

**SU(2) reduction in  $\mathcal{N} = 4$  supersymmetric mechanics**Sergey Krivonos<sup>1,\*</sup> and Olaf Lechtenfeld<sup>2,†</sup><sup>1</sup>*Bogoliubov Laboratory of Theoretical Physics, JINR, 141980 Dubna, Russia*<sup>2</sup>*Institut für Theoretische Physik, Leibniz Universität Hannover, 30167 Hannover, Germany*

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We perform an  $su(2)$  Hamiltonian reduction of the general  $su(2)$ -invariant action for a self-coupled  $(4, 4, 0)$  supermultiplet. As a result, we elegantly recover the  $\mathcal{N} = 4$  supersymmetric mechanics with spin degrees of freedom which was recently constructed in [S. Fedoruk, E. Ivanov, and O. Lechtenfeld, Phys. Rev. D **79**, 105015 (2009)]. This observation underscores the exceptional role played by the root supermultiplet in  $\mathcal{N} = 4$  supersymmetric mechanics.

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**I. INTRODUCTION**

In a recent paper [1],  $\mathcal{N} = 4$  superconformal mechanics with  $n$  bosonic and  $4n$  fermionic degrees of freedom has been endowed with a potential term through a coupling to auxiliary supermultiplets with  $4n$  bosonic and  $4n$  fermionic components [2]. This combination gave rise to an  $O\text{Sp}(4|2)$  supersymmetric  $n$ -particle Calogero model. Subsequently, the one-particle case, i.e.  $O\text{Sp}(4|2)$  superconformal mechanics, was analyzed on the classical and quantum level [3]. Simultaneously, it was demonstrated that the potential-generating strategy works perfectly for the most general  $D(2, 1; \alpha)$  superconformal one-particle mechanics [4]. It is quite satisfying how the spin degrees of freedom appear in the bosonic sector, with only first time derivatives in the action. Thus, the proposed coupling of two different  $\mathcal{N} = 4$  supermultiplets provides a simple and elegant way to incorporate spin degrees of freedom in supersymmetric mechanics.

In both previous treatments [3,4], on mass shell all components of the basic  $(1, 4, 3)$  supermultiplet are expressed through those of the “auxiliary”  $(4, 4, 0)$  one. It seems that just this auxiliary supermultiplet plays a fundamental role in the construction. It is therefore natural to inquire whether these models can be reformulated *purely* in terms of  $(4, 4, 0)$  supermultiplets. Of course, such a reformulation has to be supplied with a Hamiltonian reduction, which would reduce the four physical bosons to one boson plus spin variables. Alternatively, the passage from  $SU(2)$ -symmetric  $(4, 4, 0)$  models to general  $(1, 4, 3)$  models via gauging was described in [2] using harmonic superspace.

Incidentally, spin degrees of freedom have appeared in a bosonic system after Hamiltonian reduction (on the Lagrangian level) via the second Hopf map  $S^7/S^3 \simeq S^4$  [5]. In the bosonic sector this reduced system resembles those in [3,4], besides the presence of four additional bosonic variables.

In the present paper we realize the above ideas and rederive the  $\mathcal{N} = 4$  supersymmetric “spin mechanics” of [3,4] by an  $su(2)$  Hamiltonian reduction applied to the general  $su(2)$  invariant action for a self-coupled  $(4, 4, 0)$  supermultiplet. It is a further manifestation of the fundamental importance of the root supermultiplet [6] in  $\mathcal{N} = 4$  supersymmetric mechanics [2,7,8].

**II. SU(2) REDUCTION**

Our point of departure is a quartet of real  $\mathcal{N} = 4$  superfields  $Q^{ia}$  with  $i, a = 1, 2$  defined in the  $\mathcal{N} = 4$  superspace  $\mathbb{R}^{(1|4)} = (t, \theta_i, \bar{\theta}^i)$  and subject to the constraints

$$D^{(i} Q^{j)a} = 0, \quad \bar{D}^{(i} Q^{j)a} = 0 \quad \text{and} \quad (Q^{ia})^\dagger = Q_{ia}, \quad (2.1)$$

where the corresponding covariant derivatives have the form

$$D^i = \frac{\partial}{\partial \theta_i} + i\bar{\theta}^i \partial_t, \quad \bar{D}_i = \frac{\partial}{\partial \bar{\theta}^i} + i\theta_i \partial_t, \quad (2.2)$$

so that  $\{D^i, \bar{D}_j\} = 2i\delta_j^i \partial_t$ .

