SPARTICLE SPECTRUM AND DARK MATTER IN M-THEORY

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ABSTRACT

The phenomenological implications of the eleven dimensional limit of M-theory (strongly coupled $E_8 \times E_8$) are investigated. In particular we calculate the supersymmetric spectrum subject to constraints of correct electroweak symmetry breaking and the requirement that the lightest supersymmetric particle provides the dark matter of the universe. The $B$-soft term associated with the generation of a $\mu$ term in the superpotential is calculated and its phenomenology is discussed.
In recent years it has become clear that the five perturbative string theories and the 11 dimensional supergravity are different limits in moduli space of a unique fundamental theory. This strongly indicates an enormous degree of symmetry of the underlying theory and its intrinsically non-perturbative nature. String duality correlates the six corners of the moduli space. The duality transformation involves Planck’s constant $\hbar$ and is therefore intrinsically quantum mechanical. One then might argue that before we have the complete picture of $M$(other)theory it is premature to make any attempt at phenomenology. However, it may be that the corners of the moduli space capture most of the features of the theory relevant for low-energy phenomenology.\(^3\)

One of the most interesting dualities (and the most relevant for low-energy phenomenology) is the one in which the low-energy limit of $M$-theory (i.e 11D-Supergravity) compactified on the line segment $I \sim S^1/Z_2$ (i.e an orbifold) is equivalent to the strong coupling limit of the $E_8 \times E_8$ heterotic string [1]. In this picture one end of the line segment of length $\pi \rho$ live the observable fields contained in the first $E_8$ while the hidden sector fields live in the second $E_8$ factor on the other end. Gravitational fields propagate in the bulk.

The main phenomenological virtue of such a framework is that due to the extra dimension one may obtain unification of all interactions at a scale $M_U \sim 3 \times 10^{16}$ GeV consistent with experimental data for the low energy gauge couplings [4]. Also the analysis of gaugino condensation reveals that phenomenologically acceptable gaugino masses, comparable with the gravitino mass $m_3/2$, arise quite naturally in sharp contrast to the weakly coupled case where tiny gaugino masses were troublesome [7]. It is therefore of great importance to investigate further the phenomenological implications of the 11-dimensional low energy limit of $M$-theory and to determine any deviations from the weakly coupled case.

Several papers have recently analyzed the effective supergravity and the soft supersymmetry-breaking terms emerging in this framework [4, 5, 7, 8, 9, 10, 11]. Some properties of the sparticle spectrum which depend only on the boundary conditions and not on the details of the electroweak symmetry breaking, have been discussed in [11]. However, a detailed analysis of the spectrum subject to the constraint of correct electroweak symmetry breaking [12], as well as other phenomenological implications, has not been performed. It is the purpose of this paper to investigate the phenomenological implications of $M$-theory relevant to accelerator experiments and the cosmological properties of the lightest supersymmetric particle (LSP) as well as the prospects for its detection in underground non-baryonic dark matter experiments.

The soft supersymmetry-breaking terms are determined by the following functions of the effective supergravity theory [11, 10]:

\[
K = -\ln(S + \bar{S}) - 3\ln(T + \bar{T}) + \left(\frac{3}{T + \bar{T}} + \frac{\alpha}{S + \bar{S}}\right)|C|^2,
\]

\(^3\)The 11-D limit of $M$-theory is not gauge invariant so Horava and Witten have argued that quantum terms are needed to restore gauge invariance [1]. See however, M. Faux’s argument for a consistent classical limit of $M$-theory [2].
\[ f_{E_6} = S + \alpha T, \quad f_{E_8} = S - \alpha T, \]
\[ W = d_{pq} C^p C^q C^r \]

where \( K \) is the Kähler potential, \( W \) the perturbative superpotential, and \( f_{E_6}, f_{E_8} \) are the gauge kinetic functions for the observable and hidden sector gauge groups \( E_6 \) and \( E_8 \) respectively. Also \( S, T \) are the dilaton and Calabi-Yau moduli fields and \( C^\alpha \) charged matter fields. The superpotential and the gauge kinetic functions are exact up to non-perturbative effects. The integer \( \alpha = \frac{1}{8\pi^2} \int \omega \wedge [\text{tr}(F \wedge F) - \frac{1}{2}\text{tr}(A \wedge A)] \) for the Kähler form \( \omega \) is normalized as the generator of the integer (1,1) cohomology.

