive superspins 1 and 3/2 has been initiated by the papers [4] and [5] (see also [6–11]) while analogous description for massive supermultiplets with arbitrary superspin is absent up to now. Lagrangian description for massless higher superspin theories is developed much better [12–18].

²Superfield approach to three dimensional higher spin supersymmetric models proposed in [22].

apparently is a most simple and efficient way.³ For the dimensions $d \geq 4$ complete unfolded description of massive bosonic higher spins has been constructed in [26]. In three dimensions the authors of [27] suggested unfolded equations for the infinite chain of gauge invariant zero-forms and showed that such system can describe different representations of AdS_3 algebra such as massive, topologically massive, fractional spins and so on. Recently one of us has shown the relation of such unfolded equations with the Lagrangian formulation for massive bosonic fields [28]. We begin with the frame-like gauge invariant Lagrangian for the massive bosonic higher spin s [20] that includes a set of one-forms $\Omega^{\alpha(2k)}$, $1 \leq k \leq s - 1$ and zero-form $B^{\alpha(2)}$ (here and in what follows we use a multispinor formalism, see below for notations and conventions). Then it appears that to construct gauge invariant unfolded equations one has to introduce a set of non gauge invariant zero forms $B^{\alpha(2k)}$, $2 \leq k \leq s - 1$ playing the roles of Stueckelberg fields. At last the whole system is constructed by addition of the infinite chain of gauge invariant zero-forms $B^{\alpha(2k)}$, $k \geq s$ whose equations are in agreement with the results of [27]. Thus the complete system contains the following set of fields:

$$\Omega^{\alpha(2k)}, \quad B^{\alpha(2k)} \quad 1 \leq k \leq s - 1, \quad B^{\alpha(2k)} \quad k \geq s.$$

Using our previous results on the frame-like gauge invariant Lagrangian description of the massive spin $s + 1/2$ fermionic field [21] one can construct (see appendices C, D) an unfolded formulation that also includes a set of gauge and Stueckelberg fields as well as infinite number of gauge invariant ones:

$$\Phi^{\alpha(2k+1)}, \quad \phi^{\alpha(2k+1)} \quad 0 \leq k \leq s - 1, \quad \phi^{\alpha(2k+1)} \quad k \geq s.$$

Having in our disposal the unfolded equations for the massive bosonic and fermionic fields we look for the supertransformations leaving these equations invariant. Namely, following the example of the massless scalar supermultiplet $(1/2, 0)$ in $d = 4$ [29] (see also [30]), we consider quadratic deformations of the unfolded equations that have the form (schematically):

$$0 = DB \oplus B \oplus F\Psi^\alpha, \quad 0 = DF \oplus F \oplus B\Psi^\alpha$$

where Ψ^α is the massless spin-3/2 gravitino. The requirement that such deformations to be consistent determines all the arbitrary coefficients. After that explicit expressions for the supertransformations can be easily extracted and in the unfolded formalism they have purely algebraic form:

$$\delta B \sim F\zeta^\alpha, \quad \delta F \sim B\zeta^\alpha.$$

The paper is organized as follows. For completeness and comparison in section 2 we consider two massless supermultiplets $(s, s + 1/2)$ and $(s, s - 1/2)$ with arbitrary integer s . In-particular, these examples clearly show specific properties of $d = 3$ case related with the factorization of AdS_3 superalgebra. We want to emphasize that the supermultiplets

³In principle the description of such supermultiplets can be done analogous to [19–21]. However the consideration turns out to be more complicated.

$(s, s + 1/2)$ and $(s, s - 1/2)$ are essentially different. In the first case the higher spin of the multiplet is half-integer, while in the second case the higher spin is integer. Therefore these two supermultiplets should be investigated separately. Section 3 contains a number of concrete (relatively) low spin examples of massive supermultiplets. Here we consider $(1, 1/2)$, $(1/2, 0)$, $(3/2, 1)$ and $(2, 3/2)$ ones. The main part of the paper contains sections 4 and 5 where we consider massive arbitrary spin supermultiplets $(s, s + 1/2)$ and $(s, s - 1/2)$. The paper contains also four appendices devoted to the unfolded equations that are necessary for the main part. Appendix A gives unfolded equations for the $s = 0, 1/2, 1, 3/2, 2$, appendix B — for the massive spin- s boson, while appendices C and D — for the massive spins $s + 1/2$ and $s - 1/2$ fermions correspondingly.

Notations and conventions. We use a frame-like multispinor formalism where all objects (one-forms or zero-forms) have local indices which are completely symmetric spinor ones. To simplify expressions we will use condensed notations for the spinor indices such that e.g.

$$\Omega^{\alpha(2k)} = \Omega^{(\alpha_1 \alpha_2 \dots \alpha_{2k})}.$$

Also we will always assume that spinor indices denoted by the same letter and placed on the same level are symmetrized, e.g.

$$\Omega^{\alpha(2k)} \zeta^\alpha = \Omega^{(\alpha_1 \dots \alpha_{2k} \zeta^{\alpha_{2k+1}})}.$$

AdS_3 space will be described by the background frame (one-form) $e^{\alpha(2)}$ and the covariant derivative D normalized so that

$$D \wedge D \zeta^\alpha = -\lambda^2 E^\alpha{}_\beta \zeta^\beta$$

where two-form $E^{\alpha(2)}$ is defined as follows:

$$e^{\alpha(2)} \wedge e^{\beta(2)} = \varepsilon^{\alpha\beta} E^{\alpha\beta}.$$

In what follows the wedge product sign \wedge will be omitted.

2 Massless supermultiplets

In this section we present an explicit construction for the supermultiplets with massless higher spin fields.

2.1 Kinematics

The description of the massless bosonic spin- s field in the frame-like multispinor formalism requires a pair of one-forms $\Omega^{\alpha(2s-2)}$ and $f^{\alpha(2s-2)}$. The free Lagrangian (that is three-form in our formalism) describing such field living in AdS_3 space has the form:

$$\begin{aligned} \mathcal{L}_0 = (-1)^s & \left[(s-1) \Omega_{\alpha(2s-3)\beta} e^\beta{}_\gamma \Omega^{\alpha(2s-3)\gamma} + \Omega_{\alpha(2s-2)} D f^{\alpha(2s-2)} \right. \\ & \left. + \frac{(s-1)\lambda^2}{4} f_{\alpha(2s-3)\beta} e^\beta{}_\gamma f^{\alpha(2s-3)\gamma} \right]. \end{aligned} \tag{2.1}$$

This Lagrangian is invariant under the following local gauge transformations:

$$\begin{aligned}\delta\Omega^{\alpha(2s-2)} &= D\eta^{\alpha(2s-2)} + \frac{\lambda^2}{4}e^\alpha{}_\beta\xi^{\alpha(2s-3)\beta} \\ \delta f^{\alpha(2s-2)} &= D\xi^{\alpha(2s-2)} + e^\alpha{}_\beta\eta^{\alpha(2s-3)\beta}\end{aligned}\tag{2.2}$$

where η and ξ are zero-forms completely symmetric in their local indices.