This  $\mathcal{N} = 4$  supermultiplet describes four bosonic and four fermionic but zero auxiliary variables off shell [9,10]. Let us now introduce the composite  $\mathcal{N} = 4$  superfield<sup>1</sup>

$$X = 2(Q^{ia} Q_{ia})^{-1} \quad (2.3)$$

which, in virtue of (2.1), obeys the constraints [10]

$$D^i D_i X = \bar{D}_i \bar{D}^i X = [D^i, \bar{D}_i] X = 0. \quad (2.4)$$

The most general action for  $Q^{ia}$  is constructed by integrating an arbitrary superfunction  $\tilde{\mathcal{F}}(Q^{ia})$  over the whole  $\mathcal{N} = 4$  superspace. Here, we restrict ourselves to prepotentials of the form

<sup>1</sup>We stress that  $Q^{ia} Q_{ia} \sim e^{-U}$  in the standard parametrization [10], where  $U$  is the superdilatation. Therefore, the new superfield  $X$  is well defined.

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$$\tilde{\mathcal{F}}(Q^{ia}) = \mathcal{F}(X(Q^{ia})) \rightarrow S = -\frac{1}{8} \int dt d^4\theta \mathcal{F}(X). \quad (2.5)$$

The rationale for this selection is its manifest invariance under  $su(2)$  transformations acting on the “ $a$ ” index of  $Q^{ia}$ . This is the symmetry over which we are going to perform the Hamiltonian reduction.

In terms of components the action (2.5) reads

$$S = \int dt \left\{ G \left( \dot{x}^2 + i(\dot{\eta}^i \bar{\eta}_i - \eta^i \dot{\bar{\eta}}_i) + \frac{1}{2} x^2 \omega^{ij} \omega_{ij} \right) - i(2G + xG') \omega^{ij} \eta_i \bar{\eta}_j - \frac{1}{4} \left( G'' + 6 \frac{G'}{x} + 6 \frac{G}{x^2} \right) \eta^i \eta_i \bar{\eta}_j \bar{\eta}^j \right\}, \quad (2.6)$$

where

$$x = X|, \quad \eta^i = -iD^i X|, \quad \bar{\eta}_i = -i\bar{D}_i X|, \quad (2.7)$$

$$q^{ia} = \sqrt{X} Q^{ia}|, \quad G = \mathcal{F}''(X)|$$

and

$$\omega_{ij} = \dot{q}_i^a q_{ja} + \dot{q}_j^a q_{ia}. \quad (2.8)$$

Here, as usual,  $(\dots)|$  denotes the  $\theta_i = \bar{\theta}_i = 0$  limit.

To proceed we introduce the following substitution for the bosonic variables  $q^{ia}$  subject to  $q^{ia} q_{ia} = 2$ :

$$q^{11} = \frac{e^{-(i/2)\phi}}{\sqrt{1 + \Lambda \bar{\Lambda}}}, \quad q^{21} = -\frac{e^{-(i/2)\phi}}{\sqrt{1 + \Lambda \bar{\Lambda}}}, \quad (2.9)$$

$$q^{22} = (q^{11})^\dagger, \quad q^{12} = -(q^{12})^\dagger.$$

In terms of the new variables  $(\phi, \Lambda, \bar{\Lambda})$ , the  $su(2)$  rotations  $\delta q^{ia} = \gamma^{(ab)} q_b^i$  read [10]

$$\delta \Lambda = \gamma^{11} e^{i\phi} (1 + \Lambda \bar{\Lambda}), \quad \delta \bar{\Lambda} = \gamma^{22} e^{-i\phi} (1 + \Lambda \bar{\Lambda}),$$

$$\delta \phi = -2i\gamma^{12} + i\gamma^{22} e^{-i\phi} \Lambda - i\gamma^{11} e^{i\phi} \bar{\Lambda}. \quad (2.10)$$

It is easy to check that

$$\omega^{11} = 2 \frac{\dot{\Lambda} - i\Lambda \dot{\phi}}{1 + \Lambda \bar{\Lambda}}, \quad \omega^{22} = (\omega^{11})^\dagger \quad (2.11)$$

$$\text{and } \omega^{12} = i \frac{1 - \Lambda \bar{\Lambda}}{1 + \Lambda \bar{\Lambda}} \dot{\phi} + \frac{\dot{\Lambda} \bar{\Lambda} - \Lambda \dot{\bar{\Lambda}}}{1 + \Lambda \bar{\Lambda}}$$

are indeed invariant under (2.10), as is the whole action (2.6).

Next, we introduce the standard Poisson brackets

$$\{\pi, \Lambda\} = 1, \quad \{\bar{\pi}, \bar{\Lambda}\} = 1, \quad \{p_\phi, \phi\} = 1, \quad (2.12)$$

so that the generators of the transformations (2.10),

$$I_\phi = p_\phi, \quad I = e^{i\phi} [(1 + \Lambda \bar{\Lambda})\pi - i\bar{\Lambda} p_\phi], \quad (2.13)$$

$$\bar{I} = e^{-i\phi} [(1 + \Lambda \bar{\Lambda})\bar{\pi} + i\Lambda p_\phi],$$

will be the Noether constants of motion for the action (2.6). To perform the reduction over this  $SU(2)$  group we fix the

