Effective gravity and OSp(N,4)-invariant matter

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We reexamine the OSp(N,4)-invariant interacting model of massless chiral and gauge superfields, whose superconformal invariance was instrumental both in proving the all-order no-renormalization of the mass and chiral self-interaction Lagrangians and in determining the linear superfield renormalization needed. We show that the renormalization of the gravitational action modifies only the cosmological term, without affecting higher-order tensors. This could explain why the effect of the cosmological constant is shadowed by the effects of Newtonian gravity. [S0556-2821(98)04202-7]

I. INTRODUCTION

The introduction of locally supersymmetric theories [1] was motivated by the wish for a unified description encompassing the theories of elementary particles and gravity. The resulting supergravity is not renormalizable, but its large symmetry provides powerful constraints, such as the vanishing of the lepton anomalous magnetic moment, as required by the supersymmetry of the theory [2]. The finiteness of the fermionic (and therefore also the bosonic) contribution to it, which can be traced back to an effective chiral symmetry in the gravitino sector [3], has been checked in [4,3] and, more recently in [5], making use of supersymmetry preserving regularization schemes.

An introduction to the effective action in quantum gravity can be found in [6]. The quantum corrections to the low-energy limit of a theory coupling gravity to scalar fields have been computed [7]. The equivalence principle has been invoked, in order to reduce the terms with an arbitrary number of derivatives in the effective theory [8]. This principle also constrains the spin-1 and spin-0 partners of the graviton in the $N=2,8$ supergravity multiplets [9,10].

OSp(N,4)-invariant models in the fixed four-dimensional background of anti–de Sitter (AdS$_4$) space occur as both ground state solutions of gauged extended supergravity theories (see [11] and references therein) and vacuum configurations for superstrings [12]. The OSp(N,4)-invariant generalization of the Wess-Zumino model [13] is the simplest one. We recall that the Wess-Zumino model with softly broken supersymmetry in de Sitter space plays a role in the Affleck-Dine mechanism [14] for baryogenesis, in contrast to the maximal symmetry of AdS$_4$ space which grants the existence of global supersymmetry. This mechanism is effective for supersymmetric grand unified theories, as the quantum corrections do not affect the flat directions in the superpotential, owing to the no-renormalization theorem [15].

Very recently, the one-loop effective potential along a flat direction in this model has been calculated [16]. In [17] the one-loop effective action was computed for scalar QED, taking into account the large-scale configurations that change the topology. Also, the question of the dynamics of a super-

string propagating in AdS$_4$ space with the OSp(1,4) supersymmetry group, deserves further study, within a geometrical framework, especially in connection with the underlying algebraic structure of the W-algebra extension of two-dimensional conformal symmetry (see, e.g., [18–21]).

The renormalization procedure for OSp(N,4)-invariant theories breaks the naturality implied by the no-renormalization theorem and so allows *all* classically invariant counterterms to appear in the divergent structure of the quantum effective action [22–34]. This breakdown of the no-renormalization theorem in AdS$_4$ space forces us to introduce a linear superfield in the effective action, for the purpose of renormalization. However, the corresponding modification of the classical potential does not induce the breaking of supersymmetry invariance [25,11]. The superconformal invariance of the model with interacting chiral and real gauge superfields in AdS$_4$ space, following the line suggested in Ref. [35], allowed us to prove to all orders in the perturbative series the nonrenormalization of the mass or the cubic interaction action [34].

We organize the present work as follows. We begin in Sec. II with recalling the superfield formulation of the AdS$_4$ interacting model of chiral and gauge superfields. The superconformal invariance of the massless model allows us to implement an expansion in the curvature effects, in terms of the interaction vertices of the quantum model. In Sec. III we present the renormalization of the gravitational action, based on the use of OSp(N,4) superfield techniques, and propose some interpretation for the shadowing of the effect of the cosmological constant by the effects of Newtonian gravity. We draw our conclusions in Sec. IV.