Given eqs(1) one can determine \([10, 11]\) the soft supersymmetry breaking terms for the observable sector gaugino masses \( M_{1/2} \), scalar masses \( m_0 \) and trilinear scalar masses \( A \) as functions of the auxiliary fields \( F_S \) and \( F_T \) of the moduli \( S, T \) fields respectively.

\[ M_{1/2} = \frac{\sqrt{3} C m_{3/2}}{(S + S + \alpha(T + T))} \left( (S + \bar{S}) \sin \theta + \frac{\alpha(T + \bar{T}) \cos \theta}{\sqrt{3}} \right), \]
\[ m_0^2 = V_0 + m_{3/2}^2 - \frac{3m_{3/2}^2 C^2}{3(S + S) + \alpha(T + T)} \]
\[ \times \left\{ \alpha(T + \bar{T}) \left( 2 - \frac{\alpha(T + \bar{T})}{3(S + S) + \alpha(T + T)} \right) \sin^2 \theta \right. \]
\[ + \left. (S + S) \left( 2 - \frac{3(S + \bar{S})}{3(S + S) + \alpha(T + T)} \right) \cos^2 \theta \right\} \]
\[ - \frac{2\sqrt{3} \alpha(T + \bar{T})(S + \bar{S})}{3(S + S) + \alpha(T + T)} \sin \theta \cos \theta, \]
\[ A = \sqrt{3} C m_{3/2} \left\{ \left( -1 + \frac{3\alpha(T + \bar{T})}{3(S + S) + \alpha(T + T)} \right) \sin \theta \right. \]
\[ + \left. \sqrt{3} \left( -1 + \frac{3(S + S)}{3(S + S) + \alpha(T + T)} \right) \cos \theta \right\} \]

while the \( B \)-soft term associated with non-perturbatively generated \( \mu \) term in the superpotential is given by:

\[ B_{\mu} = m_{3/2} \left[ -3C \cos \theta - \sqrt{3} C \sin \theta \right. \]
\[ + \left. \frac{6C \cos \theta (S + \bar{S})}{3(S + S) + \alpha(T + T)} + \frac{2\sqrt{3} \sin \theta \alpha(T + \bar{T})}{3(S + S) + \alpha(T + T)} - 1 + F_S \frac{\partial \ln \mu}{\partial S} + F_T \frac{\partial \ln \mu}{\partial T} \right] \]

\[ \text{We assume very small CP-violating phases in the soft terms. This assumption is supported by the CP-structure of the soft terms in T-duality invariant string compactifications at least for the T-moduli sector [14].} \]
In (2),(3) the auxiliary fields are parametrized as follows [13]:

\[ F^S = \sqrt{3} m_{3/2} C (S + \bar{S}) \sin \theta, \]
\[ F^T = m_{3/2} C (T + \bar{T}) \cos \theta \]  

(4)

and \( \theta \) is the goldstino angle which specifies the extent to which the supersymmetry breaking resides in the dilaton versus the moduli sector. Also \( m_{3/2} \) is the gravitino mass and \( C^2 = 1 + \frac{V_0}{3m_{3/2}} \) with \( V_0 \) the tree level vacuum energy density. Note that in the limit \( \alpha(T + \bar{T}) \to 0 \) we recover the soft terms of the weakly coupled large \( T \)-limit of Calabi-Yau compactifications [13].

We now consider the supersymmetric spectrum. Our parameters are the goldstino angle \( \theta \), \( \alpha(T + \bar{T}), \text{sign} \mu \) (which is not determined by the radiative electroweak symmetry breaking constraint), where \( \mu \) is the Higgs mixing parameter in the low energy superpotential, and \( \tan \beta \) (i.e the ratio of the two Higgs vacuum expectation values \( \tan \beta = \frac{<H_2>}{<H_1>} \)) if we leave \( B \) a free parameter determined by the minimization of the Higgs potential. If \( B \) instead is given by [3], one determines the value of \( \tan \beta \). For this purpose we take \( \mu \) independent of \( T \) and \( S \) because of our lack of knowledge of \( \mu \) in \( M \)-theory. We also set \( C = 1 \) in the above expressions assuming zero cosmological constant. The soft masses start running from a mass \( R_{11}^2 \sim 7.5 \times 10^{15} \text{GeV} \) with \( R_{11} \) the extra \( M \)-theory dimension. This is perhaps the most natural choice. However values as low as \( 10^{13} \) \text{GeV} are possible and have been advocated by some authors [5]. However, the recent analysis of [7] disfavours such scenarios. For the most part of our analysis we shall consider the former value of \( R_{11} \), but we shall also comment on the consequences of the latter.

