M-theory Inspired No-scale Supergravity

Tianjun Li\textsuperscript{1,2}, Jorge L. Lopez\textsuperscript{3}, and D.V. Nanopoulos\textsuperscript{1,2}

\textsuperscript{1}Center for Theoretical Physics, Department of Physics, Texas A&M University
College Station, TX 77843–4242, USA
\textsuperscript{2}Astroparticle Physics Group, Houston Advanced Research Center (HARC)
The Mitchell Campus, The Woodlands, TX 77381, USA
\textsuperscript{3}Bonner Nuclear Lab, Department of Physics, Rice University
6100 Main Street, Houston, TX 77005, USA

Abstract

We propose a supergravity model that contains elements recently shown to arise in the strongly-coupled limit of the $E_8 \times E_8$ heterotic string (M-theory), including a no-scale–like Kähler potential, the identification of the string scale with the gauge coupling unification scale, and the onset of supersymmetry breaking at an intermediate scale determined by the size of the eleventh dimension of M-theory. We also study the phenomenological consequences of such scenario, which include a rather constrained sparticle spectrum within the reach of present-generation particle accelerators.

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Elementary particle phenomenology based on conventional weakly-coupled string theory has met with considerable success over the years. These successes include the taming of quantum gravity, the prediction of space-time supersymmetry (at least at the Planck scale), the prediction of realistic gauge groups and matter representations, and the illumination of the family and fermion mass problems. Despite these accomplishments, string theory has not been able to deliver on its original promise of an unambiguous description of high-energy physics in terms of a single parameter: the Planck mass. This problem is traced to the very large degeneracy of the string vacuum. In addition, different string theories (e.g., heterotic, Type I, Type II) are also allowed. These and other difficulties, like the supersymmetry breaking mechanism, were thought to require an understanding of the strongly-coupled limit of string. Of course, if strings are strongly coupled much doubt is cast on the results obtained at weak coupling that are not protected by non-renormalization theorems. One such result is the desirable prediction of a no-scale supergravity structure [1] with zero vacuum energy at tree level, at least in some string compactifications [2, 3].

The recent discovery of extended duality symmetries relating different string theories at various values of the string coupling [4] has provided a much needed tool to study the strongly-coupled limit of strings. These new developments have elucidated some of the problematic features of weakly-coupled strings [5], such as the relation between different string theories, a new mechanism to reconcile the gauge coupling and string unification scales, and a possible topological mechanism for supersymmetry breaking. The new symmetries require an enlarged set of possible theories, including five-branes and 11-dimensional supergravity.

In this paper we concentrate on the low-energy supersymmetry phenomenology that might arise in the strongly-coupled limit of the $E_8 \times E_8'$ heterotic string in ten dimensions. This theory has been shown [6] to map into a weakly-coupled eleven-dimensional “M-theory”, which has as its low-energy limit eleven-dimensional supergravity. The eleventh dimension ($\rho$) has an orbifold structure ($S^1/Z_2$): at one end live the observable fields contained in $E_8$, at the other end live the hidden sector fields contained in $E_8'$, and in the middle (“bulk”) propagate the gravitational fields. One can distinguish a few mass scales in such scenario: the eleven-dimensional Planck scale ($M_{11}$), the size of the eleventh dimension ($\rho$), and the compactification mass scale ($R^{-1}$). Strong-weak coupling duality relates these three scales as follows [7, 8, 9]:

$$M_{11} = (2\alpha_G V)^{-1/6},$$  \hfill (1)

$$\rho^{-1} = \left(\frac{2}{\alpha_G}\right)^{3/2} M_{Pl}^2 V^{-1/2},$$  \hfill (2)

$$V = R^d \rho^{6-d}; \quad 0 < d \leq 6,$$  \hfill (3)

where $\alpha_G$ is the unified gauge coupling ($\alpha_G \approx \frac{1}{25}$ as deduced from running the Standard Model gauge couplings), $M_{Pl} \approx 10^{19}$ GeV is the usual Planck mass, and $V$ is the volume of the compactified space (which does not include the eleventh dimension). This volume depends on the compactification manifold used, and in Eq. (3) it has
been parametrized in terms of two internal radii: $R$ representing $d$ possibly ‘large’
dimensions and $r$ representing $6 - d$ possible ‘small’ dimensions. In compactification
on symmetric Calabi-Yau manifolds $d = 6$.

The relevance of these equations to gauge coupling unification was pointed
out by Witten [7], who noticed that $M_{11}$, essentially the “string” unification scale
($M_{\text{string}}$), might be lowered down to the gauge coupling unification scale $M_{\text{LEP}}$ by
choosing a sufficiently large compactification volume $V$. Previously the only way
to reduce $M_{\text{string}}$ (at tree level) was to increase the unified gauge coupling, which
then disagreed with the result obtained by running the gauge couplings up from low
energies.