II. INTERACTING CHIRAL AND REAL GAUGE SUPERFIELDS

For the purpose of fixing notation, we briefly recall in this section a superspace approach [35,33,34] to the OSp(N,4)-invariant theory of a supergravity multiplet coupled to interacting chiral and real gauge supermultiplets.

In order to choose AdS$_4$ space as a background space for the matter and gauge model, we set the supergravity prepotential H to zero and introduce the background only through the compensator ϕ. Its equation of motion

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$$\bar{D}^2\bar{\phi}=\alpha\phi^2 \qquad (2.1)$$

can be obtained from the action for supergravity with the cosmological term

$$S=-\frac{3}{\kappa^2}\int d^4x d^4\theta E^{-1}+\left(\alpha\frac{1}{\kappa^2}\int d^4x d^2\theta\phi^3+\text{H.c.}\right), \qquad (2.2)$$

yielding the solution with a regular behavior at infinity

$$\phi=\frac{1}{1-\alpha\bar{\alpha}x^2/4}-\frac{\bar{\alpha}\theta^2}{(1-\alpha\bar{\alpha}x^2/4)^2}, \qquad (2.3)$$

with the inverse determinant $E^{-1}=\bar{\phi}\phi$. Then, by applying this solution to the construction of invariant actions of the general and chiral type, we can formulate different supersymmetric matter models in the given background. We recall the expression of the covariant derivatives in terms of the ϕ field

$$\bar{\nabla}_{\dot{\alpha}}=\phi^{-1}\bar{\phi}^{1/2}\bar{D}_{\dot{\alpha}}, \quad \nabla_\alpha=\bar{\phi}^{-1}\phi^{1/2}D_\alpha. \qquad (2.4)$$

The theory of interacting chiral (η) and real gauge superfields is described (in the gauge-chiral representation) by the action [35]

$$S(\eta,\bar{\eta},V)=\int d^4x d^4\theta E^{-1}\bar{\eta}_j[\exp(V)]^j{}_i\eta^i$$

$$+\left[\int d^4x d^2\theta\phi^3 W^2+\int d^4x d^2\theta\phi^3\right.$$

$$\left.\times\left(m\frac{1}{2}\eta^2+\lambda\frac{1}{6}\eta^3\right)+\text{H.c.}\right], \qquad (2.5)$$

with $V^i_j=V^A(T_A)^i{}_j$, and where $(T_A)^i{}_j$ is a matrix representation of the generators of the gauge group that leaves this action invariant. This model possesses a partial superconformal invariance, which has been exploited [33,34], in order to treat perturbatively the effects of the background curvature, when carrying out the renormalization procedure that yields its quantum effective action. All ϕ dependence in the free-field functional integral can be removed by carrying out a superconformal transformation, in accordance with the canonical weights of the matter and gauge fields and their sources:

$$\hat{\eta}=\phi\eta, \quad \hat{W}_\alpha=\phi^{3/2}W_\alpha, \quad \hat{J}=\phi^2 J, \quad \hat{J}_V=\phi\bar{\phi}J_V. \qquad (2.6)$$

Hence the quantum model can be described in the most natural form in terms of the transformed fields, i.e., defining the effective action in terms of the fields with carets.

The definition of $\hat{W}$ in terms of the familiar derivatives in flat background reads

$$\hat{W}_\alpha=i\bar{D}^2 D_\alpha V. \qquad (2.7)$$

The covariantization of this expression in the Yang-Mills chiral representation

$$W_\alpha=i(\bar{\nabla}^2+\alpha)\exp(-V)\nabla_\alpha\exp(V) \qquad (2.8)$$

gives W_α in terms of the background covariant derivatives (2.4). After the superconformal transformation, the gauge-fixing procedure can be carried out, along the line of the flat background theory. The resulting gauge propagator reads (in the Fermi-Feynman gauge, with $\xi=1$)

$$\langle VV\rangle_0=-\frac{1}{p^2}\delta^4(\theta-\theta'). \qquad (2.9)$$

It is worthwhile to notice that, as a consequence of superconformal invariance, the ghost propagators and vertices of the flat space-time theory, along with the usual flat space-time D algebra, remain intact in AdS_4 space.