Then using (2) (3) as boundary conditions for the soft terms, one evolves the renormalization group equations down to the weak scale and determines the sparticle spectrum compatible with the constraints of correct electroweak symmetry breaking and experimental constraints on the sparticle spectrum from unsuccessful searches at LEP, Tevatron etc, and also that the LSP provides a good dark matter candidate.

Electroweak symmetry breaking is characterized by the extrema equations

\[ \frac{1}{2} M_Z^2 = \frac{\bar{m}_{H_1}^2 - \bar{m}_{H_2}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \mu^2 \]
\[- B \mu = \frac{1}{2} (\bar{m}_{H_1}^2 + \bar{m}_{H_2}^2 + 2 \mu^2) \sin 2\beta \]  

(5)

where

\[ \bar{m}_{H_1,H_2}^2 \equiv m_{H_1,H_2}^2 + \frac{\partial \Delta V}{\partial v_{1,2}^2} \]  

(6)

and \( \Delta V = (64\pi^2)^{-1} \text{Str} M^4 [\ln(M^2/Q^2) - \frac{3}{2}] \) is the one loop contribution to the Higgs effective potential. We include contributions from the third generation of particles and sparticles.

Since \( \mu^2 \gg M_Z^2 \) for most of the allowed region of the parameter space [13], the following approximate relationships hold at the electroweak scale for the masses
of neutralinos and charginos, which of course depend on the details of electroweak symmetry breaking.

\[ m_{\chi_{1}^{\pm}} \sim m_{\chi_{0}^{2}} \sim 2m_{\chi_{1}^{0}} \]
\[ m_{\chi_{3,4}^{\pm}} \sim |\mu| \]  

(7)

In (7) \( m_{\chi_{1,2}^{\pm}} \) are the chargino mass eigenstates and \( m_{\chi_{0}^{i}}, i = 1 \ldots 4 \) are the four neutralino mass eigenstates with \( i = 1 \) denoting the lightest neutralino. The former arise after diagonalization of the mass matrix.

\[
M_{ch} = \begin{pmatrix}
M_2 & \sqrt{2}m_{W} \sin \beta \\
m_{W} \cos \beta & -\mu
\end{pmatrix}
\]  

(8)

We consider first the extreme \( M \)-theory limit \( \theta = \frac{\pi}{20} \). An interesting case is when the scalar masses are much smaller than the gaugino masses at the unification scale. \( \theta \) For instance this can be achieved by a goldstino angle \( \theta = \frac{\pi}{20} \). In this case sleptons are much lighter compare to the gluino than in the weakly coupled Calabi Yau compactifications of the heterotic string for any value of \( \theta \). As a result for high values of \( \tan \beta \) \( \theta \) there is a possibility that right handed selectrons or the lightest stau mass eigenstate become the LSP. The stau mass matrix is given by the expression

\[
\mathcal{M}_{\tilde{\tau}}^{2} = \begin{pmatrix}
\mathcal{M}_{11}^{2} & m_{\tilde{\tau}}(A_{\tau} + \mu \tan \beta) \\
m_{\tilde{\tau}}(A_{\tau} + \mu \tan \beta) & \mathcal{M}_{22}^{2}
\end{pmatrix}
\]  

(9)

This of course is phenomenologically unacceptable and results in strong constraints in the sparticle spectrum. In fig.1 we plot the critical value of gravitino mass \( m_{3/2}^{\tilde{g}} \) above (below) which \( m_{\tilde{\tau}_{L}} \) or \( m_{\tilde{e}_{R}} < m_{\chi_{0}^{1}} \) \( (m_{\tilde{\tau}_{L}} < 43 \text{GeV}) \) versus \( \tan \beta \). The acceptable parameter space lies between the upper and lower bound in fig.1. One can see that because of the above contrainst \( \tan \beta \leq 13 \) in this \( M \)-theory limit because the acceptable parameter space vanishes for larger values of \( \tan \beta \). We also plot in fig.2 the critical chargino mass \( m_{\chi_{1}^{\pm}} \) above which the right handed sleptons become the LSP, versus \( \tan \beta \). In the same figure we also draw the experimental lower bound for the chargino mass from LEP of about 83 GeV (horizontal line). The \( \tan \beta \) dependence of these constraints may be understood from the \( D \)-term contribution to the \( \tilde{\tau}_{R} \) and \( \tilde{\nu} \) mass formulas

\[
\tilde{m}_{\tilde{\tau}}^{2} = c_{i} m_{3/2}^{2} - d_{i} \frac{\tan^{2} \beta - 1}{\tan^{2} \beta + 1} M_{W}^{2}
\]  

(10)

\(^{5}\) Note that in this limit the dilaton dominated scenario , i.e \( (\theta = \frac{\pi}{20}) \) is not feasible.