Noether constants as (c.f. [5])

$$I_\phi = m \quad \text{and} \quad I = \bar{I} = 0, \quad (2.14)$$

which yields

$$p_\phi = m \quad \text{and} \quad \pi = \frac{im\bar{\Lambda}}{1 + \Lambda \bar{\Lambda}}, \quad \bar{\pi} = -\frac{im\Lambda}{1 + \Lambda \bar{\Lambda}}. \quad (2.15)$$

Conducting a Routh transformation over the variables  $(\Lambda, \bar{\Lambda}, \phi)$ , we reduce the action (2.6) to

$$\tilde{S} = S - \int dt \{ \pi \dot{\Lambda} + \bar{\pi} \dot{\bar{\Lambda}} + p_\phi \dot{\phi} \} \quad (2.16)$$

and substitute the expressions (2.15) into  $\tilde{S}$ . A slightly lengthy but straightforward calculation gives

$$\tilde{S}_{\text{red}} = \int dt \left\{ G(x^2 + i(\dot{\eta}^i \bar{\eta}_i - \eta^i \dot{\bar{\eta}}_i)) - \frac{1}{4} \left( G'' - \frac{3}{2} \frac{(G')^2}{G} \right) \eta^2 \bar{\eta}^2 - \frac{m^2}{4x^2 G} - \frac{m(2G + xG')}{2x^2 G(1 + \Lambda \bar{\Lambda})} (2\Lambda \eta_1 \bar{\eta}_1 - 2\bar{\Lambda} \eta_2 \bar{\eta}_2 - (1 - \Lambda \bar{\Lambda})(\eta_1 \bar{\eta}_2 + \eta_2 \bar{\eta}_1)) \right\}. \quad (2.17)$$

To ensure that the reduction constraints (2.15) are satisfied we add Lagrange multiplier terms,

$$S_{\text{red}} = \tilde{S}_{\text{red}} + \int dt \left\{ m \dot{\phi} + \frac{im(\dot{\Lambda} \bar{\Lambda} - \Lambda \dot{\bar{\Lambda}})}{1 + \Lambda \bar{\Lambda}} \right\}. \quad (2.18)$$

Finally, by employing new variables  $v^i = q^{i1}$  and  $\bar{v}_i = (v^i)^\dagger$  we rewrite this action in the symmetric form

$$S_{\text{red}} = \int dt \left\{ G(x^2 + i(\dot{\eta}^i \bar{\eta}_i - \eta^i \dot{\bar{\eta}}_i)) - \frac{1}{4} \left( G'' - \frac{3}{2} \frac{(G')^2}{G} \right) \eta^2 \bar{\eta}^2 - \frac{m^2}{4x^2 G} + im(\dot{v}^i \bar{v}_i - v^i \dot{\bar{v}}_i) - \frac{m(2G + xG')}{2x^2 G} v^i \bar{v}^j (\eta_i \bar{\eta}_j + \eta_j \bar{\eta}_i) \right\}$$

with  $v^i \bar{v}_i = 1$ . (2.19)

Amazingly, this final action coincides with the one presented in [4] and specializes to the one derived in [3] for the choice of  $G = 1$ , which corresponds to  $OSp(4|2)$  symmetry.

We stress that the  $su(2)$  reduction algebra, realized in (2.10), commutes with all (super)symmetries of the action (2.5). Therefore, all symmetry properties of the theory [including the  $D(2, 1; \alpha)$  invariance for a properly chosen prepotential  $\mathcal{F}$ ] are preserved in our reduction.

### III. CONCLUSION

We have demonstrated that the novel  $\mathcal{N} = 4$  supersymmetric “spin mechanics” of [1,3,4] is nicely interpreted as an  $su(2)$  reduction of a self-interacting root supermultiplet with (4, 4, 0) component content. This procedure is remarkably simple and automatically successful.

An almost straightforward application of this insight is a similar  $su(2)$  reduction applied to the  $\mathcal{N} = 4$  “nonlinear” supermultiplet [10]. The resulting system will contain only spinor variables accompanied by four fermions. In this regard, one could also investigate the nonlinear root supermultiplet and its action [11].

Finally, we mention that our reduction will almost never work for the  $\mathcal{N} = 8$  supersymmetric mechanics in the literature. The reason is simple: these systems do not

possess any internal symmetry which commutes with all eight supersymmetries. This is also the situation discussed in [5]. The one positive exception is the “real”  $\mathcal{N} = 8$ ,  $d = 1$  hypermultiplet, which is obtained by dimensional reduction from  $\mathcal{N} = 2$ ,  $d = 4$  and requires  $\mathcal{N} = 8$ ,  $d = 1$  harmonic superspace [12,13]. We expect the corresponding  $su(2)$  reduction to produce some spin extension of the recently constructed  $\mathcal{N} = 8$  superconformal mechanics [14]. We intend to turn to this issue soon.

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