Enforcing the boundary conditions needed, in order to preserve supersymmetry for the scalar and spinor propagators in AdS_4 [32], and evaluating the free-field functional integral, yields the vacuum expectation values [33]

$$\langle T\hat{\eta}(x')\hat{\bar{\eta}}(0)\rangle=\frac{1}{4\pi^2}\hat{\bar{D}}^2 D^2\delta^4(\theta-\theta')\frac{1}{(x')^2},$$

$$\langle T\hat{\eta}(x')\hat{\eta}(0)\rangle=\frac{1}{16\pi^2}(|\alpha|^2+2|\alpha|^3\theta'\theta). \qquad (2.10)$$

From the generating functional one can read the vertex contributions involving the field $\hat{\eta}$. It turns out [33,34] that there are no ϕ's at every such vertex (in $D=4$), with the only exception of a vertex quadratic in $\hat{\eta}$, which appears with a factor ϕ. The conclusion is that, for our purposes, we can handle the quantum system of $N=1$ super-Yang-Mills coupled to matter scalar superfields in AdS_4 space in a way similar to the corresponding theory in a flat background, with the only difference of including in the Feynman rules the additional quadratic vertex

$$\frac{1}{2}m\phi+\text{H.c.} \qquad (2.11)$$

The residual explicit ϕ dependence of the latter reflects the deviation of the theory from a superconformal one, owing to the introduction of a mass term. A remarkable feature of the above rescaling, which effectively removes ϕ from the superconformal invariant part of the superfield action (with the caveat of possible anomalous contributions [33]), is that it takes automatically into consideration the need to resort to some perturbative approach in the effects of the curvature of the background space, leading us naturally to introduce the above implicit expansion in the compensator ϕ.

III. THE RENORMALIZATION OF THE GRAVITATIONAL ACTION

Here we describe our main result in an attempt to explain the shadowing of the cosmological constant. We build upon our previous work [33,34] and carry out the renormalization of the gravitational action induced by $\text{OSp}(N,4)$-invariant matter multiplets in curved space. Here great care is needed, as every sign at each step is crucial.
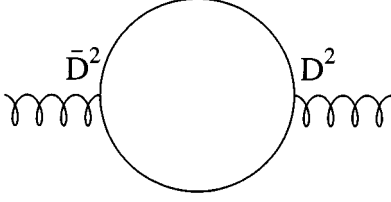
FIG. 1. One-loop renormalization of gravitational actions from chiral matter (the $\hat{\eta}$ solid lines): the curly lines denote the compensator ϕ and its H.c. $\overline{\phi}$.

We start by discussing the renormalization of the gravitational action induced by a matter chiral superfield, through the presence of the Feynman diagram in Fig. 1. This yields the following divergent contribution:

$$am^2 \frac{1}{\epsilon} \int d^4x d^4\theta \phi \overline{\phi}, \qquad (3.1)$$

with $a>0$ in any case. One can doubt about the fact of interpreting this diagram as a renormalization of the pure supergravity action

$$-\frac{3}{\kappa^2} \int d^4x d^4\theta \phi \overline{\phi}, \qquad (3.2)$$

or the cosmological superspace action

$$\frac{\alpha}{\kappa^2} \int d^4x d^2\theta \phi^3 + \text{H.c.} \qquad (3.3)$$

There are clear reasons to support this point of view, which goes in the direction of considering the above contribution as a renormalization of the second term, i.e., the cosmological action. In order to pursue this idea we can translate the result in components to read

$$am^2 \alpha_R^2 \frac{1}{\epsilon} \int d^4x \sqrt{-g}. \qquad (3.4)$$

Here, and throughout this section, we denote with subscript indices R,B the renormalized and the bare parameters of the action, respectively.