\(^{6}\) Note that the weakly coupled case scalar masses are comparable to gaugino mass \( m_{0} \sim M \) at the unification scale.

\(^{7}\) In case we determine \( B \) from the electroweak symmetry breaking \( \tan \beta \) is a free parameter; also in this case \( M_{\tilde{g}} : m_{L_{L}} : m_{\tilde{e}_{R}} \sim 1 : 0.25 : 0.14 \)
where the \( c_i \) are some RGE-dependent constants and \( d_{\tilde{e}_R} = -\tan^2 \theta_W < 0 \) whereas \( d_\nu = \frac{1}{2}(1 + \tan^2 \theta_W) > 0 \).

In the allowed region of fig.1 the LSP is the lightest neutralino \( \chi_1^0 \). Assuming \( R \)-parity conservation the LSP is stable and consequently can provide a good dark matter candidate. It is a linear combination of the superpartners of the photon, \( Z^0 \) and neutral-Higgs bosons,

\[
\chi_1^0 = c_1 \tilde{B} + c_2 \tilde{W}^3 + c_3 \tilde{H}_1^0 + c_4 \tilde{H}_2^0
\]  

(11)

The neutralino 4 \times 4 mass matrix can be written as

\[
\begin{pmatrix}
M_1 & 0 & -M_Z A_{11} & M_Z A_{21} \\
0 & M_2 & M_Z A_{12} & -M_Z A_{22} \\
-M_Z A_{11} & M_Z A_{12} & 0 & \mu \\
M_Z A_{21} & -M_Z A_{22} & \mu & 0
\end{pmatrix}
\]

with

\[
\begin{pmatrix}
A_{11} & A_{12} \\
A_{21} & A_{22}
\end{pmatrix} = \begin{pmatrix}
\sin \theta_W \cos \beta & \cos \theta_W \cos \beta \\
\sin \theta_W \sin \beta & \cos \theta_W \sin \beta
\end{pmatrix}
\]

In fact the lightest neutralino in this model is almost a pure bino (\( \tilde{B} \)), which means \( f_g \equiv |c_1|^2 + |c_2|^2 \gg 0.5 \). Most cosmological models predict that the relic abundance of neutralinos [16] satisfies

\[
0.1 \leq \Omega_{\text{LSP}} h^2 \leq 0.4
\]

(12)

We calculated the relic abundance of the lightest neutralino using standard technology [17] and found strong constraints on the resulting spectrum. In particular for \( \mu < 0 \) the lower limit on the relic abundance results in a lower limit on the gravitino mass of about 200 GeV. In figs(3-6), we plot the relic abundance of the lightest neutralino versus the gravitino mass for various values of \( \tan \beta \) for goldstino angle of \( \theta = \frac{7\pi}{20} \). We also note that for the allowed parameter space \( \Omega_{\text{LSP}} h^2 \) never exceeds the upper limit of 0.4. However, for other values of the goldstino angle the upper limit on \( \Omega_{\text{LSP}} h^2 \) can constrain the gravitino mass. The lower limit also constrains the allowed values of the \( \tan \beta \) parameter even further. For instance one can see from the plots that \( \tan \beta < 10 \) in order that \( \Omega_{\text{LSP}} h^2 \geq 0.1 \). In the allowed physical region direct detection rates are of order \( 10^{-2} - 10^{-4} \text{events/Kg/day} \). The lightest Higgs \( m_h \) is in the range \( 87 GeV \leq m_h \leq 115 GeV \), while the neutralino mass is in the range \( 77 GeV \leq m_{\chi_1^0} \leq 195 GeV \). At this point is worth mentioning that if one chooses to run the soft masses from the mass \( R_{11}^{-1} \sim O(10^{13}) \text{GeV} \) instead of \( 7.5 \times 10^{15} \text{GeV} \) the cosmological constraint (12) is very powerful and eliminates all of the parameter space since the relic abundance is always much smaller than 0.1, when \( \theta = \frac{7\pi}{20} \). For \( \mu > 0 \) the maximum gravitino mass above which \( m_{\tilde{e}_R} < m_{\chi_1^0} \) or
$m_{\tilde{g}} < m_{\chi^0_1}$ is smaller for fixed $\tan \beta$. Clearly this novel $M$-theory limit provide us with a phenomenology distinct from the weakly coupled case which should be a subject of experimental scrutiny.

