Compactifications of M-theory and their Phenomenological Consequences

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Abstract

We compactify the M-theory proposed by Horava and Witten on a Calabi-Yau manifold with boundary $S_1/Z_2$. A no-scale-like Kähler potential, the superpotential, and the gauge kinetic function are obtained in this 4-dimensional $E_6 \times E_8$ model. We also study the general phenomenological consequences of the resulting M-theory-inspired model, which may include very light gravitinos, axions, and axinos.

April 1997
1 Introduction

Recently, Horava and Witten [1] presented a systematic analysis of eleven-dimensional supergravity on a manifold with boundary, that is related to the strong coupling limit of the $E_8 \times E_8$ heterotic string. In this novel theory, many previous successes based on conventional weakly-coupled string theory may be preserved. In addition, Witten [2] offered an explanation of why Newton’s constant appears to be so small when the 4-dimensional grand unified gauge coupling ($\alpha_{\text{GUT}}$) takes experimentally acceptable values. It follows that the strengths of all interactions, including gravitation, may be naturally unified at the GUT scale, unlike the case of the weakly-coupled heterotic string. Many other interesting implications have been studied, such as gluino condensation and supersymmetry breaking [3], the strong CP problem [4], threshold scale and strong coupling effects [5, 6], and phenomenological consequences [7] which include a constrained sparticle spectrum within the reach of present-generation particle accelerators.

In this paper we compactify this theory on a Calabi-Yau manifold with Hodge numbers $h_{(1,1)} = 1$ and $h_{(2,1)} = 0$ and boundary $S_1/Z_2$. A no-scale-like Kähler potential [8], the superpotential, and the gauge kinetic function are obtained explicitly. In four dimensions this result is related to the previous weakly-coupled string no-scale supergravity result by Witten [9] through a field transformation (Sec. 2), which means that they are equivalent in four dimensions. One might then argue that many heterotic string models obtained previously may also exist in the M-theory regime.

In addition, we consider the physical couplings and scales in the Einstein frame, the eleven-dimensional metric and fivebrane units, and give the intermediate supersymmetry-breaking scale determined by the size of the eleventh dimension of this theory [3, 8, 10] for different grand unified theories (Sec. 3). We also argue that there may exist a very light gravitino in this scenario, which may explain the $e e \gamma \gamma + E_T^{\text{miss}}$ event [11] observed by the CDF Collaboration, and may have some further implications at LEP 2. Finally, we comment on very light axions and axinos that might exist in this scenario.

2 Formal derivation

We follow the notation of Ref. [1], in which the bosonic part of the eleven-dimensional supergravity Lagrangian is given by

$$L_B = \frac{1}{\kappa^2} \int_{M^{11}} d^{11}x \sqrt{g} \left( -\frac{1}{2} R - \frac{1}{48} G_{IJKL} G^{IJKL} - \frac{\sqrt{2}}{3456} \epsilon^{I_1 I_2 \cdots I_{11}} C_{I_1 I_2 I_3} G_{I_4 \cdots I_7} G_{I_8 \cdots I_{11}} \right)$$

$$- \sum_{i=1,2} \frac{1}{2\pi (4\pi \kappa^2)^{7/3}} \int_{M^{10}_i} d^{10}x \sqrt{g} \frac{1}{4} F_{ABA}^a F^{aAB}$$ (1)

where $G_{11,ABC} = (\partial_{11} C_{ABC} \pm 23 \text{ permutations}) + \frac{\kappa^2}{\sqrt{2} \lambda^2} \delta(x^{11}) \omega_{ABC}$, $\lambda^2 = 2\pi (4\pi \kappa^2)^{2/3}$, and the gauge group at the boundary is $E_8 \times E_8$. 