On the other hand, we know that the final expression for the gravitational actions has to be

$$S = -\frac{1}{\kappa^2} \int d^4x \sqrt{-g} R(r) + \frac{6}{\kappa^2} \int d^4x \sqrt{-g} \alpha_R^2, \qquad (3.5)$$

where, let us say, R is restricted to metrics of the anti−de Sitter type, with arbitrary radius r. Then it is clear to us that the contribution (3.4) can be consistently interpreted as a renormalization of only the cosmological term, i.e., the second term on the right-hand side (RHS) of Eq. (3.5). This interpretation is also compatible with the fact that the renormalization of the higher order gravitational tensors (i.e., the contributions of order α_R^4) does not take place. In our opinion, this is not a coincidence, rather it means that any
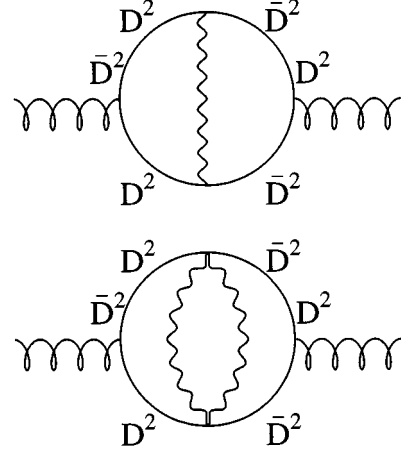

FIG. 2. $O(g^2)$ two-loop corrections to the cosmological renormalization factor Z_α in (3.7); the wavy lines denote the gauge superfields.

vacuum diagram can be seen as a renormalization of only one parameter in the gravitational action, namely, α_R.

So, let us write then

$$am^2 \frac{1}{\epsilon} \int d^4x d^4\theta \phi \overline{\phi} = am^2 \frac{1}{\epsilon} \alpha_R \int d^4x d^2\theta \phi^3, \quad (3.6)$$

where we make use of the equation of motion (2.1) for the chiral compensator superfield ϕ. Considering this contribution alone, we would have

$$\alpha_B = \alpha_R - am^2 \alpha_R \frac{1}{\epsilon} \kappa^2 \equiv Z_\alpha \alpha_R, \qquad (3.7)$$

but then it is obvious that this renormalization factor Z_α cannot lead to a running α_R, since the above m and κ parameters do not depend on the renormalization scale μ (the m parameter that appears here cannot be anything, other than the bare mass m_B).

In this way, we reach the conclusion that the relevant term in Z_α is the two-loop contribution to the vacuum, which forces us to introduce gauge interactions. In fact, when introducing gauge interactions, of all the plethora of diagrams that one can imagine to the order g^2 in the gauge coupling constant, we believe there are only two that survive, after performing the D algebra of the covariant derivatives, i.e., the diagrams in Fig. 2. We give, in the following, the computation of the first graph (at the top of Fig. 2), as the other one (at the bottom of Fig. 2) looks really frightening to compute, and anyhow it cannot change the conclusion of this story. Using integration by parts for the superspace covariant derivatives and discarding a finite remainder given in Fig. 3, one can identify the divergent parts of the first graph in Fig. 2 and the diagram in Fig. 4.

The diagram in Fig. 4 yields the amplitude $\mathcal{V}$:

$$\mathcal{V} = m^2 g^2 \int \frac{d^Dk}{(2\pi)^D} \int \frac{d^Dq}{(2\pi)^D} \frac{1}{k^2} \frac{1}{q^2} \frac{1}{(k-p)^2} \frac{1}{(k+q)^2}, \qquad (3.8)$$
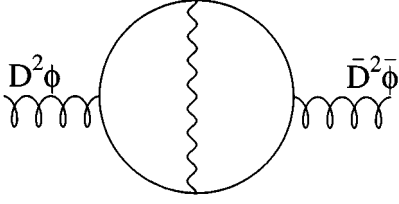
FIG. 3. A finite contribution to the gravitational action produced, when integrating by parts the superspace covariant derivatives in the Feynman diagram at the top of Fig. 2, in order to obtain the graph of Fig. 4.