Let us introduce new variables:

$$\begin{aligned}\hat{\Omega}^{\alpha(2s-2)} &= \Omega^{\alpha(2s-2)} + \frac{\lambda}{2}f^{\alpha(2s-2)} \\ \hat{f}^{\alpha(2s-2)} &= \Omega^{\alpha(2s-2)} - \frac{\lambda}{2}f^{\alpha(2s-2)}\end{aligned}\tag{2.3}$$

and similarly for the parameters of the gauge transformations:

$$\begin{aligned}\hat{\eta}^{\alpha(2s-2)} &= \eta^{\alpha(2s-2)} + \frac{\lambda}{2}\xi^{\alpha(2s-2)} \\ \hat{\xi}^{\alpha(2s-2)} &= \eta^{\alpha(2s-2)} - \frac{\lambda}{2}\xi^{\alpha(2s-2)}.\end{aligned}\tag{2.4}$$

Then the Lagrangian can be rewritten as the sum of the two independent parts:

$$\begin{aligned}\mathcal{L}_0 &= \frac{(-1)^s}{2\lambda}[(s-1)\lambda\hat{\Omega}_{\alpha(2s-3)\beta}e^\beta{}_\gamma\hat{\Omega}^{\alpha(2s-3)\gamma} + \hat{\Omega}_{\alpha(2s-2)}D\hat{\Omega}^{\alpha(2s-2)} \\ &\quad + (s-1)\lambda\hat{f}_{\alpha(2s-3)\beta}e^\beta{}_\gamma\hat{f}^{\alpha(2s-3)\gamma} - \hat{f}_{\alpha(2s-2)}D\hat{f}^{\alpha(s-2)}]\end{aligned}\tag{2.5}$$

while the gauge transformations take the form:

$$\begin{aligned}\delta\hat{\Omega}^{\alpha(2s-2)} &= D\hat{\eta}^{\alpha(2s-2)} + \frac{\lambda}{2}e^\alpha{}_\beta\hat{\eta}^{\alpha(2s-3)\beta} \\ \delta\hat{f}^{\alpha(2s-2)} &= D\hat{\xi}^{\alpha(2s-2)} - \frac{\lambda}{2}e^\alpha{}_\beta\hat{\xi}^{\alpha(2s-3)\beta}.\end{aligned}\tag{2.6}$$

Free Lagrangian for the massless fermionic field with spin s has the form:

$$\mathcal{L}_0 = \frac{i}{2}(-1)^{s-1/2}[\Phi_{\alpha(2s-2)}D\Phi^{\alpha(2s-2)} + (s-1)\lambda\Phi_{\alpha(2s-3)\beta}e^\beta{}_\gamma\Phi^{\alpha(2s-3)\gamma}].\tag{2.7}$$

It is invariant under the following local gauge transformations:

$$\delta\Phi^{\alpha(2s-2)} = D\xi^{\alpha(2s-2)} + \frac{\lambda}{2}e^\alpha{}_\beta\xi^{\alpha(2s-3)\beta}.\tag{2.8}$$

2.2 Supermultiplet $(s, s + 1/2)$

The free Lagrangian now is the sum of the free Lagrangians for the bosonic and fermionic fields:

$$\mathcal{L}_0 = \mathcal{L}_0(\Omega^{\alpha(2s-2)}, f^{\alpha(2s-2)}) + \mathcal{L}_0(\Phi^{\alpha(2s-1)}).\tag{2.9}$$

Let us consider the following ansatz for the global supertransformations:

$$\begin{aligned}\delta\Omega^{\alpha(2s-2)} &= i\alpha_1\Phi^{\alpha(2s-2)\beta}\zeta_\beta \\ \delta f^{\alpha(2s-2)} &= i\alpha_2\Phi^{\alpha(2s-2)\beta}\zeta_\beta \\ \delta\Phi^{\alpha(2s-1)} &= \beta_1\Omega^{\alpha(2s-2)}\zeta^\alpha + \beta_2f^{\alpha(2s-2)}\zeta^\alpha.\end{aligned}\tag{2.10}$$

Recall that in AdS_3 space by global supertransformations we mean the ones with the parameters satisfying the relation

$$D\zeta^\alpha = -\frac{\lambda}{2}e^\alpha{}_\beta\zeta^\beta. \quad (2.11)$$

Invariance of the Lagrangian requires:

$$\alpha_1 = \frac{\lambda}{2}(2s-1)\beta_1, \quad \alpha_2 = (2s-1)\beta_1, \quad \beta_2 = \frac{\lambda}{2}\beta_1.$$

In terms of hatted variables it gives:

$$\begin{aligned} \delta\hat{\Omega}^{\alpha(2s-2)} &= i(2s-1)\lambda\beta_1\Phi^{\alpha(2s-2)\beta}\zeta_\beta, & \delta\hat{f}^{\alpha(2s-2)} &= 0 \\ \delta\Phi^{\alpha(2s-1)} &= \beta_1\hat{\Omega}^{\alpha(2s-2)}\zeta^\beta. \end{aligned} \quad (2.12)$$

2.3 Supermultiplet $(s, s - 1/2)$

In this case the free Lagrangian has the form:

$$\mathcal{L}_0 = \mathcal{L}_0(\Omega^{\alpha(2s-2)}, f^{\alpha(2s-2)}) + \mathcal{L}_0(\Phi^{\alpha(2s-3)}). \quad (2.13)$$

Let us consider the following ansatz for the supertransformations:

$$\begin{aligned} \delta\Omega^{\alpha(2s-2)} &= i\alpha_1\Phi^{\alpha(2s-3)}\zeta^\alpha \\ \delta f^{\alpha(2s-2)} &= i\alpha_2\Phi^{\alpha(2s-3)}\zeta^\alpha \\ \delta\Phi^{\alpha(2s-3)} &= \beta_1\Omega^{\alpha(2s-3)\beta}\zeta_\beta + \beta_2f^{\alpha(2s-3)\beta}\zeta_\beta. \end{aligned} \quad (2.14)$$

Invariance of the Lagrangian requires:

$$\alpha_1 = \frac{\lambda}{2}\alpha_2, \quad \beta_1 = 2(s-1)\alpha_2, \quad \beta_2 = 2(s-1)\frac{\lambda}{2}\alpha_2.$$

In terms of hatted variables it gives:

$$\begin{aligned} \delta\hat{\Omega}^{\alpha(2s-2)} &= i\lambda\alpha_2\Phi^{\alpha(2s-3)}\zeta^\alpha, & \delta\hat{f}^{\alpha(2s-2)} &= 0 \\ \delta\Phi^{\alpha(2s-3)} &= 2(s-1)\alpha_2\hat{\Omega}^{\alpha(2s-3)\beta}\zeta_\beta. \end{aligned} \quad (2.15)$$

In both cases the results are consistent with the factorization of superalgebra in AdS_3 :

$$\mathcal{A} \sim \text{OSp}(1,2) \otimes \text{Sp}(2).$$

3 Low spin examples

In this section we consider examples of the massive low spins supermultiplets with $0 \leq s \leq 2$. Unfolded equations for all these fields are given in appendix A.

3.1 Supermultiplet (1, 1/2)

The unfolded formulation of this supermultiplet requires bosonic zero-forms $B^{\alpha(2k)}$, $k \geq 1$ as well as fermionic ones $\phi^{\alpha(2k+1)}$, $k \geq 0$. As it was already explained in the Introduction, our general procedure is to consider the deformation of the initial unfolded equations corresponding to switching on a background gravitino field Ψ^α satisfying the relation

$$D\Psi^\alpha = -\frac{\lambda}{2}e^\alpha{}_\beta\Psi^\beta$$

and to require that deformed equations remain to be consistent. This in turn allows one easily extract the explicit form of the global supertransformations leaving unfolded equations invariant.