A very interesting result that also emerges from ‘M-phenomenology’ is the ap-
ppearance of the no-scale supergravity structure [8, 11]. No-scale supergravity is then
seen to emerge both in the weakly- and in the strongly-coupled limits of the heterotic
string. In particular, the Kähler potential seems to be of the form that guarantees
universality of the scalar masses (i.e., $m_0 = 0$) at the scale where supersymmetry-
breaking effects are first felt in the observable sector ($\Lambda_{\text{susy}}$). A more model-dependent
question (as it involves the superpotential) is the form of the parameter $A_0$, which for
simplicity we will take to vanish ($A_0 = 0$), as is typical in weakly-coupled string real-
izations of no-scale supergravity [1, 2]. Even more model dependent are the values of
the Higgs mixing parameter $\mu$ and its associated supersymmetry breaking parameter
$B_0$. These will be left unspecified in our present analysis. The gaugino masses will
be taken to be universal ($m_{1/2}$) and will parametrize the low-energy spectrum. These
may be obtained in principle from knowledge of the gauge kinetic function and the
gravitino mass. After using the radiative electroweak symmetry breaking conditions
at the electroweak scale, we need to specify only two parameters to fully describe the
low-energy supersymmetry spectrum: $m_{1/2}$ and the ratio of Higgs vacuum expectation
values $\tan \beta$.

Also of great interest is the mechanism of supersymmetry breaking that ap-
ppears to be emerging in M-phenomenology. Horava [11] has argued that supersym-
metry breaking, say via gaugino condensation in the hidden sector, is not felt im-
mediately in the observable sector because of a topological obstruction (essentially
the eleventh dimension of length $\rho$ that separates the two sectors). Supersymmetry
breaking becomes apparent only after the renormalization scale is low enough to not
reveal the presence of the eleventh dimension anymore (i.e., when this extra dimen-
sion becomes ‘compact’). In practice one is to allow for non-vanishing supersymmetry
breaking parameters only for scales $Q < \rho^{-1}$. This effect can leave a deep imprint
on the low-energy sparticle spectrum, which depends quantitatively on the amount
of ‘running’ of these parameters. As has been noted in connection with the relevance
of gaugino condensation [11, 9] (and as we emphasize below), $\rho^{-1}$ is expected to have
a typical intermediate-scale value. The effect of taking $\Lambda_{\text{susy}} \sim \rho^{-1}$ is most notice-
able in the case of $m_0 = 0$ which we consider here. On purely phenomenological

\footnote{Earlier attempts at reconciling $M_{\text{string}}$ with $M_{\text{LEP}}$ are reviewed in Ref. [11].}
\footnote{An explicit calculation, along the lines of Ref. [9], also supports this conclusion [12].}
grounds, the effect of considering \( \Lambda_{\text{susy}} \) as a new parameter in unified models was (to our knowledge) first discussed in Ref. [13].

Before dwelling into the low-energy phenomenology, let us elucidate the value of \( \Lambda_{\text{susy}} \sim \rho^{-1} \) that we wish to consider. Following Ref. [7] we set \( M_{11} = M_{\text{LEP}} \sim 10^{16} \text{ GeV} \) in Eq. (3). This allows us to solve for \( V = M_{\text{LEP}}^2 / 2\alpha_G \). Substituting this value in the expression for \( \rho^{-1} \) in Eq. (2) we obtain

\[
\rho^{-1} = \frac{4}{\alpha_G} \left( \frac{M_{\text{LEP}}}{M_{\text{Pl}}} \right)^3 M_{\text{Pl}} \sim \left( 10^{12} - 10^{13} \right) \text{ GeV} ,
\]

which is very suggestive in the context of gaugino condensation where \( m_{\text{susy}} \sim \rho^{-3} / M_{\text{Pl}}^2 \sim M_Z \). It also follows that

\[
R^{-1} = \left( 2\alpha_G \right)^{1/d} \left( \frac{M_{\text{LEP}}}{r^{-1}} \right)^{(6-d)/d} M_{\text{LEP}}
\]

Note that in the reasonable case of \( r^{-1} \sim M_{11} \sim M_{\text{LEP}} \), the \( d \) dependence of \( R^{-1} \) is greatly reduced and \( R^{-1} \sim M_{\text{LEP}} \) is also obtained.