where we work in dimension $D = 4 - \epsilon$. The first integration can be carried out, yielding the amplitude

$$\mathcal{V} = m^2 g^2 \frac{1}{(2\pi)^{4-\epsilon}} \pi^{2-\epsilon/2} \Gamma\left(\frac{\epsilon}{2}\right) B\left(1 - \frac{\epsilon}{2}, 1 - \frac{\epsilon}{2}\right) \mathcal{I},$$

(3.9)

where

$$B(1-z, 1-w) \equiv \int_0^1 dy \frac{1}{y^z (1-y)^w}$$

(3.10)

and we define the momentum integral

$$\mathcal{I} \equiv \int \frac{d^D k}{(2\pi)^D} \frac{1}{k^2} \frac{1}{(k-p)^2} \frac{1}{(k^2)^{\epsilon/2}}.$$

(3.11)

Also the integral $\mathcal{I}$ can be evaluated. Hence, using standard properties of the B function, we get the result

$$\mathcal{V} = m^2 g_R^2 \frac{1}{(2\pi)^{8-2\epsilon}} \pi^{4-\epsilon} \frac{2}{\epsilon} \Gamma(\epsilon) \frac{1}{1-\epsilon}$$

$$\times \left[\Gamma\left(1 - \frac{\epsilon}{2}\right) \right]^3 \frac{1}{\Gamma(2 - 3\epsilon/2)} \left(\frac{\mu^2}{p^2}\right)^\epsilon.$$

(3.12)

Here we introduced the expression of the renormalized gauge coupling constant g_R in terms of the renormalization scale
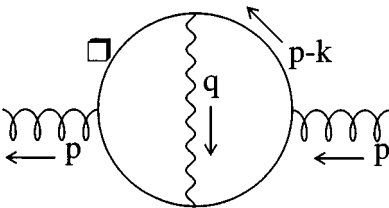
$$g_R = g \mu^{-\epsilon}.$$

(3.13)

FIG. 4. $O(g^2)$ two-loop Feynman diagram corresponding to the amplitude $\mathcal{V}$ in Eq. (3.8); the two independent momenta (loop variables) of the internal lines are explicitly indicated, together with the momenta of the external compensators.

Two remarks are in order about this result. First of all, the leading divergence is of order ϵ^{-2}, which makes it relevant for the purpose of running α_R. Second, it contains subleading nonlocal divergencies of the form $\ln(p/\mu)$. We will have to prove that they cancel with similar terms coming from the other graph in Fig. 2, for this whole thing to make sense. What is important about this other graph, apart from the cancellation of the nonlocal divergencies, is that the leading divergence comes with a positive sign (as is the case for the amplitude $\mathcal{V}$ computed above). So let us assume that the whole contribution from Fig. 2 is of the form

$$b m^2 g_R^2 \alpha_R \left(\frac{1}{\epsilon}\right)^2 \int d^4 x d^2 \theta \phi^3,$$

(3.14)

with $b > 0$, plus perhaps $1/\epsilon$ subleading divergencies that are not relevant, when studying the renormalization group equation for α_R to the order g_R^2. Then, we have that

$$Z_\alpha = 1 - a m^2 \kappa^2 \frac{1}{\epsilon} - b m^2 \kappa^2 g_R^2 \left(\frac{1}{\epsilon}\right)^2,$$

(3.15)

$$\frac{d\alpha_B}{d\mu} = 0 = Z_\alpha \frac{d\alpha_R}{d\mu} + \frac{dZ_\alpha}{d\mu} \alpha_R,$$