To perform the dimensional reduction to five dimensions (with the 4-dimensional boundary where the Yang-Mills fields live) under the Calabi-Yau manifold with Hodge numbers $h_{(1,1)} = 1$ and $h_{(2,1)} = 0$, we follow Refs. [9, 12] and keep only the SU(3) singlets in the internal indices:

$$g_{\mu\nu} \rightarrow g_{\mu\nu} ; \quad g_{ij} = e^\sigma \delta_{ij}$$

(2)

$$C_{\mu \bar{i} j} = i C_{\mu} \delta_{\bar{i} j} ; \quad C_{ijk} = C'_{ijk}$$

(3)

Then, performing the Weyl rescaling:

$$g_{\mu\nu} \rightarrow e^{-2\sigma} g_{\mu\nu}$$

(4)

and just paying attention to the observable sector, which we assume is at the boundary $x^{11} = 0$, we obtain the canonically normalized Einstein action:

$$L_B = \frac{1}{\kappa^2} \int_{M^6} d^6x \sqrt{g} \left( -\frac{1}{2} R - \frac{9}{4} (\partial_{\mu} \sigma)^2 - \frac{1}{48} e^{6\sigma} G_{\mu\nu\rho\sigma} G^{\mu\nu\rho\sigma} - 27 f_{\mu\nu} f^{\mu\nu} - 36 e^{-3\sigma} |\partial_{\mu} \hat{C}'|^2 
- 54 \sqrt{2} e^{6\sigma \delta} C_{\mu} f_{\nu\rho} f_{\sigma \delta} + \frac{3}{4} \sqrt{2} i \epsilon_{\mu \nu \rho \sigma \delta} \hat{C}^{\rho} \partial_{\mu} \hat{C}' G_{\nu \rho \sigma \delta} \right) 
- \frac{V}{2\pi (4\pi \kappa^2)^2} \int_{M^4} d^4x \sqrt{g} \left( -\frac{1}{4} e^{3\sigma} f Tr[F_{\mu\nu} F^{\mu\nu}] - 3 D_{\mu} C^* D^{\mu} C^x - \frac{8}{3} e^{-3\sigma} |\partial W'|^2 
- \frac{9}{2} f e^{-3\sigma} \sum_i (C^*, \lambda^i C)^2 \right)$$

(5)

where

$$W' = d_{xyz} C^x C^y C^z$$

(6)

$$f_{5\mu} = \partial_5 C_{\mu} - \partial_{\mu} C_5 + i \kappa^2 \delta(x^5) (C_x \leftrightarrow D_{\mu} C^x)$$

(7)

$$\partial_{\mu} \hat{C}' = \partial_{\mu} C' + \frac{\sqrt{2}}{3} \delta_{\mu 5} \delta(x^5) \kappa^2 \lambda^2 W'$$

(8)

The definitions of the $C_x$, $d_{xyz}$ and $f$ are the same as in the Ref. [9], and $V$ is the coordinate volume of the Calabi-Yau manifold, i.e., $V = \int d^6x$. In addition, the (observable) gauge group at the boundary ($x^{11} = 0$ or $x^5 = 0$) is now $E_6$ with spin connection embedding, as in Ref. [9].

If we define a pseudoscalar $\tilde{D}$ by a duality transformation:

$$\frac{1}{4!} e^{6\sigma} G_{\mu\nu\rho\sigma} = \epsilon_{\mu\nu\rho\sigma}(\partial^\delta D + \frac{3}{4} \sqrt{2i} \hat{C}' \partial^\delta \hat{C}')$$

(9)
we have the following Kähler potential in the five-dimensional bulk:

\[ K = -\ln [S + \bar{S} - 72 \bar{C}'C'] \] (10)

where

\[ S = e^{3\sigma} + i24\sqrt{2}D + 36 \bar{C}'C' \] (11)

This Kähler potential parametrizes the $SU(2,1)$ quaternionic manifold \[12, 13\]. Also, $D$ is the invisible axion.