For other values of the goldstino angle for which scalar masses are comparable to the gaugino masses (a case which is more similar to the weakly coupled limit \cite{13} in which $m_0 \sim \frac{1}{\sqrt{3}} M_{1/2}$) we do not obtained constraints from the bounds on the mass of right handed selectrons and staus, but in this case the upper limit on the relic abundance leads to an upper limit on the gravitino mass. For instance, for a goldstino angle of $\theta = \frac{\pi}{4}$ (see fig.7) and $\tan \beta = 2.5$, the requirement that $\Omega_{LSP} h^2 \leq 0.4$ results in $m_{3/2} \leq 365$ GeV. This results in an upper limit of the lightest Higgs $m_h \leq 100$ GeV. The lower limit is now $m_{3/2} \geq 115$ GeV. However, the LEP limit on the chargino mass of 82 GeV requires that $m_{3/2} \geq 150$ GeV. In this case the tightest neutralino is in the range $52$ GeV $\leq m_{\chi^0_1} \leq 148$ GeV for $\mu < 0$. For $\mu > 0$ we have $65$ GeV $\leq m_{\chi^0_1} \leq 153$ GeV. Detection rates of the LSP for $\Omega_{LSP} h^2 < 0.4$ results in $m_{3/2} \leq 365$ GeV. Direct detection rates of the lightest neutralino are in the range $6 \times 10^{-2} - 10^{-4}$ events/Kg/day for $\mu < 0$ and $O(10^{-3}) - O(10^{-5})$ events/Kg/day for $\mu > 0$. For higher $\tan \beta$ values on can obtain higher detection rates.

In conclusion, we have analyzed the supersymmetric spectrum and the properties of the lightest neutralino (LSP) in the 11-dimensional limit of M-theory. The most striking result in the case of small scalar masses compare with gluino masses, is that one obtains a limit on $\tan \beta \leq 13$, since above that value the right handed selectron or the lightest stau is the LSP, which is phenomenologically unacceptable since the LSP should be electrically neutral. Also the cosmological constraint on the relic abundance of the LSP results in a lower limit on the gravitino mass $m_{3/2} \geq 200$ GeV. This further constrains $\tan \beta$; $\tan \beta < 10$. In this case the upper limit on the relic abundance is not relevant since $\Omega_{LSP} h^2 < 0.4$ for goldstino angle $\theta = \frac{\pi}{4}$ and for all the allowed values of $\tan \beta$. Also the lower limit on the relic abundance excludes the case of $R_{11}^{-1} \sim O(10^{13})$ GeV. The scenario with $B_\mu$ given by \cite{3}, and $\mu$ independent of, $S$ and, $T$ is excluded since one has to go to non-perturbative Yukawa couplings in order to obtain a value of $\tan \beta$ consistent with $B_\mu$ as in \cite{3}. For other values of the goldstino angle (which resemblance more the weakly coupled Calabi-Yau compactifications for which $m_0 = \frac{1}{\sqrt{3}} M_{1/2}$) the upper bound on the relic abundance results in an upper bound on the gravitino mass. For a goldstino angle, $\theta = \frac{\pi}{4}$ and $\tan \beta = 2.5$ then we find $m_{3/2} \leq 365$ GeV. Direct detection rates of the lightest neutralino are in the range of $10^{-1} - 10^{-4}$ events/Kg/day. The $M$-theory limit yields interesting phenomenology which should be the target of experimental investigation.

\footnote{Note however, that this M-theory limit is somewhat similar to the O-I orbifold model \cite{13, 18} and for a particular goldstino angle, different from the dilaton-dominated limit. On the other hand the O-I model has non-universal soft supersymmetry-breaking terms at the string scale.}
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Figure 1: Allowed parameter space in $m_{3/2}, \tan \beta$ plane

$\theta = \frac{7\pi}{20}, \alpha(T+\bar{T}) = 2, \mu < 0$
Figure 2: Maximum lightest chargino mass vs tan $\beta$

\[ \theta = \frac{7\pi}{20}, \alpha(T + \overline{T}) = 2, \mu < 0 \]
Figure 3: Relic abundance versus $m_{3/2}$ for $\tan \beta = 2.5, \mu < 0$
Figure 4: Relic abundance of LSP vs $m_{3/2}$ for $\tan \beta = 2.5, \mu > 0$
Figure 5: Relic abundance of $LSP_{13}$ vs $m_{3/2}$ for $\tan \beta = 5, \mu < 0$
Figure 6: Relic abundance of LSP in case of $\tan \beta = 10$
Figure 7: Relic abundance of LSP in case of $\theta = \frac{\pi}{4}$ and $\tan \beta = 2.5$. 

$\Omega_{LSP} h^2 = \frac{\pi}{4}, a(T+T) = 2, \tan \beta = 2.5, \mu < 0$