Let us begin with the deformations for the bosonic equations:

$$0 = DB^{\alpha(2k)} - e_{\beta(2)}B^{\alpha(2k)\beta(2)} - A_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)}B^{\alpha(2k-2)} - E_k \phi^{\alpha(2k)\beta}\Psi_\beta - F_k \phi^{\alpha(2k-1)}\Psi^\alpha. \quad (3.1)$$

Their consistency in the linear approximation (i.e. taking into account quadratic terms in the consistency condition only) requires:

$$E_k = E_1, \quad F_k = -\frac{1}{2} \left[(2k-1)C_{k-1} - 2(k-1)A_{k-1} - \frac{\lambda}{2} \right] E_1. \quad (3.2)$$

For the A_k and C_k corresponding to spin-1 and spin-1/2 (see appendix A) this gives

$$m = m_1 + \frac{\lambda}{2}, \quad F_k = -\frac{(k+1)}{2k(2k+1)}[m_1 - k\lambda]E_1. \quad (3.3)$$

Similarly, deformations for the fermionic equations have the form:

$$0 = D\phi^{\alpha(2k+1)} - e_{\beta(2)}\phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)}\phi^{\alpha(2k-1)} - G_k B^{\alpha(2k+1)\beta}\Psi_\beta - H_k B^{\alpha(2k)}\Psi^\alpha. \quad (3.4)$$

Their consistency requires:

$$G_k = G_0, \quad H_k = -\frac{1}{2} \left[2kA_k - (2k-1)C_{k-1} - \frac{\lambda}{2} \right] G_0 \quad (3.5)$$

and gives the same relation on masses. An explicit expression for H_k looks like:

$$H_k = \frac{k}{2(k+1)(2k+1)}[m_1 + (k+1)\lambda]G_0. \quad (3.6)$$

Thus we have found the supertransformations leaving unfolded equations for massive spin-1 and spin-1/2 invariant:

$$\begin{aligned} \delta B^{\alpha(2k)} &= E_1 \phi^{\alpha(2k)\beta} \zeta_\beta + F_k \phi^{\alpha(2k-1)} \zeta^\alpha \\ \delta \phi^{\alpha(2k+1)} &= G_0 B^{\alpha(2k+1)\beta} \zeta_\beta + H_k B^{\alpha(2k)} \zeta^\alpha. \end{aligned} \quad (3.7)$$

In this, we have two arbitrary constants E_1 and G_0 . Calculating the commutator of these supertransformations one can fix $E_1 G_0$ as a normalization for the superalgebra. But to fix relative values for E_1 and G_0 one has to construct appropriate Lagrangian formalism.

3.2 Supermultiplet (1/2, 0)

In this case we need the bosonic zero-forms $\pi^{\alpha(2k)}$ and the fermionic ones $\phi^{\alpha(2k+1)}$, $k \geq 0$. Deformations for their unfolded equations:

$$\begin{aligned} 0 &= D\pi^{\alpha(2k)} - e_{\beta(2)}\pi^{\alpha(2k)\beta(2)} - B_k e^{\alpha(2)}\pi^{\alpha(2k-2)} - E_k \phi^{\alpha(2k)\beta} \Psi_\beta - F_k \phi^{\alpha(2k-1)} \Psi^\alpha \\ 0 &= D\phi^{\alpha(2k+1)} - e_{\beta(2)}\phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)}\phi^{\alpha(2k-1)} \\ &\quad - G_k \pi^{\alpha(2k+1)\beta} \Psi_\beta - H_k \pi^{\alpha(2k)} \Psi^\alpha \end{aligned} \quad (3.8)$$

as well as all calculations are the same as in the previous case except that now all $A_k = 0$. We obtain:

$$\begin{aligned} m_0^2 &= m^2 - m\lambda - \frac{3}{4}\lambda^2 \\ F_k &= \frac{1}{2(2k+1)} \left[m + (2k+1)\frac{\lambda}{2} \right] E_0 \\ H_k &= -\frac{1}{2(2k+1)} \left[m - (2k+1)\frac{\lambda}{2} \right] G_0. \end{aligned} \quad (3.9)$$

The supertransformations also have the same form:

$$\begin{aligned} \delta\pi^{\alpha(2k)} &= E_0 \phi^{\alpha(2k)\beta} \zeta_\beta + F_k \phi^{\alpha(2k-1)} \zeta^\alpha \\ \delta\phi^{\alpha(2k+1)} &= G_0 \pi^{\alpha(2k+1)\beta} \zeta_\beta + H_k \pi^{\alpha(2k)} \zeta^\alpha. \end{aligned} \quad (3.10)$$

3.3 Supermultiplet (3/2, 1)

Unfolded formulation for the massive spin-3/2 requires one-form Φ^α , Stueckelberg zero-form ϕ^α as well as a number of gauge invariant zero-forms $\phi^{\alpha(2k+1)}$, $k \geq 1$. Let us consider the following deformations for the fermionic equations:

$$\begin{aligned} 0 &= D\Phi^\alpha + M e^\alpha{}_\beta \Phi^\beta + 2m E^\alpha{}_\beta \phi^\beta - \alpha_1 e_{\beta(2)} B^{\beta(2)} \Psi^\alpha \\ 0 &= D\phi^\alpha + 2m \Phi^\alpha + M e^\alpha{}_\beta \phi^\beta - e_{\beta(2)} \phi^{\alpha\beta(2)} - G_0 B^{\alpha\beta} \Psi_\beta \\ 0 &= D\phi^{\alpha(2k+1)} - e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)} \phi^{\alpha(2k-1)} \\ &\quad - G_k B^{\alpha(2k+1)\beta} \Psi_\beta - H_k B^{\alpha(2k)} \Psi^\alpha. \end{aligned} \quad (3.11)$$

Consistency of the first equation requires:

$$\alpha_1 = \frac{mG_0}{(2M + \lambda)}, \quad M = m_1 - \frac{\lambda}{2} \quad (3.12)$$

while the consistence of the remaining equations gives:

$$G_k = G_0, \quad H_k = -\frac{(k+2)}{2(k+1)(2k+1)} [m_1 - (k+1)\lambda] G_0. \quad (3.13)$$

For the massive spin-1 unfolded equations and their deformations have the same form as before:

$$\begin{aligned} 0 &= DB^{\alpha(2k)} - e_{\beta(2)} B^{\alpha(2k)\beta(2)} - A_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)} B^{\alpha(2k-2)} \\ &\quad - E_k \phi^{\alpha(2k)\beta} \Psi_\beta - F_k \phi^{\alpha(2k-1)} \Psi^\alpha. \end{aligned} \quad (3.14)$$

Their consistency gives the same mass relation and

$$E_k = E_1, \quad F_k = \frac{(k-1)}{2k(2k+1)}[m_1 + k\lambda]E_1. \quad (3.15)$$

Thus, the supertransformations have the form ($k \geq 1$):

$$\begin{aligned} \delta\Phi^\alpha &= -\alpha_1 e_{\beta(2)} B^{\beta(2)} \zeta^\alpha, & \delta\phi^\alpha &= G_0 B^{\alpha\beta} \zeta_\beta \\ \delta\phi^{\alpha(2k+1)} &= G_0 B^{\alpha(2k+1)\beta} \zeta_\beta + H_k B^{\alpha(2k)} \zeta^\alpha \\ \delta B^{\alpha(2k)} &= E_1 \phi^{\alpha(2k)\beta} \zeta_\beta + F_k \phi^{\alpha(2k-1)} \zeta^\alpha \end{aligned} \quad (3.16)$$

and also contain two arbitrary constants E_1 and G_0 .