We now compactify the above 5-dimensional-with-boundary Lagrangian on $S_1/Z_2$. For the fields with 11-dimensional origin, i.e., the fields in the bulk, we keep only the zero modes, and considering the boundary condition we can expand $\delta(x^5)$ as:

\[ \delta(x^5) = \frac{1}{2\pi \rho} + \frac{1}{\pi \rho} \sum_{n=1}^{\infty} \cos \frac{n\pi \rho}{\rho} \] (12)

where $\rho$ is the coordinate radius of $S_1$, i.e., $\rho = \frac{1}{2\pi} \int dx^5$. Furthermore, choosing the following elfbein form:

\[ e_A^M = \begin{pmatrix} e_\mu^a & 0 \\ 0 & \varphi \end{pmatrix} \] (13)

and performing the Weyl rescaling:

\[ g_{\mu\nu} \rightarrow \varphi^{-1} g_{\mu\nu} \] (14)

we obtain the 4-dimensional Lagrangian in the Einstein frame (note that at the boundary $C'=0$ and $C_{\mu\nu\rho} = 0$):

\[ L_B = \frac{V}{\kappa^2} 2\pi \rho \int d^4x \sqrt{g} \left( -\frac{1}{2} R - \frac{1}{12} e^{6\sigma} G_{11\mu
u\rho} G^{11\mu\nu\rho} - \frac{9}{4} (\partial_{\mu}\sigma)^2 - \frac{3}{4} \left( \frac{\partial_{\mu}\varphi}{\varphi} \right)^2 + 54 \varphi^{-2} (\partial_\mu C_5 - i \frac{\kappa^2}{6\sqrt{2} 2\pi \rho \lambda^2} C^\alpha \dot{D}_\mu C^\alpha)^2 - \frac{1}{2\pi \rho \lambda^2} [3 \varphi^{-1} D_\mu C^\alpha D^\mu C^\alpha + \frac{1}{4} f e^{3\sigma} \text{Tr} [F_{\mu\nu} F^{\mu\nu}] + \frac{8}{3} \varphi^{-2} e^{-3\sigma} |\partial W'|^2 + 9 \frac{2f}{2} \varphi^{-2} e^{-3\sigma} \sum_i (C^\alpha, \lambda^i C)^2 - 8 \left( \frac{1}{2\pi \rho} \right)^2 \left( \frac{\kappa^4}{\lambda} e^{-3\sigma} \varphi^{-3} |W'|^2 \right) \right) \] (15)

Finally, if we define $g_c^2 = 2\pi \rho \lambda^2 / \kappa^2$ and perform the transformation:

\[ A_\mu \rightarrow g_c A_\mu \; ; \; C_\alpha \rightarrow g_c C_\alpha \] (17)
we obtain the standard supergravity Lagrangian:

\[ L_B = \frac{V}{\kappa^2} 2\pi \rho \int d^4 x \sqrt{g} \left( -\frac{1}{2} R - \frac{1}{12} e^{6\sigma} G_{11\mu\nu} G^{11\mu\nu} - \frac{9}{4} (\partial_\mu \sigma)^2 - \frac{3}{4} \frac{(\partial_\mu \varphi)^2}{\varphi} \right) \\
-54\varphi^{-2}(\partial_\mu C_5) - \frac{i}{6\sqrt{2}} C_x^{*} \tilde{D}_\mu C^x)^2 \\
-3\varphi^{-1} D_\mu C_x^{*} D^\mu C^x - \frac{1}{4} f e^{3\sigma} T r [F_{\mu\nu} F^{\mu\nu}] - \frac{8}{3} g_c^2 \varphi^{-2} e^{-3\sigma} |\partial W'|^2 \\
-\frac{9}{2f} g^2 e^{-2} e^{-3\sigma} \sum_i (C^*, \lambda_i C)^2 - 8g^2 e^{-3\sigma} |W'|^2 \right) 
\]

(18)

From this expression, neglecting the overall factor \( \frac{V}{\kappa^2} 2\pi \rho \), and defining the pseudoscalar by the duality transformation:

\[
\frac{1}{4!} e^{6\sigma} G_{11\mu\nu} = \epsilon_{\mu\nu\rho\sigma} (\partial^\sigma D)
\]

(19)
we obtain the following Kähler potential:

\[ K = - \ln [S + \bar{S}] - 3 \ln [T + \bar{T} - 2C_x^{*}C^x] , \]

(20)

where

\[ S = e^{3\sigma} + i24\sqrt{2}D \]

(21)
and

\[ T = \varphi - i6\sqrt{2}C_5 + C_x^{*}C^x \]

(22)
Here \( D \) and \( C_5 \) are the pseudoscalars and the invisible axions \([7]\). In addition, we have the following gauge kinetic function:

\[ Re f_{\alpha\beta} = f Re S \delta_{\alpha\beta} \]

