Dimensional Reduction of the Heterotic Supergravity on nearly-Kähler cosets*

Athanasios Chatzistavrakidis†
Institute of Nuclear Physics, NCSR Demokritos,
GR-15310, Ag. Paraskevi, Athens, GREECE

Pantelis Manousselis‡
Physics Department, National Technical University of Athens,
GR-15780 Zografou Campus, Athens, GREECE

George Zoupanos§
Physics Department, National Technical University of Athens,
GR-15780 Zografou Campus, Athens, GREECE

Abstract

We study the reduction of the $\mathcal{N}=1$, ten-dimensional Heterotic Supergravity to four dimensions, at zeroth order in α', when the internal space is a coset manifold. Then we apply the formalism using as internal manifolds all possible six-dimensional non-symmetric coset spaces, namely $G_2/SU(3), Sp_4/SU(2)\times U(1)$ and $SU(3)/U(1)\times U(1)$, which are known to preserve the supersymmetric spectrum. The effective actions describe $\mathcal{N}=1$ four-dimensional supergravities of the no-scale type. We find the scalar potentials and comment on the different possibilities concerning the vacua of these models. We note that all the spaces we use admit a nearly-Kähler structure therefore these models constitute specific examples of nearly-Kähler reduction.

1. Introduction

Supergravity theories have been studied extensively over the past thirty years. In particular, exploring the possibility that superstring theories describe the real world, the task of providing a suitable compactification which

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† e-mail address: cthan@mail.ntua.gr
‡ e-mail address: pman@central.ntua.gr
§ e-mail address: zoupanos@cern.ch

would lead to a realistic four-dimensional theory has been pursued in many diverse ways.

The early attempts to reduce such theories made extensive use of Calabi-Yau (CY) manifolds, i.e. manifolds with $SU(3)$ holonomy [1]. A reduction and truncation procedure was developed in [2] and [3], where the internal space is not specified but the general characteristics of CY manifolds are kept. However, there exist some problems with the use of CY in the reduction procedure due to their complicated geometry. Among others their metric is not known explicitly and their Euler characteristic is too large to accommodate an acceptable number of fermion generations. Moreover, in CY compactifications the resulting low-energy field theory in four dimensions contains a number of massless chiral fields, characteristic of the internal geometry, known as moduli. These fields correspond to flat directions of the effective potential and therefore their values are left undetermined. Since these values specify the masses and couplings of the four-dimensional theory, the theory has limited predictive power.

In the context of flux compactifications, recent developments led to the study of a new class of internal spaces called manifolds with $SU(3)$-structure. The case of $SU(3)$-structures is of special interest since the "local Lorentz" (structure) group $SO(6)$ of the internal space can be reduced down to $SU(3)$ in a way that a single spinor can be globally defined on it. This spinor is not necessarily covariantly constant with respect to the Levi-Civita connection, as in the case of Calabi-Yau manifolds, but it can be constant with respect to a connection with torsion. This condition allows for a wider class of internal spaces, such as nearly-Kähler and half-flat manifolds. The Heterotic String theory has been recently studied in this context in [4] and [5]. Simple G-structure manifolds are six dimensional cosets possessing an $SU(3)$ structure. They have been identified as supersymmetric solutions in the case of type II theories (see e.g. [6],[7]). In the studies of compactification of the Heterotic Supergravity the use of cosets spaces was introduced in [8], and recently developed further in [9]-[11]. Particularly, in [11] it was shown that supersymmetric compactifications of the Heterotic String theory of the form $AdS_4 \times S/R$ exist when background fluxes and general condensates are present. In addition, effective theories where constructed in [7], [12], [13] in the case of type II supergravity.

In this contribution we discuss the dimensional reduction of the Heterotic String at zeroth order in α' in the case where the internal manifold is a coset space. In section 2 we provide a brief reminder of the Heterotic Supergravity and discuss the basics of manifolds with $SU(3)$-structure. In section 3 we present the general reduction procedure that we follow and determine the resulting four-dimensional Lagrangian. In section 4 we apply the previously found results in the case of all three six-dimensional non-symmetric coset spaces, which in addition admit a nearly-Kähler structure and we discuss the supergravity description from the four-dimensional point of view. Section 5 contains our conclusions and a discussion on the vacuum of the models we studied, making contact with the results of no-scale supergravity.