3.4 Supermultiplet (2, 3/2)

Unfolded formulation for the massive spin-2 requires one-form $\Omega^{\alpha(2)}$, Stueckelberg zero-form $B^{\alpha(2)}$ and the number of gauge invariant zero-forms $B^{\alpha(2k)}$, $k \geq 1$. Deformations for the bosonic equations we take in the following form:

$$\begin{aligned} 0 &= D\Omega^{\alpha(2)} + m_2 E^\alpha{}_\beta B^{\alpha\beta} + \frac{M_2}{2} e^\alpha{}_\beta \Omega^{\alpha\beta} - \beta_1 \Phi^\alpha \Psi^\alpha - \beta_2 e^{\alpha(2)} \phi_\beta \Psi^\beta \\ 0 &= DB^{\alpha(2)} + m_2 \Omega^{\alpha(2)} + \frac{M_2}{2} e^\alpha{}_\beta B^{\alpha\beta} - e_{\beta(2)} B^{\alpha(2)\beta(2)} - E_1 \phi^{\alpha(2)\beta} \Psi_\beta - F_1 \phi^\alpha \Psi^\alpha \\ 0 &= DB^{\alpha(2k)} - e_{\beta(2)} B^{\alpha(2k)\beta(2)} - A_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)} B^{\alpha(2k-2)} \\ &\quad - E_k \phi^{\alpha(2k)\beta} \Psi_\beta - F_k \phi^{\alpha(2k-1)} \Psi^\alpha. \end{aligned} \quad (3.17)$$

The consistency of the first equation leads to

$$M_2 = M_1 - \frac{\lambda}{2}, \quad \beta_1 = \frac{m_1 m_2}{(M_2 + \lambda)} E_1, \quad \beta_2 = -\frac{m_2}{2} E_1 \quad (3.18)$$

while the consistency of the remaining equations produces

$$E_k = E_1, \quad F_k = -\frac{(k+2)}{2k(2k+1)}[M_2 - k\lambda]E_1. \quad (3.19)$$

For the massive spin-3/2 we need the same set of fields as before while the deformations for their equations look like:

$$\begin{aligned} 0 &= D\Phi^\alpha + M_1 e^\alpha{}_\beta \Phi^\beta + 2m_1 E^\alpha{}_\beta \phi^\beta - \alpha_1 \Omega^{\alpha\beta} \Psi_\beta - \alpha_2 e_{\beta(2)} B^{\beta(2)} \Psi^\alpha \\ 0 &= D\phi^\alpha + 2m_1 \Phi^\alpha + M_1 e^\alpha{}_\beta \phi^\beta - e_{\beta(2)} \phi^{\alpha\beta(2)} - G_0 B^{\alpha\beta} \Psi_\beta \\ 0 &= D\phi^{\alpha(2k+1)} - e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)} \phi^{\alpha(2k-1)} \\ &\quad - G_k B^{\alpha(2k+1)\beta} \Psi_\beta - H_k B^{\alpha(2k)} \Psi^\alpha. \end{aligned} \quad (3.20)$$

Consistency of the first one requires

$$\alpha_1 = -\frac{m_2}{2m_1} G_0, \quad \alpha_2 = \frac{m_1}{2(2M_1 - \lambda)} G_0 \quad (3.21)$$

while the consistency of the remaining ones leads to

$$G_k = G_0, \quad H_k = \frac{(k-1)}{2(k+1)(2k+1)}[M_2 + (k+1)\lambda]G_0. \quad (3.22)$$

The complete set of the supertransformations has the form:

$$\begin{aligned} \delta\Omega^{\alpha(2)} &= -\beta_1\Phi^\alpha\zeta^\alpha - \beta_2e^{\alpha(2)}\phi_\beta\zeta^\beta \\ \delta B^{\alpha(2k)} &= E_1\phi^{\alpha(2k)\beta}\zeta_\beta + F_k\phi^{\alpha(2k-1)}\zeta^\alpha \\ \delta\Phi^\alpha &= -\alpha_1\Omega^{\alpha\beta}\zeta_\beta - \alpha_2e_{\beta(2)}B^{\beta(2)}\zeta^\alpha \\ \delta\phi^{\alpha(2k+1)} &= G_0B^{\alpha(2k)\beta}\zeta_\beta + H_kB^{\alpha(2k)}\zeta^\alpha. \end{aligned} \quad (3.23)$$

4 Supermultiplet $(s, s + 1/2)$

The unfolded formulation for the bosonic spin- s field (appendix B) requires a set of one-forms $\Omega^{\alpha(2k)}$ and Stueckelberg zero-forms $B^{\alpha(2k)}$, $1 \leq k \leq s-1$ as well as an infinite number of gauge invariant zero-forms $B^{\alpha(2k)}$, $k \geq s$. Similarly, the unfolded formulation for the fermionic spin $s + 1/2$ field (appendix C) requires a set of one-forms $\Phi^{\alpha(2k+1)}$ and Stueckelberg zero-forms $\phi^{\alpha(2k+1)}$, $0 \leq k \leq s-1$ as well as an infinite number of gauge invariant zero-forms $\phi^{\alpha(2k+1)}$, $k \geq s$.

Let us begin with the deformations of the equations for the gauge invariant fermionic zero-forms ($k \geq s$):

$$\begin{aligned} 0 &= D\phi^{\alpha(2k+1)} - e_{\beta(2)}\phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)}\phi^{\alpha(2k-1)} \\ &\quad - G_k B^{\alpha(2k+1)\beta}\Psi_\beta - H_k B^{\alpha(2k)}\Psi^\alpha. \end{aligned} \quad (4.1)$$

Their consistency gives the mass relation:

$$M_1 = M_2 - \frac{\lambda}{2} \quad (4.2)$$

and fixes all coefficients in terms of one constant G_s :

$$G_k = G_s, \quad H_k = -\frac{(s+k+1)}{2(k+1)(2k+1)}[M_2 - (k+1)\lambda]G_s, \quad k \geq s. \quad (4.3)$$

Note that the expression for H_k is consistent with the particular cases considered above.

Now we consider the deformations of the equations for the fermionic one-forms:

$$\begin{aligned} 0 &= D\Phi^{\alpha(2k+1)} + d_k e_{\beta(2)}\Phi^{\alpha(2k+1)\beta(2)} + c_k e^\alpha{}_\beta \Phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)}e^{\alpha(2)}\Phi^{\alpha(2k-1)} \\ &\quad + \alpha_k \Omega^{\alpha(2k+1)\beta}\Psi_\beta + \beta_k \Omega^{\alpha(2k)}\Psi^\alpha \end{aligned} \quad (4.4)$$

$$0 = D\Phi^\alpha + d_0 e_{\beta(2)}\Phi^{\alpha\beta(2)} + c_0 e^\alpha{}_\beta \Phi^\beta + 4d_{(-1)}^2 E^\alpha{}_\beta \phi^\beta + \alpha_0 \Omega^{\alpha\beta}\Psi_\beta + \beta_0 e_{\beta(2)}B^{\beta(2)}\Psi^\alpha.$$

Their consistency requires (here $\hat{\alpha}^2 = 2\hat{\beta}^2 = (M_2 - s\lambda)G_s^2$):

$$\begin{aligned} \alpha_k^2 &= \frac{(s-k-1)}{(2k+3)}[M_2 + (k+1)\lambda]\hat{\alpha}^2 \\ \beta_k^2 &= \frac{(s+k+1)}{(k+1)(2k+1)^2}[M_2 - (k+1)\lambda]\hat{\beta}^2 \\ \beta_0 &= \frac{(s+1)}{2}[M_2 - \lambda]\alpha_0. \end{aligned} \quad (4.5)$$

At last, the deformations of the equations for the Stueckelberg fermionic zero-forms look like:

$$\begin{aligned}
0 &= D\phi^{\alpha(2k+1)} + \Phi^{\alpha(2k+1)} + c_k e^\alpha_\beta \phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)} e^{\alpha(2)} \phi^{\alpha(2k-1)} + d_k e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} \\
&\quad - G_k B^{\alpha(2k+1)\beta} \Psi_\beta - H_k B^{\alpha(2k)} \Psi^\alpha \tag{4.6} \\
0 &= D\phi^{\alpha(2s-1)} + \Phi^{\alpha(2s-1)} + c_{(s-1)} e^\alpha_\beta \phi^{\alpha(2s-2)\beta} + \frac{d_{(s-2)}}{(s-1)(2s-1)} e^{\alpha(2)} \phi^{\alpha(2s-3)} \\
&\quad - e_{\beta(2)} \phi^{\alpha(2s-1)\beta(2)} - G_{s-1} B^{\alpha(2s-1)\beta} \Psi_\beta - H_{s-1} B^{\alpha(2s-2)} \Psi^\alpha .
\end{aligned}$$