(23)
and the superpotential \( W \),

\[ W = 8\sqrt{\frac{2}{3}} g_c d_{xyz} C^{x} C^{y} C^{z} \]

(24)
Furthermore, if we perform the field transformation:

\[ \varphi \rightarrow \phi^{3/4} e^{\sigma} ; e^{\sigma} \rightarrow \phi^{-1/4} e^{\sigma} \]

(25)
we obtain \( S, T \) fields which are similar to the previous result in Ref. \([8]\), as noticed in the Ref. \([13]\). This is an interesting result, and allows us to argue that the weakly-coupled heterotic string models derived previously may exist in the M-theory proposed by Horava and Witten. The above 4-dimensional results should be related to the 4-dimensional results from the weakly-coupled heterotic string compactification by a field transformation, although we note that there may not exist a 10-dimensional effective field theory (EFT) \([3]\).
3 Phenomenological consequences

Let us now discuss the physical couplings and the physical radius of the eleventh dimension ($\rho_p$) in the various frames. If we define the Planck mass in $d$ dimensions as $8\pi G_N^{(d)} = M_d^{d-2} = \kappa_d^2$ and $M_4 = M_{Pl} = 2.4 \times 10^{18}$ GeV, from the 11-dimensional Lagrangian [Eq. (1)], we have $M_{11} = \kappa^{-2/9}$, and from the above 4-dimensional Lagrangian [Eq. (15)] in the Einstein frame, we obtain:

$$8\pi G_N^{(4)} E = \frac{\kappa^2}{2\pi \rho V}$$

and

$$\left[\alpha_{GUT}\right]_E = \frac{1}{2V_p f} \left(4\pi \kappa^2\right)^{2/3}$$

where $V_p$ is the physical volume of the Calabi-Yau manifold, i.e., $V_p = V ReS$. Similar relations have been obtained by Witten and others [2, 3, 4, 5, 6] in the metric of the eleven-dimensional theory, these are:

$$8\pi G_N^{(4)} W = \frac{\kappa^2}{2\pi \rho_p V_p'}$$

and

$$\left[\alpha_{GUT}\right]_W = \frac{1}{2V_p f} \left(4\pi \kappa^2\right)^{2/3}$$

where $\rho_p$ is physical radius, i.e., $\rho_p = \frac{1}{2\pi} \int dx^{11} \sqrt{g_{11,11}}$. We have also included the constant $f$, which arises from the following normalization: if we assume that the 4-dimensional grand unified group is $G$ (i.e., $E_8$ is broken to $G$) and if $T$ is a generator of $G$, $\text{Tr}_{E_8}$ and $\text{Tr}_G$ are traces in the adjoint representations of $G$ and $E_8$, then $\text{Tr}_{E_8} T^2 = f \text{Tr}_G T^2$ [9].

Because eleven-dimensional supergravity can be derived from the eleven-dimensional supermembrane world volume action by imposing kappa symmetry [15], and there are arguments in favor of a membrane/fivebrane duality in eleven dimensions [10], it has been argued that the fivebrane units are the natural or fundamental units of M-theory [10]. Therefore, we also consider the above relations in the fivebrane units [10, 14]. We continue to use our above notation, i.e., we do not use the membrane quantization condition [10], since we want to obtain an explicit expression for $\rho_p^{-1}$ and the eleven-dimensional fundamental constant $\kappa$; the result is the same in both approaches. We obtain the relevant 4-dimensional Lagrangian in fivebrane units [10, 14]:

$$L_5 = -\int d^4 x \left(\frac{1}{2\kappa^2}2\pi \rho V e^{\sigma} e^{2\phi/3} R + \frac{f}{2\pi} \left(4\pi \kappa^2\right)^{-2/3} V e^{3\sigma} \text{tr} F_{\mu\nu} F^{\mu\nu}\right)$$

therefore we have:

$$8\pi G_N^{(4)}_{5B} = \frac{\kappa^2}{2\pi \rho_p V_p} e^{2\sigma}$$

and

$$\left[\alpha_{GUT}\right]_{5B} = \frac{1}{2V_p f} \left(4\pi \kappa^2\right)^{2/3}$$