2. General Framework

2.1. The Heterotic Supergravity

In this section we briefly review the field content and the Lagrangian of the Heterotic Supergravity in order to fix our notation and conventions.

The field content of the Heterotic Supergravity consists of the $\mathcal{N}=1$, $D=10$ supergravity multiplet, which accommodates the fields g_{MN}, ψ_M, B_{MN}, λ and φ (i.e. the graviton, the gravitino which is a Rarita-Schwinger field, the two-form potential, the dilatino which is a Majorana-Weyl spinor, and the dilaton which is a scalar). Capital Latin letters denote here ten-dimensional indices.

The corresponding Lagrangian of the ten-dimensional $\mathcal{N}=1$ Heterotic Supergravity in the Einstein frame can be written as [14]

$$\mathcal{L}=\mathcal{L}_b+\mathcal{L}_f+\mathcal{L}_{int}, \tag{1}$$

where the different sectors of the theory, ignoring the gauge fields and the gaugini at lowest order, are [1]

$$\hat{e}^{-1}\mathcal{L}_b=-\frac{1}{2\hat{\kappa}^2}(\hat{R}\hat{*}\mathbf{1}+\frac{1}{2}e^{-\hat{\phi}}\hat{H}_{(3)}\wedge\hat{*}\hat{H}_{(3)}+\frac{1}{2}d\hat{\phi}\wedge\hat{*}d\hat{\phi}), \tag{2}$$

$$\hat{e}^{-1}\mathcal{L}_f=-\frac{1}{2}\hat{\bar{\psi}}_M\hat{\Gamma}^{MNP}D_N\hat{\psi}_P-\frac{1}{2}\hat{\bar{\lambda}}\hat{\Gamma}^M D_M\hat{\lambda}, \tag{3}$$

$$\begin{aligned}\hat{e}^{-1}\mathcal{L}_{int}=&-\frac{1}{2}\hat{\bar{\psi}}_M\hat{\Gamma}^N\hat{\Gamma}^M\hat{\lambda}\partial_N\hat{\phi}+e^{-\hat{\phi}/2}\hat{H}_{PQR}\Big(\hat{\bar{\psi}}_M\hat{\Gamma}^{MPQRN}\hat{\psi}_N\\&+6\hat{\bar{\psi}}^P\hat{\Gamma}^Q\hat{\psi}^R-\sqrt{2}\bar{\psi}_M\hat{\Gamma}^{PQR}\hat{\Gamma}^M\hat{\lambda})\Big)+\text{four-fermion terms},\end{aligned} \tag{4}$$

where we have placed hats in all the ten dimensional fields to distinguish them from their four dimensional counterparts which will appear after the reduction. The three-form $\hat{H}_{(3)}$ is the field strength for the B-field, namely $\hat{H}=d\hat{B}$. The Gamma matrices are the generators of the ten dimensional Clifford algebra, hence we place hats on them too, while those with more than one index denote antisymmetric products of Γs. Also $\hat{\kappa}$ is the gravitational coupling constant in ten dimensions, with dimensions $(\text{mass})^4$. $\hat{e}$ is the determinant of the metric, while $\hat{*}$ is the Hodge star operator in ten dimensions. Let us mention that the gravitational constant is defined as $\hat{\kappa}^2=8\pi G_N$, where G_N is the Newton constant. As such its relation to the Planck mass is $\hat{\kappa}=\frac{1}{m_{Pl}}$, since the Planck mass is $m_{Pl}=\frac{1}{\sqrt{8\pi G_N}}$.

2.2. Manifolds with $SU(3)$-structure

Calabi-Yau manifolds were proposed as internal spaces for compactifications in view of the requirement that a four-dimensional $\mathcal{N}=1$ supersymmetry is preserved. Namely they admit a single globally defined spinor,

[1] Here we use differential form notation for the kinetic terms of the bosons, which will prove to be useful in the course of the reduction.

which is covariantly constant with respect to the (torsionless) Levi-Civita connection. However, there exists a larger class of manifolds for which the spinor is covariantly constant with respect to a connection with torsion. These are called manifolds with $SU(3)$ structure and clearly Calabi-Yau manifolds are a subclass in the category of $SU(3)$-structure manifolds.