We now proceed to the analysis of the low-energy sparticle spectrum under the assumptions of \( \Lambda_{\text{susy}} = \rho^{-1} \) and \( m_0 = A_0 = 0 \). As remarked above, the only free parameters are \( m_{1/2} \) and \( \tan \beta \). We find that the requirement of radiative electroweak symmetry breaking plus two basic phenomenological requirements, allow solutions in the \( (m_{1/2}, \tan \beta) \) plane for only one sign of \( \mu \) and only within a completely bounded region. For the case of \( \Lambda_{\text{susy}} = 10^{13} \text{ GeV} \), this region is shown in Fig. 1, where to facilitate comparison with experiment we also show the region in the \( (m_{\chi^\pm}, \tan \beta) \) plane. The upper limit on \( m_{1/2} \) (for a fixed value of \( \tan \beta \)) follows from the requirement that the lightest supersymmetric particle be neutral [14]. Above the upper boundary the right-handed selectron (\( \tilde{e}_R \)) becomes lighter than the lightest neutralino (\( \chi \)). The bottom boundary is obtained by imposing the absolute lower limit on the sneutrino mass from LEP 1 searches \( (m_{\tilde{\nu}} > 43 \text{ GeV}) \). The area to the right of the right-most tip of the region is excluded by these two conflicting constraints. The \( \tan \beta \) dependence of these constraints may be understood from the D-term contribution to the \( \tilde{e}_R \) and \( \tilde{\nu} \) mass formulas

\[
\tilde{m}_i^2 = c_i m_{1/2}^2 - d_i \tan^2 \beta - 1 \tan^2 \beta + 1 M_W^2 ,
\]

where the \( c_i \) are some RGE-dependent constants and \( d_{\tilde{e}_R} = -\tan^2 \theta_W < 0 \) whereas \( d_\nu = \frac{1}{2} (1 + \tan^2 \theta_W) > 0 \). The dotted line indicates the lower bound on \( \tan \beta \) that is consistent with the top-quark mass \( (m_t = 175 \text{ GeV}) \) and perturbative Yukawa couplings up to the unification scale [14]. In practice, the LEP 172 lower bound on the chargino mass \( (m_{\chi^\pm} > 83 \text{ GeV}) \) gives the strongest constraint on the

\footnote{The procedure is standard, see e.g., Refs. [13, 14].}

\footnote{In practice, because we allow for a non-zero tau lepton mass, the lightest stau mass eigenstate (\( \tilde{\tau}_1 \)) is in effect slightly lighter than the right-handed selectron.}

\footnote{This limit applies as long as \( m_{\chi^\pm} - (m_{\chi} \text{ or } m_{\nu}) > 3 \text{ GeV} \), which is always satisfied in this model.
parameter space (dashed line on bottom panel in Fig. 1). Nonetheless, a portion of the parameter space remains allowed, and in fact it is within the reach of future LEP 2 energy upgrades, as we discuss below.

To give a more detailed picture of the low-energy spectrum, in Fig. 2 we display representative sparticle masses as a function of the chargino mass for $\Lambda_{\text{susy}} = 10^{13}$ GeV and $\tan \beta = 3$. This choice of $\tan \beta$ allows the widest range of sparticle masses (see Fig. 1). This figure shows that the spectrum ‘terminates’ when $m_\chi$ approaches $m_{\tilde{e}_R}$ from below, as mentioned above in connection with the upper boundary in Fig. 1. It is interesting to note the significant splitting of the top-squark ($\tilde{t}_1, \tilde{t}_2$) masses around the average squark ($\tilde{q}$) mass.

In the LEP 172 allowed region in Figs. 1 and 2 we find $m_{\chi^\pm} < 95$ GeV and $m_{\tilde{e}_R} < 70$ GeV. Both of these particles appear within the reach of LEP 2. More to the point, one might wonder whether the such light right-handed selectron masses might have already excluded by LEP 2 searches, as they have been certainly kinematically accessible. We have calculated the cross section $\sigma(e^+e^- \rightarrow \tilde{e}_R^+\tilde{e}_R^-)$ at LEP 161, for which explicit limits have been released by the OPAL Collaboration [16]. We find $\sigma < 0.2 \text{ pb}$, which in $L = 10.1 \text{ fb}^{-1}$ would have yielded a maximum of two events. Indeed, the experimental sensitivity to this mode is at the 0.5 pb level [16]. Thus, past LEP 2 searches in the selectron channels do not restrict the allowed parameter space any further.

One should also consider the predictions for trilepton events at the Tevatron. We find $\sigma(p\bar{p} \rightarrow \chi^\pm\chi') \approx (1.0 - 0.7) \text{ pb}$ for $m_{\chi^\pm} = (83 - 95)$ GeV. The leptonic decays of the chargino and neutralino are maximally enhanced because of the lighter right-handed sleptons and sneutrinos, respectively. That is, $B(\chi^\pm \rightarrow \ell \nu \chi) \approx 2/3$ and $B(\chi' \rightarrow \ell^+\ell^-\chi) \approx 1/2$, where $\ell = e + \mu$. Combining these numbers we arrive at a single channel (i.e., any single one of $eee, eep, e\mu\mu$, or $\mu\mu\mu$) cross section of $(0.16 - 0.11) \text{ pb}$. This result is slightly below the sensitivity reached at the Tevatron in trilepton searches [17], and thus these also do not constrain the allowed parameter space any further.