therefore we have:
where \( V_p = e^{3\sigma} V \) and \( \rho_p = \rho e^{2\beta/3} \). We note that in all three cases (Einstein, Witten, and fivebrane), \( \alpha_{\text{GUT}} \) is the same, as the term \( \sqrt{g} \text{tr} F^2 \) is invariant under the rescaling of \( g_{\mu\nu} \) in four dimensions. Therefore, we have in general:

\[
M_{11} = \left[ 2(4\pi)^{-2/3} V_p f \alpha_{\text{GUT}} \right]^{-1/6} \tag{33}
\]

If we define \( V_p = L^{d} l^{6-d} \), ( \( 0 \leq d \leq 6 \) ), where \( L^{-1} \) is the compactification scale and \( l \) is the small internal length, we obtain:

\[
L^{-1} = \left[ 2(4\pi)^{-2/3} f \alpha_{\text{GUT}} \right]^{1/d} \left( \frac{M_{11}}{l^{-1}} \right)^{(6-d)/d} M_{11}, \tag{34}
\]

which tells us that \( L^{-1} \sim M_{11} \), when \( l^{-1} \sim M_{11} \).

Let us now discuss \( \rho_p^{-1} \) in the eleven-dimensional metric and fivebrane units (it is not natural to think of the Lagrangian in the Einstein frame as fundamental). We obtain:

\[
\begin{align*}
\left[ \rho_p^{-1} \right]_W &= 8\pi^2 (2f \alpha_{\text{GUT}})^{-3/2} \left( M_{\text{Pl}}^W \right)^{-2} V_p^{-1/2} \tag{35} \\
\left[ \rho_p^{-1} \right]_{5B} &= 8\pi^2 (2f \alpha_{\text{GUT}})^{-3/2} \left( M_{\text{Pl}}^{5B} \right)^{-2} V_p^{-1/2} e^{-2\sigma} \tag{36}
\end{align*}
\]

The eleven-dimensional length \( (\pi\rho_p)^{-1} \), is of great phenomenological importance because it is related to the scale of supersymmetry breaking \( [3, 7] \). To obtain numerical results we set \( M_{\text{Pl}}^W = M_{\text{Pl}}^{5B} = M_{\text{Pl}} = 2.4 \times 10^{18} \text{ GeV}, \alpha_{\text{GUT}} = \frac{1}{25}, V_p = M_{\text{GUT}}^6, \) and \( M_{\text{GUT}} = 10^{16} \text{ GeV} \). We also set \( f = 1 \) for simplicity and to facilitate comparison with previous papers which only consider this case. We find \([[(\pi\rho_p)^{-1}]_W \sim 1.9 \times 10^{14} \text{ GeV} \) and \([[(\pi\rho_p)^{-1}]_{5B} \sim 1.2 \times 10^{13} \text{ GeV} \) (as \( e^{2\sigma} \approx 16.6 \) if we set \( V^{-1/6} = M_{11} \)).

We now relax the \( f = 1 \) choice and consider the case of realistic grand unified groups. The results for \([[(\pi\rho_p)^{-1}]_W \) and \([[(\pi\rho_p)^{-1}]_{5B} \) for \( G = E_6, \text{SO}(10), \text{SU}(5), \) and \( \text{SU}(5) \times \text{U}(1) \) are listed in Table \( \text{II} \). (We have assumed \( V^{-1/6} = M_{11} \) in all these cases.) We then generally conclude that the supersymmetry-breaking scale is expected to be in \((10^{12} - 10^{14}) \text{ GeV} \) range.