Their consistency give

$$\begin{aligned}
G_k &= \alpha_k, & H_k &= \beta_k, & k &\leq s-2 \tag{4.7} \\
H_{s-1} &= \beta_{s-1}, & G_s &= G_{s-1} = \frac{(2s-1)\beta_{s-1}}{M_2 - s\lambda} .
\end{aligned}$$

Let us now turn to the bosonic equations. Once again it is convenient to begin with the deformations of the equations for the gauge invariant zero-forms:

$$\begin{aligned}
0 &= DB^{\alpha(2k)} - e_{\beta(2)} B^{\alpha(2k)\beta(2)} - A_k e^\alpha_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)} B^{\alpha(2k-2)} \\
&\quad - E_k \phi^{\alpha(2k)\beta} \Psi_\beta - F_k \phi^{\alpha(2k-1)} \Psi^\alpha . \tag{4.8}
\end{aligned}$$

They give the same relation on masses and:

$$E_k = E_s, \quad F_k = -\frac{(s-k)}{2k(2k+1)} [M_2 + k\lambda] E_s, \quad k \geq s . \tag{4.9}$$

Note that the expression for F_k is also consistent with the particular cases considered above.

Deformations of the equations for the bosonic one-forms look like:

$$\begin{aligned}
0 &= D\Omega^{\alpha(2k)} + b_k e_{\beta(2)} \Omega^{\alpha(2k)\beta(2)} + a_k e^\alpha_\beta \Omega^{\alpha(2k-1)\beta} + \frac{b_{k-1}}{k(2k-1)} e^{\alpha(2)} \Omega^{\alpha(2k-2)} \\
&\quad + \gamma_k \Phi^{\alpha(2k)\beta} \Psi_\beta + \delta_k \Phi^{\alpha(2k-1)} \Psi^\alpha \\
0 &= D\Omega^{\alpha(2)} + b_1 e_{\beta(2)} \Omega^{\alpha(2)\beta(2)} + a_1 e^\alpha_\beta \Omega^{\alpha\beta} + 2b_0^2 E^\alpha_\beta B^{\alpha\beta} \tag{4.10} \\
&\quad + \gamma_1 \Phi^{\alpha(2)\beta} \Psi_\beta + \delta_1 \Phi^\alpha \Psi^\alpha + \gamma_0 e^{\alpha(2)} \phi^\beta \Psi_\beta .
\end{aligned}$$

They give (here $\hat{\gamma}^2 = 2\hat{\delta}^2 = E_s^2/2(M_2 - s\lambda)$):

$$\begin{aligned}
\gamma_k^2 &= \frac{(s+k+1)}{(k+1)} [M_2 - (k+1)\lambda] \hat{\gamma}^2 \\
\delta_k^2 &= \frac{(s-k)}{k^2(2k+1)} [M_2 + k\lambda] \hat{\delta}^2 \tag{4.11} \\
\gamma_0 &= -(s+1) [M_2 - \lambda] \delta_1 .
\end{aligned}$$

At last, the deformations of the equations for the Stueckelberg bosonic zero-forms:

$$\begin{aligned}
 0 &= DB^{\alpha(2k)} + \Omega^{\alpha(2k)} + a_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} + \frac{b_{k-1}}{k(2k-1)} e^{\alpha(2)} B^{\alpha(2k-2)} + b_k e_{\beta(2)} B^{\alpha(2k)\beta(2)} \\
 &\quad - E_k \phi^{\alpha(2k)\beta} \Psi_\beta - F_k \phi^{\alpha(2k-1)} \Psi^\alpha \\
 0 &= DB^{\alpha(2s-2)} + \Omega^{\alpha(2s-2)} + a_{s-1} e^\alpha{}_\beta B^{\alpha(2s-3)\beta} + \frac{b_{s-2}}{(s-1)(2s-3)} e^{\alpha(2)} B^{\alpha(2s-4)} \\
 &\quad - e_{\beta(2)} B^{\alpha(2s-2)\beta(2)} - E_{s-1} \phi^{\alpha(2s-2)\beta} \Psi_\beta - F_{s-1} \phi^{\alpha(2s-3)} \Psi^\alpha.
 \end{aligned} \tag{4.12}$$

Their consistency requires:

$$E_k = \gamma_k, \quad F_k = \delta_k, \quad k \leq s-1, \quad E_s = E_{s-1}. \tag{4.13}$$

Thus we have expressed all the coefficients in terms of just two arbitrary constants (one can choose (E_s, G_s) or $(\hat{\alpha}, \hat{\gamma})$). The complete set of supertransformations leaving all unfolded equations invariant has the form:

$$\begin{aligned}
 \delta\Omega^{\alpha(2k)} &= \gamma_k \Phi^{\alpha(2k)\beta} \zeta_\beta + \delta_k \Phi^{\alpha(2k)} \zeta^\alpha \\
 \delta\Omega^{\alpha(2)} &= \gamma_1 \Phi^{\alpha(2)\beta} \zeta_\beta + \delta_1 \Phi^\alpha \zeta^\alpha + \gamma_0 e^{\alpha(2)} \phi^\beta \zeta_\beta \\
 \delta B^{\alpha(2k)} &= E_k \phi^{\alpha(2k)\beta} \zeta_\beta + F_k \phi^{\alpha(2k-1)} \zeta^\alpha \\
 \delta\Phi^{\alpha(2k+1)} &= \alpha_k \Omega^{\alpha(2k+1)\beta} \zeta_\beta + \beta_k \Omega^{\alpha(2k)} \zeta^\alpha \\
 \delta\Phi^\alpha &= \alpha_0 \Omega^{\alpha\beta} \zeta_\beta + \beta_0 e_{\beta(2)} B^{\beta(2)} \zeta^\alpha \\
 \delta\phi^{\alpha(2k+1)} &= G_k B^{\alpha(2k+1)\beta} \zeta_\beta + H_k B^{\alpha(2k)} \zeta^\alpha.
 \end{aligned} \tag{4.14}$$

5 Supermultiplet $(s, s - 1/2)$

In this case the bosonic equations are the same as before while the unfolded formulation for the fermionic spin $s - 1/2$ field (appendix D) requires a set of one-forms $\Phi^{\alpha(2k+1)}$ and Stueckelberg zero-forms $\phi^{\alpha(2k+1)}$, $0 \leq k \leq s - 2$ as well as an infinite number of gauge invariant zero-forms $\phi^{\alpha(2k+1)}$, $k \geq s - 1$.

Deformations for gauge invariant fermionic zero-forms:

$$\begin{aligned}
 0 &= D\phi^{\alpha(2k+1)} - e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)} \phi^{\alpha(2k-1)} \\
 &\quad - G_k B^{\alpha(2k+1)\beta} \Psi_\beta - H_k B^{\alpha(2k)} \Psi^\alpha.
 \end{aligned} \tag{5.1}$$

Their consistency gives a mass relation:

$$M_1 = M_2 + \frac{\lambda}{2} \tag{5.2}$$

as well as fixes all coefficients in terms of G_{s-1} :

$$G_k = G_{s-1}, \quad H_k = -\frac{(s-k-1)}{2(k+1)(2k+1)} [M_2 + (k+1)\lambda] G_{s-1}. \tag{5.3}$$