In particular, in order to define a single spinor on a six-dimensional manifold one has to reduce the structure group $SO(6)$. The simplest procedure one can follow is to reduce this group to $SU(3)$, since then the decomposition of the spinor of $SO(6)$ reads $\mathbf{4} = \mathbf{3} + \mathbf{1}$ and the spinor we are looking for is the singlet, let us call it η. Then, we can use η to define the $SU(3)$-structure forms, which are a two-form J and a three-form Ω defined as

$$J_{mn} = \mp i\eta_{\pm}^{\dagger}\gamma_{mn}\eta_{\pm}, \quad \Omega_{mnp} = \eta_{-}^{\dagger}\gamma_{mnp}\eta_{+}, \quad \Omega^{*}_{mnp} = -\eta_{+}^{\dagger}\gamma_{mnp}\eta_{-}. \tag{5}$$

These forms are not closed forms but instead they satisfy

$$\begin{aligned} dJ &= \frac{3}{4}i(\mathcal{W}_1\Omega^{*} - \overline{\mathcal{W}}_1\Omega) + \mathcal{W}_4 \wedge J + \mathcal{W}_3, \\ d\Omega &= \mathcal{W}_1 J \wedge J + \mathcal{W}_2 \wedge J + \mathcal{W}_5^{*} \wedge \Omega. \end{aligned} \tag{6}$$

The expressions (6) define the five intrinsic torsion classes, which are a zero-form $\mathcal{W}_1$, a two-form $\mathcal{W}_2$, a three-form $\mathcal{W}_3$ and two one-forms $\mathcal{W}_4$ and $\mathcal{W}_5$. These classes completely characterize the intrinsic torsion of the manifold.

The classes $\mathcal{W}_1$ and $\mathcal{W}_2$ are vanishing when the manifold is complex and furthermore a Kähler manifold has vanishing $\mathcal{W}_3$ and $\mathcal{W}_4$. A Calabi-Yau manifold has all the torsion classes equal to zero. An interesting class of $SU(3)$-structure manifolds are called nearly-Kähler manifolds. In this case all the torsion classes but $\mathcal{W}_1$ are vanishing. This suggests that the manifold is not Kähler and not even complex. A complete list of other classes of $SU(3)$-structure manifolds can be found in [15].

Manifolds with $SU(3)$-structure in general [4],[5] and nearly-Kähler manifolds in particular [9],[12] have attracted a lot of interest in flux compactifications over the last years. Since we are interested in coset spaces it is worth noting that certain six-dimensional coset spaces provide simple realizations, and to our knowledge the only ones, of manifolds with a nearly-Kähler structure. These spaces are the three non-symmetric six-dimensional cosets, namely $G_2/SU(3), Sp_4/(SU(2) \times U(1))_{non-max}$ and $SU(3)/U(1) \times U(1)$, plus the coset $SU(2)^3/SU(2)$. It is therefore interesting to perform an explicit reduction over these spaces and find the resulting effective actions, a task which we shall perform in the forthcoming sections.

3. Reduction Procedure

In the present section we focus on the bosonic part of the Heterotic Supergravity Lagrangian and perform its reduction from ten to four dimensions over the coset spaces S/R. Since the Kähler potential and the superpotential of the four-dimensional theory can be obtained from the bosonic part, this procedure will be sufficient to find the supergravity description in four dimensions.

In order to reduce the theory we need ansatz for the bosons, namely for the metric, the dilaton and the B-field. Starting with the metric, our ansatz reads

$$d\hat{s}^2_{(10)} = e^{2\alpha\varphi(x)}\eta_{mn}e^m e^n + e^{2\beta\varphi(x)}\gamma_{ab}(x)e^a e^b, \tag{7}$$

where m, n are four-dimensional indices and a, b are flat internal indices. In this ansatz we included exponentials which rescale the metric components. This is always needed in order to obtain an action without any prefactor for the Einstein-Hilbert part. We shall see that we need to specify the values of α and β in order to fulfil this requirement.

Let us note that this ansatz is dictated by two further requirements. Firstly, the metric is required to be S-invariant [16]. Secondly, the requirement of consistency of our reduction enforces the vanishing of the Kaluza-Klein fields and allows only the scalar fluctuations [17],[18].