Before addressing further phenomenological features of this scenario, we would like to connect up with our previous phenomenologically oriented study of M-theory–inspired no-scale supergravity in Ref. \([7]\). In that paper we assumed a supergravity model with a no-scale-supergravity–like Kähler potential (implying vanishing universal scalar masses \( m_0 = 0 \)), as suggested by earlier work in Refs. \([3, 4]\). In the present paper we have shown explicitly that such assumption is justified. Moreover, the transmission of supersymmetry-breaking effects from the hidden to the observable sector was assumed to follow the mechanism outlined by Horava \([3]\), whereby such effects are only felt for scales below \((\pi\rho_p)^{-1} \). Our previous sampling of such scales agrees well with our present results in Table \( \text{II} \). These two ingredients were shown \([7]\) to lead to a rather restricted spectrum of superparticle masses within the reach of the present generation of accelerator experiments. It is also worth pointing out another, perhaps more intuitive, explanation of the \( m_0 = 0 \) result: since in M-theory the observable
Table 1: Supersymmetry-breaking scales in the eleven-dimensional metric (W) and fivebrane units (5B) for various choices of (observable) unified gauge groups (G). The parameter f defined in the text is also listed in each case. All scales in GeV.

| G           | f    | (πρ_p)^{-1} W | (πρ_p)^{-1} 5B |
|-------------|------|---------------|----------------|
| E_6         | 2.5  | 4.9 × 10^{13} | 5.3 × 10^{12}  |
| SO(10)      | 3.75 | 2.6 × 10^{13} | 3.8 × 10^{12}  |
| SU(5)       | 6    | 1.3 × 10^{13} | 2.6 × 10^{12}  |
| SU(5)×U(1)  | 6    | 1.3 × 10^{13} | 2.6 × 10^{12}  |

sector fields live in the twisted sector of an orbifold, general arguments [17] indicate that such fields should not feel any supersymmetry-breaking effects.

Because the supersymmetry breaking scale is low, there may exist light gravitinos in the M-theory regime. As has been explored in detailed recently, such light gravitinos may explain the eeγγ + E_T,miss event reported by the CDF Collaboration, and might have some interesting consequences at LEP2 [18]. In this connection, it was noticed some time ago that reactions such as gg → ḡG, ḠG (where g, ḡ, Ḡ stand for gauge boson, gaugino, and gravitino respectively) may exceed the tree-level unitarity limit because of the non-renormalizability of the low-energy effective gravity theory [19]. The critical energy was estimated to be E_{cr} ∼ cM_p m_{3/2}/m_{ḡ}, with c ∼ 10^2. As tree-level unitarity is violated for E > E_{cr}, above E_{cr} the theory is expected to become strongly interacting or change its structure (or both) in order to restore unitarity. The eleventh dimension threshold (πρ_p)^{-1} appears to fulfill such requirements, and thus one might require (πρ_p)^{-1} < E_{cr}, implying a lower bound on m_{3/2} for any given supersymmetry-breaking scale. Taken at face value, this constraint and the results in Table 1 appear to require m_{3/2} > (10^2 − 10^4) eV, which may be consistent with the estimated m_{3/2} < 250 eV required for the χ → γ + Ḡ decay to occur within the detector [18]. However, as we are not sure whether the above unitarity constraints can be used directly in M-theory (as they depend on specific processes that might be forbidden in this scenario), we find the above level of consistency rather encouraging.

There also exist axions in this scenario, for example D: if we define e^{σθ} θ μ ν σ C_{11μν} = ε_{μνσρ} θ σ θ, we have a (θ F^{μν} F_{μν}) term from (G_{11μνρ} G^{11μνρ}). The strong CP problem may be solved by such axions [4]. In this scenario, because the decay constant of these axions is very high ∼ 10^{16} GeV [4, 20], their axino superpartners, if they are very light, might not provide the alternative explanation to the CDF missing energy event proposed in Ref. [21].
4 Conclusions

We have constructed an explicit 4-dimensional $E_6 \times E_8$ model from the M-theory of Horava and Witten. We have also calculated the physical couplings and physical eleventh dimension length, which is presumed to be related to the supersymmetry breaking scale. We have done these calculations in various metrics and shown how they are all related to each other. Finally we discussed the phenomenological consequences of such scenario, which may include light gravitinos, axions, and axinos.

Acknowledgments

T. Li would like to thank K. Benakli, J. X. Lu, and C. N. Pope for useful discussions. This work has been partially supported by the World Laboratory. The work of J. L. has been supported in part by DOE grant DE-FG05-93-ER-40717 and that of D.V.N. by DOE grant DE-FG05-91-ER-40633.
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