Deformations for the fermionic one-forms look like:

$$\begin{aligned}
 0 &= D\Phi^{\alpha(2k+1)} + d_k e_{\beta(2)} \Phi^{\alpha(2k+1)\beta(2)} + c_k e^\alpha{}_\beta \Phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)} e^{\alpha(2)} \Phi^{\alpha(2k-1)} \\
 &\quad + \alpha_k \Omega^{\alpha(2k+1)\beta} \Psi_\beta + \beta_k \Omega^{\alpha(2k)} \Psi^\alpha \\
 0 &= D\Phi^\alpha + d_0 e_{\beta(2)} \Phi^{\alpha\beta(2)} + c_0 e^\alpha{}_\beta \Phi^\beta + 4d_{(s-1)}^2 E^\alpha{}_\beta \phi^\beta + \alpha_0 \Omega^{\alpha\beta} \Psi_\beta + \beta_0 e_{\beta(2)} B^{\beta(2)} \Psi^\alpha.
 \end{aligned} \tag{5.4}$$

Their consistency requires (here $\hat{\alpha}^2 = 2\hat{\beta}^2 = G_{s-1}^2/(M_2 - (s-1)\lambda)$):

$$\begin{aligned}
 \alpha_k^2 &= \frac{(s+k+1)}{(2k+3)} [M_2 - (k+1)\lambda] \hat{\alpha}^2 \\
 \beta_k^2 &= \frac{(s-k-1)}{(k+1)(2k+1)^2} [M_2 + (k+1)\lambda] \hat{\beta}^2 \\
 \beta_0 &= \frac{(s-1)}{2} [M_2 + \lambda] \alpha_0.
 \end{aligned} \tag{5.5}$$

At the same time deformations for the Stueckelberg fermionic zero-forms:

$$\begin{aligned}
 0 &= D\phi^{\alpha(2k+1)} + \Phi^{\alpha(2k+1)} + c_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)} e^{\alpha(2)} \phi^{\alpha(2k-1)} + d_k e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} \\
 &\quad - G_k B^{\alpha(2k+1)\beta} \Psi_\beta - H_k B^{\alpha(2k)} \Psi^\alpha \\
 0 &= D\phi^{\alpha(2s-3)} + \Phi^{\alpha(2s-3)} + c_{(s-2)} e^\alpha{}_\beta \phi^{\alpha(2s-4)\beta} + \frac{d_{(s-3)}}{(s-2)(2s-3)} e^{\alpha(2)} \phi^{\alpha(2s-5)} \\
 &\quad - e_{\beta(2)} \phi^{\alpha(2s-3)\beta(2)} - G_{s-2} B^{\alpha(2s-3)\beta} \Psi_\beta - H_{s-2} B^{\alpha(2s-4)} \Psi^\alpha.
 \end{aligned} \tag{5.6}$$

They give:

$$G_k = \alpha_k, \quad H_k = \beta_k, \quad k \leq s-2, \quad G_{s-1} = G_{s-2}. \tag{5.7}$$

Deformations for bosonic gauge invariant zero-forms:

$$\begin{aligned}
 0 &= DB^{\alpha(2k)} - e_{\beta(2)} B^{\alpha(2k)\beta(2)} - A_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)} B^{\alpha(2k-2)} \\
 &\quad - E_k \phi^{\alpha(2k)\beta} \Psi_\beta - F_k \phi^{\alpha(2k-1)} \Psi^\alpha.
 \end{aligned} \tag{5.8}$$

Their consistency gives the same mass relation as before and leads to the following expressions for all coefficients in terms of E_s :

$$E_k = E_s, \quad F_k = -\frac{(s+k)}{2k(2k+1)} [M_2 - k\lambda] E_s, \quad k \geq s. \tag{5.9}$$

Deformations for bosonic one-forms:

$$\begin{aligned}
 0 &= D\Omega^{\alpha(2k)} + b_k e_{\beta(2)} \Omega^{\alpha(2k)\beta(2)} + a_k e^\alpha{}_\beta \Omega^{\alpha(2k-1)\beta} + \frac{b_{k-1}}{k(2k-1)} e^{\alpha(2)} \Omega^{\alpha(2k-2)} \\
 &\quad + \gamma_k \Phi^{\alpha(2k)\beta} \Psi_\beta + \delta_k \Phi^{\alpha(2k-1)} \Psi^\alpha, \quad \gamma_{s-1} = 0 \\
 0 &= D\Omega^{\alpha(2)} + b_1 e_{\beta(2)} \Omega^{\alpha(2)\beta(2)} + a_1 e^\alpha{}_\beta \Omega^{\alpha\beta} + 2b_0^2 E^\alpha{}_\beta B^{\alpha\beta} \\
 &\quad + \gamma_1 \Phi^{\alpha(2)\beta} \Psi_\beta + \delta_1 \Phi^\alpha \Psi^\alpha + \gamma_0 e^{\alpha(2)} \phi^\beta \Psi_\beta.
 \end{aligned} \tag{5.10}$$

Their consistence requires (here $\hat{\gamma}^2 = 2\hat{\delta}^2 = \frac{1}{2}(M_2 - (s-1)\lambda)E_s^2$):

$$\begin{aligned}\gamma_k^2 &= \frac{(s-k-1)}{(k+1)}[M_2 + (k+1)\lambda]\hat{\gamma}^2 \\ \delta_k^2 &= \frac{(s+k)}{k^2(2k+1)}[M_2 - k\lambda]\hat{\delta}^2 \\ \gamma_0 &= -(s-1)[M_2 + \lambda]\delta_1.\end{aligned}\tag{5.11}$$

At last deformations for the Stueckelberg bosonic zero-forms:

$$\begin{aligned}0 &= DB^{\alpha(2k)} + \Omega^{\alpha(2k)} + a_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} + \frac{b_{k-1}}{k(2k-1)} e^{\alpha(2)} B^{\alpha(2k-2)} + b_k e_{\beta(2)} B^{\alpha(2k)\beta(2)} \\ &\quad - E_k \phi^{\alpha(2k)\beta} \Psi_\beta - F_k \phi^{\alpha(2k-1)} \Psi^\alpha \\ 0 &= DB^{\alpha(2s-2)} + \Omega^{\alpha(2s-2)} + a_{s-1} e^\alpha{}_\beta B^{\alpha(2s-3)\beta} + \frac{b_{s-2}}{(s-1)(2s-3)} e^{\alpha(2)} B^{\alpha(2s-4)} \\ &\quad - e_{\beta(2)} B^{\alpha(2s-2)\beta(2)} - E_{s-1} \phi^{\alpha(2s-2)\beta} \Psi_\beta - F_{s-1} \phi^{\alpha(2s-3)} \Psi^\alpha.\end{aligned}\tag{5.12}$$

We obtain:

$$\begin{aligned}E_k &= \gamma_k, & F_k &= \delta_k, & k &\leq s-2, & F_{s-1} &= \delta_{s-1} \\ E_s &= E_{s-1} = \frac{\gamma_{s-1}}{b_{s-2}} = \frac{2(s-1)\delta_{s-1}}{(M_2 - (s-1)\lambda)}.\end{aligned}\tag{5.13}$$

Supertransformations look the same as before (4.14) (taking into account that $\Phi^{\alpha(2s-1)}$ is absent now), but with the new expressions for all coefficients.

6 Summary and conclusion

In this paper we have presented the systematic derivation of the unfolded equations for three dimensional supersymmetric higher spin theory. Final results are formulated in sections 4 and 5. In particular, the equations (4.14) solves a problem of the supersymmetry transformations for the supermultiplets $(s, s + 1/2)$ and $(s, s - 1/2)$ respectively. Thus, we have the complete system of the unfolded equations and the corresponding supersymmetry transformations for the most general massive $(1, 0)$ supersymmetric field model in AdS_3 space.

As it is typical for the unfolded formalism, our results contain an infinite number of fields and so an infinite number of equations which are non-Lagrangian ones. As we pointed out in the Introduction, the same problem can be analysed on the base of approach developed in the papers [19–21]. Although this approach looks like more complicated, it operates with a finite number of fields and there are no reasons to expect that the final equations of motion should be non-Lagrangian. Therefore we suppose that the unfolded equations, obtained here, can be somehow reformulated, perhaps with eliminating some auxiliary fields, so that we will obtain the Lagrangian formulation. We guess that this aspect deserves a special consideration.