One may also want to study the effect of different choices for $\Lambda_{\text{susy}}$. These affect the shape of the bounded region in Fig. 1, although not very significantly in the interval specified by Eq. (4). As a means of quantifying this behavior, one can determine the maximum allowed (i.e., passing LEP 1 cuts) chargino masses for a range of $\Lambda_{\text{susy}}$ values. These are shown in Fig. 3. The figure shows that LEP 2 data presently do not constrain $\Lambda_{\text{susy}}$ in any way, although future higher energy runs will soon begin to constrain $\Lambda_{\text{susy}}$.

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Moreover, near the upper end of the parameter space the experimental detection efficiency should be greatly reduced because $m_{\tilde{e}_R}$ approaches $m_{\chi}$.
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References

[1] E. Cremmer, S. Ferrara, C. Kounnas, and D. V. Nanopoulos, Phys. Lett. B 133 (1983) 61; J. Ellis, A. Lahanas, D. V. Nanopoulos, and K. Tamvakis, Phys. Lett. B 134 (1984) 429; J. Ellis, C. Kounnas, and D. V. Nanopoulos, Nucl. Phys. B 241 (1984) 406 and B 247 (1984) 373. For a review see A. Lahanas and D. V. Nanopoulos, Phys. Rep. 145 (1987) 1.

[2] E. Witten, Phys. Lett. B 155 (1985) 151.

[3] J. L. Lopez, D. V. Nanopoulos, and K. Yuan, Phys. Rev. D 50 (1994) 4060; J. L. Lopez and D. V. Nanopoulos, Int. J. Mod. Phys. A 11 (1996) 3439.

[4] For a recent review see e.g., J. Polchinski, Rev. Mod. Phys. 68 (1996) 1245.

[5] For a recent review see e.g., S. Ferrara, hep-th/9610083.

[6] P. Horava and E. Witten, Nucl. Phys. B 475 (1996) 94.

[7] E. Witten, Nucl. Phys. B 471 (1996) 135.

[8] T. Banks and M. Dine, Nucl. Phys. B 479 (1996) 173.

[9] E. Caceres, V. Kaplunovsky, and I. M. Mandelberg, hep-th/9606036; I. Antoniadis and M. Quiros, hep-th/9609209.

[10] K. Dienes, hep-th/9602043.

[11] P. Horava, Phys. Rev. D 54 (1996) 7561.

[12] T. Li, J. L. Lopez, and D. V. Nanopoulos, to appear.

[13] S. Kelley, J. L. Lopez, and D. V. Nanopoulos, Mod. Phys. Lett. A 8 (1993) 177.

[14] S. Kelley, et. al., Nucl. Phys. B 398 (1993) 3.

[15] LEP Experiments Committee Open Session, CERN, November 19, 1996. Presentations by F. Ragusa (ALEPH Collaboration) http://alephwww.cern.ch/ALPUB/seminar/lepc_nov96, J. Timmermans (DELPHI Collaboration), D. Stickland (L3 Collaboration), A. Ball (OPAL Collaboration).

[16] K. Ackerstaff, et. al. (OPAL Collaboration), CERN-PPE/96-182.

[17] K. W. Merritt, hep-ex/9701009.
Figure 1: The allowed region in \((m_{1/2}, \tan \beta)\) [top panel] and correspondingly \((m_{\chi^\pm}, \tan \beta)\) [bottom panel] in no-scale supergravity \((m_0 = A_0 = 0)\) with \(\Lambda_{\text{susy}} = 10^{13}\) GeV. Above the top boundary \(m_{\tilde{e}_R} \approx m_{\tilde{\tau}_1} < m_{\chi}\), whereas below the bottom boundary \(m_{\tilde{\nu}_L} < 43\) GeV. The dashed line [bottom panel] represents the lower bound on the chargino mass from LEP 172 searches.
Figure 2: Calculated values of representative sparticle masses versus the chargino mass for $\Lambda_{\text{susy}} = 10^{13}$ GeV and $\tan \beta = 3$. The spectrum terminates when $m_\chi$ approaches $m_{\tilde{e}_R}$ from below. The dashed line represents the lower bound on the chargino mass from LEP 172 searches.
Figure 3: The maximum allowed value of the chargino mass as a function of $\Lambda_{\text{susy}}$ in no-scale supergravity ($m_0 = A_0 = 0$). The dashed line represents the lower bound on the chargino mass from LEP 172 searches.