Following the standard procedure of reducing this action in the case of a coset space (see e.g. [19]) and choosing $\alpha = -\frac{\sqrt{3}}{4}, \beta = -\frac{\alpha}{3}$, we find that the reduced Lagrangian reads

$$\mathcal{L} = -\frac{1}{2\kappa^2}\left(R * \mathbf{1} - P_{ab} \wedge * P_{ab} + \frac{1}{2}d\varphi \wedge * d\varphi\right) - V, \tag{8}$$

with the potential V having the form

$$V = -\frac{1}{8\kappa^2}e^{2(\alpha-\beta)\varphi}(\gamma_{ab}\gamma^{cd}\gamma^{ef}f^a_{\ \ ce}f^b_{\ \ df} + 2\gamma^{ab}f^c_{\ \ da}f^d_{\ \ cb} + 4\gamma^{ab}f_{iac}f^{ic}_b) * \mathbf{1}, \tag{9}$$

where the index i runs in R. In the reduced Lagrangian the fields P_{ab} are defined as

$$P_{ab} = \frac{1}{2}[(\Phi^{-1})^c_a d\Phi^b_c + (\Phi^{-1})^c_b d\Phi^a_c], \tag{10}$$

with

$$\gamma_{cd} = \delta_{ab}\Phi^a_c\Phi^b_d, \tag{11}$$

while the corresponding kinetic term in (8) will provide the kinetic terms for the extra scalars apart from φ that can appear through the metric $\gamma_{ab}(x)$.

On the other hand, the higher-dimensional dilaton is trivially reduced by $\hat{\phi}(x, y) = \phi(x)$, since it is already a scalar in ten dimensions. This leads to a kinetic term $-\frac{1}{4\kappa^2}d\phi \wedge * d\phi$ in the reduced Lagrangian.

Finally, concerning the three-form sector of the theory we expand the B-field on invariant forms of the coset space, namely our ansatz reads

$$\hat{B} = B(x) + b^i(x)\omega_i(y), \tag{12}$$

where the index i counts the number of invariant two-forms. Then it is straightforward to see that the higher-dimensional three-form can be written in terms of four-dimensional fields as

$$\hat{H} = dB + db^i \wedge \omega_i + b^i d\omega_i. \tag{13}$$

Let us note here that unlike the case of CY compactifications, where the expansion forms are harmonic and hence closed, here we expand in forms that are not closed thus there is an extra term appearing in eqn. (13).

Plugging the ansatz into the corresponding term in the Lagrangian we find that the reduced Lagrangian for this sector is

$$\begin{aligned}\mathcal{L} = -\frac{1}{2\kappa^2}e^{-\phi}\Big[&\frac{1}{2}e^{-4\alpha\varphi}d\theta \wedge *d\theta + \frac{m}{2}e^{-4\beta\varphi}db^i \wedge *db^i \\ &+e^{4\alpha\varphi}(n_1(b^i)^2 + n_2\epsilon_{ij}b^ib^j)vol_4\Big]\wedge vol_6,\end{aligned} \tag{14}$$

where θ is the pseudoscalar obtained by duality transformation on dB, while m, n_1 and n_2 are fixed constants defined by

$$\omega_i \wedge *\omega_j = m\delta_{ij}vol_6, \quad d\omega_i \wedge *d\omega_j = (n_1\delta_{ij} + n_2\epsilon_{ij})vol_6. \tag{15}$$

4. Examples

In this section we specialize the general discussion of the previous one in the case of non-symmetric six-dimensional coset spaces. There are known to exist three cases, namely $G_2/SU(3)$, $Sp_4/(SU(2) \times U(1))_{non-max}$ and $SU(3)/U(1) \times U(1)$.

Geometry and $SU(3)$-structure

$G_2/SU(3)$: This manifold has one G_2-invariant two-form, $\omega_1 = e^{12} - e^{34} - e^{56}$, and two G_2-invariant three forms, $\rho_1 = e^{136} + e^{145} - e^{235} + e^{246}$ and $\rho_2 = e^{135} - e^{146} + e^{236} + e^{245}$ [20]. On the other hand there are no invariant one-forms. The invariant forms of the coset space are intimately connected to its $SU(3)$-structure forms J and Ω. Indeed

$$J = R\omega_1, \quad \Omega = \sqrt{R^3}(\rho_2 + i\rho_1), \tag{16}$$

where R is the radius of the space.

The only non-vanishing torsion class is

$$\mathcal{W}_1 = -\frac{2i}{\sqrt{3R}}, \tag{17}$$

which means that this space naturally admits a nearly-Kähler structure without any further conditions.