(3.16)

$$\frac{d\alpha_R}{d\mu} = -\frac{1}{Z_\alpha} \frac{dZ_\alpha}{d\mu} \alpha_R$$

$$\approx -\frac{1}{1 - a m^2 \kappa^2/\epsilon} \left[-2 b m^2 \kappa^2 g_R \frac{dg_R}{d\mu} \left(\frac{1}{\epsilon}\right)^2 \right] \alpha_R$$

$$+ O(g_R^3)$$

$$= -\frac{1}{1 - a m^2 \kappa^2/\epsilon} \left(b m^2 \kappa^2 g_R^2 \frac{1}{\mu} \frac{1}{\epsilon} \right) \alpha_R + O(g_R^3).$$

(3.17)

In the limit $\epsilon \to 0$, we have then

$$\frac{d\alpha_R}{d\mu} = \frac{b}{a} g_R^2 \frac{1}{\mu} \alpha_R,$$

(3.18)

so that one can easily guess the kind of theories in which the effective cosmological constant goes to zero in the infrared as a power of μ. We can then write that, if $g_R^2 = \text{const}$,

$$\alpha_R = \alpha_0 \left(\frac{\mu}{\mu_0}\right)^{b g_R^2/a}.$$

(3.19)

The important thing, which we checked repeatedly, is that the exponent is *positive*, though under which circumstances it could be bigger than 1 it is more obscure.

IV. CONCLUSION

We cannot add much more to the above considerations apart from the fact that, if the interpretation $\mu^2 \approx R$ in a gravitational measurement is plausible, then this could explain why the effect of the cosmological constant is always

shadowed, no matter what the value of α_0 might be, by the effects of Newtonian gravity.

The calculation of the stress-tensor anomaly in AdS$_4$ supersymmetry showed [36] that the choice of the vacuum, around which the model is quantized, does not affect the renormalization of the purely geometrical tensors in the effective action, nor the trace anomaly induced by matter multiplets invariant under the supersymmetry group OSp(N,4) in curved space. These quantities are independent also from the boundary conditions for the free-field propagators, as proven in Ref. [36]. We remark, in passing, that the conformal anomaly, as an integrability condition for the supersymmetric σ models corresponding to superstring theories, was obtained in Refs. [37–41].

Next, we wish to compare the work of Elizalde and Odintsov (EO) [42] with our result (3.19) on the running and the consequent exponential shadowing of the cosmological constant. The two results appear to be similar, with ours as a particular case, at least at first sight. Notice, however, that our paper is of wider interest in what concerns keeping a global (anti–de Sitter) supersymmetry of the curved background space. Indeed our work and [42] may be considered as complementary, in many respects.

The paper by EO contains a phenomenological analysis of theories that are finite in flat spacetime. A class of such theories is considered in [42] interacting with an external gravitational field (including a nonminimal term linear in the curvature and quadratic in the massless scalar matter field). Our calculation applies specifically to the (globally) supersymmetric background (ground state) solution for supersymmetric (SUSY) grant unified theories (GUTs), extended supergravity (SUGRA), superconformal invariant theories (superstrings). As in our calculation the (global) supersymmetry of this ground state solution is maintained, consequently the contribution of the supersymmetric particles is determined explicitly.

In addition, the role of a nonvanishing mass for the matter fields is perfectly clear and is included in our analysis (EO only consider a massless theory). Finally, in EO a fine-tuning is needed, in order to avoid in their Eq. (19) the unrealistic growth of Newton's constant, and/or to obtain the screening of the cosmological constant. We obtain such screening for α_R directly and without fine-tuning.

We stress that we consider the superconformal invariant theories to be minimally interacting with the external gravitational background. Our main motivation is to keep the global supersymmetry of the background, which requires referring to the minimal $n = -1/3$ compensator ϕ. It was shown long ago in Ref. [35] (p. 336, Sec. 5.7) that nonminimal n implies the spontaneous breakdown of $N=1$ supersymmetry.

In our paper we have recalled several developments that simplified the procedure in practical calculations for the renormalization of our class of theories. In particular the superconformal rescaling is an improvement that can be useful in future applications. In this respect we hope that the discussion of the method given in this work will find further use.

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