Predictions of the Constrained Exceptional Supersymmetric Standard Model

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Abstract

We discuss the predictions of a constrained version of the exceptional supersymmetric standard model (cE\textsubscript{6}SSM), based on a universal high energy soft scalar mass \(m_0\), soft trilinear coupling \(A_0\) and soft gaugino mass \(M_{1/2}\). We predict a supersymmetry (SUSY) spectrum containing a light gluino, a light wino-like neutralino and chargino pair and a light bino-like neutralino, with other sparticle masses except the lighter stop being much heavier. In addition, the cE\textsubscript{6}SSM allows the possibility of light exotic colour triplet charge 1/3 fermions and scalars, leading to early exotic physics signals at the LHC. We focus on the possibility of a \(Z'\) gauge boson with mass close to 1 TeV, and low values of \((m_0, M_{1/2})\), which would correspond to an LHC discovery using “first data”, and propose a set of benchmark points to illustrate this.

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1. Introduction

The minimal supersymmetric standard model (MSSM) \[1\] provides a very attractive supersymmetric extension of the standard model (SM). Its superpotential contains the bilinear term $\mu H_d H_u$, where $H_d, u$ are the two Higgs doublets which develop vacuum expectation values (VEVs) at the weak scale and $\mu$ is the supersymmetric Higgs mass parameter which can be present before SUSY is broken. However, despite its attractiveness, the MSSM suffers from the $\mu$ problem: one would naturally expect $\mu$ to be either zero or of the order of the Planck scale, while, in order to get the correct pattern of electroweak symmetry breaking (EWSB), $\mu$ is required to be in the TeV range.

It is well known that the $\mu$ term of the MSSM can be generated effectively by the low energy VEV of a singlet field $S$ via the interaction $\lambda S H_d H_u$. However, although an extra singlet superfield seems like a minor modification to the MSSM, which does no harm to either gauge coupling unification or neutralino dark matter, its introduction leads to an additional accidental global $U(1)$ (Peccei-Quinn (PQ) \[2\]) symmetry which will result in a weak scale massless axion when it is spontaneously broken by $\langle S \rangle$ \[3\]. Since such an axion has not been observed experimentally, it must be removed somehow. This can be done in several ways resulting in different non-minimal SUSY models, each involving additional fields and/or parameters \[4, 5\]. For example, the classic solution to this problem is to introduce a singlet term $S^3$, as in the next-to-minimal supersymmetric standard model (NMSSM) \[4\], which reduces the PQ symmetry to the discrete symmetry $Z_3$. However the subsequent breaking of a discrete symmetry at the weak scale can lead to cosmological domain walls which would overclose the Universe.

A cosmologically safe solution to the axion problem of singlet models, which we follow in this Letter, is to promote the PQ symmetry to an Abelian $U(1)'$ gauge symmetry \[6\]. The idea is that the extra gauge boson will eat the troublesome axion via the Higgs mechanism resulting in a massive $Z'$ at the TeV scale. The necessary $U(1)'$ gauge group could be a relic of the breaking of some unified gauge group at high energies. Recall that the unification of gauge couplings in SUSY models allows one to embed the gauge group of the SM into Grand Unified Theories (GUTs) based on simple gauge groups such as $SU(5)$, $SO(10)$ or $E_6$. In particular the $E_6$ symmetry can be broken to the rank–5 subgroup $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)'$ where in general $U(1)' = U(1)_\chi \cos \theta + U(1)_\psi \sin \theta$ \[7\], and the two anomaly-free $U(1)_\psi$ and $U(1)_\chi$ symmetries originate from the breakings $E_6 \rightarrow SO(10) \times U(1)_\psi$, $SO(10) \rightarrow SU(5) \times U(1)_\chi$ (for recent review see \[8\]).

Within the class of $E_6$ models there is a unique choice of Abelian gauge group that allows zero charges for right-handed neutrinos and thus large Majorana masses and a high scale see-saw mechanism. This is the $U(1)_N$ gauge symmetry given by $\theta = \arctan \sqrt{15}$,
and defines the so-called exceptional supersymmetric standard model (E₆SSM) [9]. The extra $U(1)_N$ gauge symmetry survives to low energies and forbids a bilinear term $\mu H_d H_u$ in the superpotential but allows the interaction $\lambda S H_d H_u$. At the electroweak (EW) scale, the scalar component of the SM singlet superfield $S$ acquires a non-zero VEV, $\langle S \rangle = s/\sqrt{2}$, breaking $U(1)_N$ and yielding an effective $\mu = \lambda s/\sqrt{2}$ term. Thus the $\mu$ problem in the E₆SSM is solved in a similar way to the NMSSM, but without the accompanying problems of singlet tadpoles or domain walls. In this model low energy anomalies are cancelled by complete 27 representations of $E₆$ which survive to low energies, with $E₆$ broken at the high energy GUT scale.