$Sp_4/(SU(2) \times U(1))_{non-max}$: Here there exist two Sp_4-invariant two-forms $\omega_1 = e^{12} + e^{56}$ and $\omega_2 = e^{34}$ and two three-forms $\rho_1 = e^{136} - e^{145} + e^{235} + e^{246}$ and $\rho_2 = e^{135} + e^{146} - e^{236} + e^{245}$ [20]. As in the previous case no invariant one-forms exist. The structure forms are

$$J = -R_1\omega_1 + R_2\omega_2, \quad \Omega = \sqrt{R_1^2 R_2}(\rho_2 + i\rho_1), \tag{18}$$

with R_1 and R_2 the radii of the space. The non-vanishing torsion classes in this case are

$$\mathcal{W}_1 = -\frac{2i}{3}\frac{2R_1 + R_2}{\sqrt{R_1^2 R_2}}, \tag{19}$$

$$\mathcal{W}_2 = -\frac{4i}{3}\frac{1}{\sqrt{R_1^2 R_2}}\left[R_1(R_1 - R_2)e^{12} - 2R_2(R_2 - R_1)e^{34} + R_1(R_1 - R_2)e^{56}\right] \tag{20}$$

and it is obvious that the space has a nearly-Kähler limit when the condition $R_1 = R_2$ is satisfied.

$SU(3)/U(1) \times U(1)$: The coset space $SU(3)/U(1) \times U(1)$ has three $SU(3)$-invariant two-forms $\omega_1 = e^{12}, \omega_2 = e^{45}, \omega_3 = e^{67}$, two invariant three-forms $\rho_1 = e^{147} - e^{156} + e^{246} + e^{257}, \rho_2 = e^{146} + e^{157} - e^{247} + e^{256}$ and no invariant one-forms [20]. The forms which specify the $SU(3)$-structure are

$$J = -R_1\omega_1 + R_2\omega_2 - R_3\omega_3, \qquad \Omega = \sqrt{R_1 R_2 R_3}(\rho_2 + i\rho_1), \tag{21}$$

where the three radii of the space are involved, while the torsion classes are

$$\mathcal{W}_1 = -\frac{2i}{3}\frac{R_1 + R_2 + R_3}{\sqrt{R_1 R_2 R_3}}, \tag{22}$$

$$\begin{aligned}\mathcal{W}_2 = -\frac{4i}{3}\frac{1}{\sqrt{R_1 R_2 R_3}}&[R_1(2R_1 - R_2 - R_3)e^{12} - R_2(2R_2 - R_1 - R_3)e^{34} \\ &+ R_3(2R_3 - R_1 - R_2)e^{56}].\end{aligned} \tag{23}$$

It is again straightforward that under the condition of equal radii this space admits a nearly-Kähler structure.

Supergravity description in four dimensions

$G_2/SU(3)$: For $G_2/SU(3)$ the most general G_2-invariant metric is $g_{ab} = e^{2\beta\varphi}\delta_{ab}$, namely there is only one scale. Thus, using the fact that $\gamma_{ab} = \delta_{ab}$ as well as the structure constants of this coset space, we easily find that the four dimensional potential in this case is

$$V = -\frac{1}{2\kappa^2}(6e^{\frac{8\alpha}{3}\varphi} - 6e^{-\phi + 4\alpha\varphi}b^2). \tag{24}$$

Note that in eq (24) we have calculated that n_1 in eq (15) becomes 12 in the present case. In order to bring the reduced Lagrangian in the standard four-dimensional supergravity form we define the complex superfields, consisting of all the scalar moduli,

$$S = e^{\phi_0} + i\lambda, \qquad T = e^{-\varphi_0/\sqrt{3}} + ib. \tag{25}$$

Here for convenience we have redefined two of the moduli as

$$\phi_0 = \frac{1}{2}(-\phi - 4\alpha\varphi), \qquad \varphi_0 = \frac{1}{2}(-\varphi + 4\alpha\phi). \tag{26}$$

Then we define the Kähler potential and the superpotential

$$K = \frac{1}{\kappa^2}\left[-ln(S+\overline{S}) - 3ln(T+\overline{T})\right], \qquad W = 4\sqrt{3}T. \tag{27}$$

Consequently, we can easily verify that the four-dimensional potential is indeed given by the supergravity expression

$$V(\Phi,\bar{\Phi}) = \frac{1}{\kappa^4}e^{\kappa^2 K}(K^{i\bar{j}}\frac{DW}{D\Phi^i}\frac{D\overline{W}}{D\bar{\Phi}^{\bar{j}}} - 3\kappa^2 W\overline{W}), \tag{28}$$

where the complex superfields are collectively denoted by Φ and the derivatives involved are the Kähler covariant derivatives. Also one can check that the kinetic terms are also exactly retrieved as

$$-\frac{1}{\kappa^2}K_{i\bar{j}}d\Phi^i \wedge *d\bar{\Phi}^{\bar{j}} \tag{29}$$

with the same Kähler potential, as required by supergravity.