The equations and the supersymmetry transformations are obtained here in the component formulation. Both the equations and the supersymmetry transformations include all the fields which are necessary for consistent supersymmetric higher spin dynamics. However, they do not contain the auxiliary fields needed for off-shell supersymmetry. We suppose that a solution of off-shell supersymmetry problem can be realized on the base of appropriate superfield formalism. Developing such a formalism is one of the open problem in the three dimensional massive supersymmetric higher spin theory.

After the first version of this paper has been appeared in ArXiv, we were informed by S.M. Kuzenko that he and M. Tsulaia have constructed the massive higher spin off-shell $\mathcal{N} = 1$ supermultiplets on AdS_3 .

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A Unfolded equations for low spins

In this appendix we have collected all unfolded equations for the bosonic and fermionic fields with spin $0 \leq s \leq 2$.

A.1 Spin 0

Unfolded formulation for the spin-0 is very well known. It requires an infinite set of zero-forms $\pi^{\alpha(2k)}$, $k \geq 0$ satisfying the equations:

$$0 = D\pi^{\alpha(2k)} - e_{\beta(2)}\pi^{\alpha(2k)\beta(2)} - B_k e^{\alpha(2)}\pi^{\alpha(2k-2)}. \quad (\text{A.1})$$

Their consistency requires:

$$(2k + 3)B_{k+1} = (2k - 1)B_k + \frac{\lambda^2}{2} \quad (\text{A.2})$$

and leads to the solution:

$$B_k = -\frac{m_0^2 - (k^2 - 1)\lambda^2}{2(4k^2 - 1)}. \quad (\text{A.3})$$

A.2 Spin 1/2

Similarly, the unfolded formulation for the spin-1/2 requires an infinite number of fermionic zero-forms $\Phi^{\alpha(2k+1)}$, $k \geq 0$ satisfying the equations:

$$0 = D\Phi^{\alpha(2k+1)} - e_{\beta(2)}\Phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \Phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)}\Phi^{\alpha(2k-1)}. \quad (\text{A.4})$$

Their consistency requires:

$$(2k+5)C_{k+1} = (2k+1)C_k, \quad (k+2)D_{k+1} = kD_k - C_k^2 + \frac{\lambda^2}{4} \quad (\text{A.5})$$

and leads to the solution:

$$C_k = -\frac{m}{(2k+3)(2k+1)}, \quad D_k = -\frac{m^2}{2(2k+1)^2} + \frac{\lambda^2}{8}. \quad (\text{A.6})$$

A.3 Spin 1

In this case one also needs an infinite number of bosonic zero-forms $B^{\alpha(2k)}$ but now with $k \geq 1$. The unfolded equations have the form:

$$0 = DB^{\alpha(2k)} - e_{\beta(2)}B^{\alpha(2k)\beta(2)} - A_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)}B^{\alpha(2k-2)}. \quad (\text{A.7})$$

Their consistency implies the relations on the coefficients A , B :

$$(k+2)A_{k+1} = kA_k, \quad (2k+3)B_{k+1} = (2k-1)B_k - 2A_k^2 + \frac{\lambda^2}{2} \quad (\text{A.8})$$

which have the following solution:

$$A_k = -\frac{m_1}{2k(k+1)}, \quad B_k = -\frac{(k^2-1)}{2(4k^2-1)} \left[\frac{m_1^2}{k^2} - \lambda^2 \right]. \quad (\text{A.9})$$

A.4 Spin 3/2

The unfolded formulation for massive spin-3/2 already has the general pattern. Namely, it requires one-form Φ^α and Stueckelberg zero-form ϕ^α as well as an infinite number of gauge invariant zero-forms $\phi^{\alpha(2k+1)}$, $k \geq 1$. The unfolded equations for the first two fields look like:

$$\begin{aligned} 0 &= D\Phi^\alpha + M e^\alpha{}_\beta \Psi^\beta + 2m E^\alpha{}_\beta \phi^\beta \\ 0 &= D\phi^\alpha + 2m\Phi^\alpha + M e^\alpha{}_\beta \phi^\beta - e_{\beta(2)}\phi^{\alpha\beta(2)} \end{aligned} \quad (\text{A.10})$$

where:

$$M^2 = m^2 + \frac{\lambda^2}{4}.$$

Equations for the remaining fields have the same form as in the spin-1/2 case:

$$0 = D\phi^{\alpha(2k+1)} - e_{\beta(2)}\phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)}\phi^{\alpha(2k-1)}. \quad (\text{A.11})$$

Consistency conditions are also the same as for the spin-1/2 case but their solution now:

$$C_k = -\frac{3M}{(2k+3)(2k+1)}, \quad D_k = -\frac{(k+2)(k-1)}{8k(k+1)} \left[\frac{4M^2}{(2k+1)^2} - \lambda^2 \right]. \quad (\text{A.12})$$

A.5 Spin 2

In this case we also need the one-form $\Omega^{\alpha(2)}$ and Stueckelberg zero-form $B^{\alpha(2)}$ as well as an infinite number of gauge invariant zero-forms $B^{\alpha(2k)}$, $k \geq 2$. The equations for the first two fields:

$$\begin{aligned} 0 &= D\Omega^{\alpha(2)} + mE^\alpha{}_\beta B^{\alpha\beta} + \frac{M}{2}e^\alpha{}_\beta \Omega^{\alpha\beta} \\ 0 &= DB^{\alpha(2)} + m\Omega^{\alpha(2)} + \frac{M}{2}e^\alpha{}_\beta B^{\alpha\beta} - e_{\beta(2)}B^{\alpha(2)\beta(2)} \end{aligned} \quad (\text{A.13})$$

where:

$$M^2 = m^2 + \lambda^2.$$

The equations for the remaining fields have the same form as in the spin-1 case:

$$0 = DB^{\alpha(2k)} - e_{\beta(2)}B^{\alpha(2k)\beta(2)} - A_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)} B^{\alpha(2k-2)}. \quad (\text{A.14})$$

Consistency conditions are also the same as for spin-1, but their solution now:

$$A_k = -\frac{M}{k(k+1)}, \quad B_k = -\frac{(k^2-4)}{2(4k^2-1)} \left[\frac{M^2}{k^2} - \lambda^2 \right]. \quad (\text{A.15})$$

B Unfolded equations for the bosonic spin- s field

The unfolded formulation for the bosonic spin- s field requires a set of one-forms $\Omega^{\alpha(2k)}$ and Stueckelberg zero-forms $B^{\alpha(2k)}$, $1 \leq k \leq s-1$ as well as an infinite number of gauge invariant zero-forms $B^{\alpha(2k)}$, $k \geq s$.