In this Letter we discuss some of the predictions of particular relevance to the LHC from a constrained version of the E₆SSM (cE₆SSM), based on a universal high energy soft scalar mass $m_0$, soft trilinear coupling $A_0$ and soft gaugino mass $M_{1/2}$. Our primary focus is on the most urgent regions of parameter space which involve low values of $(m_0, M_{1/2})$ and low $Z'$ gauge boson masses which would correspond to an early LHC discovery using “first data”. To illustrate these features we propose and discuss a set of “early discovery” benchmark points, each associated with a $Z'$ gauge boson mass around 1 TeV and $(m_0, M_{1/2})$ below 1 TeV, which would lead to an early indication of the cE₆SSM at the LHC. We find a SUSY spectrum consisting of a light gluino of mass $\sim M_3$, a light wino-like neutralino and chargino pair of mass $\sim M_2$, and a light bino-like neutralino of mass $\sim M_1$, where $M_i$ are the low energy gaugino masses, which are typically driven small by renormalisation group (RG) running. Sfermions are generally heavier, but there can be an observable top squark. There may also be light exotic colour triplet charge 1/3 fermions and scalars, whose masses are controlled by independent Yukawa couplings. Some first results have already been trailed at conferences [10] and a longer paper containing full details of the analysis is about to appear [11].

In section 2 we briefly review the E₆SSM, then in section 3 we introduce the cE₆SSM. Section 4 describes the experimental and theoretical constraints and section 5 discusses the aforementioned predictions of the cE₆SSM elucidated by five “early discovery” benchmark points. Section 6 concludes the Letter.

2. The E₆SSM

One of the most important issues in models with additional Abelian gauge symmetries is the cancellation of anomalies. In $E₆$ theories, if the surviving Abelian gauge group factor is a subgroup of $E₆$, and the low energy spectrum constitutes a complete 27 representation of $E₆$, then the anomalies are cancelled automatically. The 27ᵢ of $E₆$, each containing a quark and lepton family, decompose under the $SU(5) \times U(1)_N$ subgroup of $E₆$ as follows:
The first and second quantities in the brackets are the $SU(5)$ representation and extra $U(1)_N$ charge while $i$ is a family index that runs from 1 to 3. From Eq. (1) we see that, in order to cancel anomalies, the low energy (TeV scale) spectrum must contain three extra copies of $5^* + 5$ of $SU(5)$ in addition to the three quark and lepton families in $5^* + 10$. To be precise, the ordinary SM families which contain the doublets of left-handed quarks $Q_i$ and leptons $L_i$, right-handed up- and down-quarks ($u^c_i$ and $d^c_i$) as well as right-handed charged leptons, are assigned to $(10, 1)_i + (5^*, 2)_i$. Right-handed neutrinos $N^c_i$ should be associated with the last term in Eq. (1), $(1, 0)_i$. The next-to-last term in Eq. (1), $(1, 5)_i$, represents SM-type singlet fields $S_i$ which carry non-zero $U(1)_N$ charges and therefore survive down to the EW scale. The three pairs of $SU(2)$-doublets ($H^u_i$ and $H^d_i$) that are contained in $(5^*, -3)_i$ and $(5, -2)_i$ have the quantum numbers of Higgs doublets, and we shall identify one of these pairs with the usual MSSM Higgs doublets, with the other two pairs being “inert” Higgs doublets which do not get VEVs. The other components of these $SU(5)$ multiplets form colour triplets of exotic quarks $D_i$ and $D^c_i$ with electric charges $-1/3$ and $+1/3$ respectively. The matter content and correctly normalised Abelian charge assignment are in Tab. 1.

| $Q$ | $u^c$ | $d^c$ | $L$ | $e^c$ | $N^c$ | $S$ | $H_2$ | $H_1$ | $D$ | $\overline{D}$ | $H'$ | $\overline{H}'$ |
|-----|------|------|-----|------|------|-----|-------|-------|-----|----------|------|----------|
| $\sqrt{\frac{5}{3}} Q_i^{Y}$ | $\frac{1}{6}$ | $-\frac{2}{3}$ | $\frac{1}{3}$ | $-\frac{1}{2}$ | 1 | 0 | 0 | $\frac{1}{2}$ | $-\frac{1}{2}$ | $-\frac{1}{3}$ | $\frac{1}{3}$ | $-\frac{1}{2}$ | $\frac{1}{2}$ |
| $\sqrt{40} Q_i^{N}$ | 1 | 1 | 2 | 2 | 1 | 0 | 5 | $-2$ | $-3$ | $-2$ | $-3$ | 2 | $-2$ |

Table 1: The $U(1)_Y$ and $U(1)_N$ charges of matter fields in the $E_6$ SSM, where $Q_i^{N}$ and $Q_i^{Y}$ are here defined with the correct $E_6$ normalisation factor required for the RG analysis.