$Sp_4/SU(2)\times U(1)$*:* This coset space admits two independent scales so we need to parametrize the metric by an extra scalar field, say χ. Also, there are now two scalars from the B-field, b_1 and b_2. In the same spirit as in the previous case we find that the four-dimensional supergravity description is obtained by defining the Kähler potential and superpotential

$$K = \frac{1}{\kappa^2}\left[-ln(S+\overline{S}) - ln[(T_1+\overline{T}_1)^2(T_2+\overline{T}_2)]\right], \quad W = 2T_1 + T_2, \tag{30}$$

where the complex superfields are defined as

$$S = e^{\phi_0} + i\lambda, \quad T_1 = e^{-\varphi_0/\sqrt{2}} + ib_1, \quad T_2 = e^{-\chi} + ib_2. \tag{31}$$

The nearly-Kähler limit corresponds to the case $T_1 = T_2$.

$SU(3)/U(1)\times U(1)$: Here there are three independent scales, namely there are three scalars parametrizing the metric fluctuations φ, χ, ψ and on the other hand there are three fields b_1, b_2, b_3 in this case.

The same logic leads in this case to the Kähler potential

$$K = \frac{1}{\kappa^2}\left[-ln(S+\overline{S}) - ln[(T_1+\overline{T}_1)(T_2+\overline{T}_2)(T_3+\overline{T}_3)]\right] \qquad (32)$$

and superpotential

$$W = -T_1 + T_2 - T_3, \qquad (33)$$

with $T_1 = e^{-\varphi} + ib_1, T_2 = e^{-\chi} + ib_2, T_3 = e^{-\psi} + ib_3$. The nearly-Kähler limit is obtained again when $T_1 = T_2 = T_3$.

5. Discussion and Conclusions

In this contribution we explicitly reduced the Heterotic Supergravity from ten dimensions to four at zeroth order in α', i.e. ignoring the gauge fields, using coset spaces as internal manifolds. We have examined three specific models based on the three six-dimensional non-symmetric coset spaces and we have determined the resulting four-dimensional effective actions.

From our results we observe that in all three cases the structure of the Kähler potential and the superpotential of the resulting four-dimensional supergravities is the same and they differ only on the number of scalar moduli appearing in each case. In particular, the volume of the internal space is parametrized by one scalar field for $G_2/SU(3)$, while in the cases of $Sp_4/(SU(2\times U(1))_{non-max}$ and $SU(3)/U(1)\times U(1)$ the volume depends on two and three scalar fields respectively. On the other hand, the scalar fields emerging from the internal components of the B-field are one, two and three respectively in the three examples we studied. It is worth noting that the structure of the Kähler potential is exactly the one that appears in the no-scale models of supergravity [21].

As far as the vacuum of the models is concerned, as long as supersymmetry remains unbroken it is impossible to find a de Sitter vacuum [22]. On the other hand, taking as example the model based on $G_2/SU(3)$ we see that a Minkowski vacuum could be obtained only when the fields acquire the expectation values $< b >= 0$ and $< \varphi >\rightarrow \infty$, which is not a plausible condition since it leads to a run-away argument. In particular, the condition $< \varphi >\rightarrow \infty$ in the vacuum means that the volume of the internal space tends to infinity. The same result holds in the other two cases.

In order to enrich our models and look for realistic phenomenology it is necessary to work on the next order in α' and include gauge fields and non-vanishing fluxes [23]. Then, working in the context of the Coset Space Dimensional Reduction [24] we shall be able to find interesting supergravity GUTs in four dimensions with an appropriate number of fermion generations and soft breaking of supersymmetry. Studying the enhanced potential and exploring Minkowski vacua in this case one can hope to be led to positive results. This work is currently in progress.

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