Equations for one-forms ($2 \leq k \leq s-1$, $b_{s-1} = 0$):

$$\begin{aligned} 0 &= D\Omega^{\alpha(2k)} + b_k e_{\beta(2)} \Omega^{\alpha(2k)\beta(2)} + a_k e^\alpha{}_\beta \Omega^{\alpha(2k-1)\beta} + \frac{b_{k-1}}{k(2k-1)} e^{\alpha(2)} \Omega^{\alpha(2k-2)} \\ 0 &= D\Omega^{\alpha(2)} + b_1 e_{\beta(2)} \Omega^{\alpha(2)\beta(2)} + a_1 e^\alpha{}_\beta \Omega^{\alpha\beta} + 2b_0^2 E^\alpha{}_\beta B^{\alpha\beta}. \end{aligned} \quad (\text{B.1})$$

Here

$$\begin{aligned} M_2^2 &= m_2^2 + (s-1)^2 \lambda^2 \\ a_k &= \frac{M_2 s}{2k(k+1)}, \quad b_k^2 = \frac{(s-k-1)(s+k+1)}{2(k+1)(2k+3)} [M_2^2 - (k+1)^2 \lambda^2]. \end{aligned} \quad (\text{B.2})$$

Equations for the Stueckelberg zero-forms ($1 \leq k \leq s-2$):

$$\begin{aligned} 0 &= DB^{\alpha(2k)} + \Omega^{\alpha(2k)} + a_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} + \frac{b_{k-1}}{k(2k-1)} e^{\alpha(2)} B^{\alpha(2k-2)} + b_k e_{\beta(2)} B^{\alpha(2k)\beta(2)} \\ 0 &= DB^{\alpha(2s-2)} + \Omega^{\alpha(2s-2)} + a_{s-1} e^\alpha{}_\beta B^{\alpha(2s-3)\beta} + \frac{b_{s-2}}{(s-1)(2s-3)} e^{\alpha(2)} B^{\alpha(2s-4)} \\ &\quad - e_{\beta(2)} B^{\alpha(2s-2)\beta(2)}. \end{aligned} \quad (\text{B.3})$$

Equations for gauge invariant zero-forms:

$$0 = DB^{\alpha(2k)} - e_{\beta(2)}B^{\alpha(2k)\beta(2)} - A_k e^\alpha{}_\beta B^{\alpha(2k-1)\beta} - B_k e^{\alpha(2)} B^{\alpha(2k-2)}. \quad (\text{B.4})$$

Here

$$A_k = -\frac{Ms}{2k(k+1)}, \quad B_k = -\frac{(k^2 - s^2)}{2(4k^2 - 1)} \left[\frac{M^2}{k^2} - \lambda^2 \right]. \quad (\text{B.5})$$

C Unfolded equations for the fermionic spin $(s + 1/2)$ field

The unfolded formulation for the fermionic spin $s + 1/2$ field requires a set of one-forms $\Phi^{\alpha(2k+1)}$ and Stueckelberg zero-forms $\phi^{\alpha(2k+1)}$, $0 \leq k \leq s - 1$ as well as an infinite number of gauge invariant zero-forms $\phi^{\alpha(2k+1)}$, $k \geq s$.

Equations for one-forms ($1 \leq k \leq s - 1$, $d_{s-1} = 0$):

$$\begin{aligned} 0 &= D\Phi^{\alpha(2k+1)} + d_k e_{\beta(2)} \Phi^{\alpha(2k+1)\beta(2)} + c_k e^\alpha{}_\beta \Phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)} e^{\alpha(2)} \Phi^{\alpha(2k-1)} \\ 0 &= D\Phi^\alpha + d_0 e_{\beta(2)} \Phi^{\alpha\beta(2)} + c_0 e^\alpha{}_\beta \Phi^\beta + 4d_{(-1)}^2 E^\alpha{}_\beta \phi^\beta. \end{aligned} \quad (\text{C.1})$$

Here

$$\begin{aligned} M_1^2 &= m_1^2 + \left(s - \frac{1}{2} \right)^2 \lambda^2 \\ c_k &= \frac{(2s+1)M_1}{(2k+1)(2k+3)}, \quad d_k^2 = \frac{(s-k-1)(s+k+2)}{2(2k+3)(k+2)} \left[M_1^2 - \frac{(2k+3)^2}{4} \lambda^2 \right]. \end{aligned} \quad (\text{C.2})$$

Equations for the Stueckelberg zero-forms ($0 \leq k \leq s - 2$):

$$\begin{aligned} 0 &= D\phi^{\alpha(2k+1)} + \Phi^{\alpha(2k+1)} + c_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)} e^{\alpha(2)} \phi^{\alpha(2k-1)} + d_k e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} \\ 0 &= D\phi^{\alpha(2s-1)} + \Phi^{\alpha(2s-1)} + c_{(s-1)} e^\alpha{}_\beta \phi^{\alpha(2s-2)\beta} + \frac{d_{(s-2)}}{(s-1)(2s-1)} e^{\alpha(2)} \phi^{\alpha(2s-3)} \\ &\quad - e_{\beta(2)} \phi^{\alpha(2s-1)\beta(2)}. \end{aligned} \quad (\text{C.3})$$

Equations for gauge invariant zero-forms:

$$0 = D\phi^{\alpha(2k+1)} - e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)} \phi^{\alpha(2k-1)}. \quad (\text{C.4})$$

Here

$$C_k = -\frac{(2s+1)M_1}{(2k+1)(2k+3)}, \quad D_k = -\frac{(k-s)(k+s+1)}{2k(k+1)} \left[\frac{M_1^2}{(2k+1)^2} - \frac{\lambda^2}{4} \right]. \quad (\text{C.5})$$

D Unfolded equations for the fermionic spin ($s - 1/2$) field

The unfolded formulation for the fermionic spin $s - 1/2$ field requires a set of one-forms $\Phi^{\alpha(2k+1)}$ and Stueckelberg zero-forms $\phi^{\alpha(2k+1)}$, $0 \leq k \leq s - 2$ as well as an infinite number of gauge invariant zero-forms $\phi^{\alpha(2k+1)}$, $k \geq s - 1$.

Equations for the one-forms ($1 \leq k \leq s - 2$, $d_{s-2} = 0$):

$$\begin{aligned} 0 &= D\Psi^{\alpha(2k+1)} + d_k e_{\beta(2)} \Phi^{\alpha(2k+1)\beta(2)} + c_k e^\alpha{}_\beta \Phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)} e^{\alpha(2)} \Phi^{\alpha(2k-1)} \\ 0 &= D\Phi^\alpha + d_0 e_{\beta(2)} \Phi^{\alpha\beta(2)} + c_0 e^\alpha{}_\beta \Phi^\beta + 4d_{(-1)}{}^2 E^\alpha{}_\beta \phi^\beta. \end{aligned} \quad (\text{D.1})$$

Here

$$\begin{aligned} M_1^2 &= m_1^2 + \left(s - \frac{3}{2}\right)^2 \lambda^2 \\ c_k &= \frac{(2s-1)M_1}{(2k+1)(2k+3)}, \quad d_k^2 = \frac{(s-k-2)(s+k+1)}{2(2k+3)(k+2)} \left[M_1^2 - \frac{(2k+3)^2}{4} \lambda^2 \right]. \end{aligned} \quad (\text{D.2})$$

Equations for the Stueckelberg zero-forms ($0 \leq k \leq s - 3$):

$$\begin{aligned} 0 &= D\phi^{\alpha(2k+1)} + \Phi^{\alpha(2k+1)} + c_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} + \frac{d_{k-1}}{k(2k+1)} e^{\alpha(2)} \phi^{\alpha(2k-1)} + d_k e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} \\ 0 &= D\phi^{\alpha(2s-3)} + \Phi^{\alpha(2s-3)} + c_{(s-2)} e^\alpha{}_\beta \phi^{\alpha(2s-4)\beta} + \frac{d_{(s-3)}}{(s-2)(2s-3)} e^{\alpha(2)} \phi^{\alpha(2s-5)} \\ &\quad - e_{\beta(2)} \phi^{\alpha(2s-3)\beta(2)}. \end{aligned} \quad (\text{D.3})$$

Equations for gauge invariant zero-forms:

$$0 = D\phi^{\alpha(2k+1)} - e_{\beta(2)} \phi^{\alpha(2k+1)\beta(2)} - C_k e^\alpha{}_\beta \phi^{\alpha(2k)\beta} - D_k e^{\alpha(2)} \phi^{\alpha(2k-1)}. \quad (\text{D.4})$$

Here

$$C_k = -\frac{(2s-1)M_1}{(2k+1)(2k+3)}, \quad D_k = -\frac{(k-s+1)(k+s)}{2k(k+1)} \left[\frac{M_1^2}{(2k+1)^2} - \frac{\lambda^2}{4} \right]. \quad (\text{D.5})$$

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