We also require a further pair of superfields $H'$ and $\overline{H'}$ with a mass term $\mu' H' \overline{H'}$ from incomplete extra $27'$ and $\overline{27'}$ representations to survive to low energies to ensure gauge coupling unification. Because $H'$ and $\overline{H'}$ originate from $27'$ and $\overline{27'}$ these supermultiplets do not spoil anomaly cancellation in the considered model. Our analysis reveals that the unification of the gauge couplings in the $E_6$ SSM can be achieved for any phenomenologically acceptable value of $\alpha_3(M_Z)$, consistent with the measured low energy central value $^{[12]}$.

Since right–handed neutrinos have zero charges they can acquire very heavy Majorana masses. The heavy Majorana right-handed neutrinos may decay into final states with

\[ 27_i \rightarrow (10, 1)_i + (5^*, 2)_i + (5^*, -3)_i + (5, -2)_i + (1, 5)_i + (1, 0)_i. \]
lepton number $L = \pm 1$, thereby creating a lepton asymmetry in the early Universe. Because the Yukawa couplings of exotic particles are not constrained by the neutrino oscillation data, substantial values of CP–violating lepton asymmetries can be induced even for a relatively small mass of the lightest right–handed neutrino ($M_1 \sim 10^6$ GeV) so that successful thermal leptogenesis may be achieved without encountering any gravitino problem \[14\].

The superpotential of the E$_6$SSM involves a lot of new Yukawa couplings in comparison to the SM. In general these new interactions violate baryon number conservation and induce non-diagonal flavour transitions. To suppress baryon number violating and flavour changing processes one can postulate a $Z^H_2$ symmetry under which all superfields except one pair of $H_d^i$ and $H_u^i$ (say $H_d \equiv H_d^3$ and $H_u \equiv H_u^3$) and one SM-type singlet field ($S \equiv S_3$) are odd. The $Z^H_2$ symmetry reduces the structure of the Yukawa interactions, and an assumed hierarchical structure of the Yukawa interactions allows to simplify the form of the E$_6$SSM superpotential substantially. Keeping only Yukawa interactions whose couplings are allowed to be of order unity leaves us with

$$W_{\text{E}_6\text{SSM}} \simeq \lambda S(H_d H_u) + \lambda_i S(H_d^i H_u^i) + \kappa_i S(D_i \bar{D}_i)$$

$$+ h_t(H_u Q) t^c + h_b(H_d Q) b^c + h_\tau(H_d L) \tau^c + \mu'(H' \bar{H}') ,$$

where $\alpha, \beta = 1, 2$ and $i = 1, 2, 3$, and where the superfields $L = L_3$, $Q = Q_3$, $t^c = u_3^c$, $b^c = d_3^c$ and $\tau^c = e_3^c$ belong to the third generation and $\lambda_i, \kappa_i$ are dimensionless Yukawa couplings with $\lambda \equiv \lambda_3$. Here we assume that all right–handed neutrinos are relatively heavy so that they can be integrated out\[^2\]. The $SU(2)_L$ doublets $H_u$ and $H_d$, which are even under the $Z^H_2$ symmetry, play the role of Higgs fields generating the masses of quarks and leptons after EWSB. The singlet field $S$ must also acquire a large VEV to induce sufficiently large masses for the $Z'$ boson. The couplings $\lambda_i$ and $\kappa_i$ should be large enough to ensure the exotic fermions are sufficiently heavy to avoiding conflict with direct particle searches at present and former accelerators. They should also be large enough so that the evolution of the soft scalar mass $m^2_S$ of the singlet field $S$ results in negative values of $m^2_S$ at low energies, triggering the breakdown of the $U(1)_N$ symmetry.

However the $Z^H_2$ can only be approximate (otherwise the exotics would not be able to decay). To prevent rapid proton decay in the E$_6$SSM a generalised definition of $R$–parity should be used. We give two examples of possible symmetries that can achieve that. If $H_d^i$, $H_u^i$, $S_i$, $D_i$, $\bar{D}_i$, and the quark superfields ($Q_i$, $u_i^c$, $d_i^c$) are even under a discrete $Z^L_2$ symmetry while the lepton superfields ($L_i$, $e_i^c$, $N_i^c$) are odd (Model I) then the allowed superpotential is invariant with respect to a $U(1)_B$ global symmetry. The exotic $\bar{D}_i$ and $D_i$ are then identified as diquark and anti-diquark, i.e. $B_D = -2/3$ and $B_{\bar{D}} = 2/3$. An

\[^2\]We shall ignore the presence of right-handed neutrinos in the subsequent RG analysis.
alternative possibility is to assume that the exotic quarks $D_i$ and $\overline{D_i}$ as well as lepton superfields are all odd under $Z_2^B$ whereas the others remain even. In this case (Model II) the $\overline{D_i}$ and $D_i$ are leptoquarks [9].

After the breakdown of the gauge symmetry, $H_u, H_d$ and $S$ form three CP–even, one CP-odd and two charged states in the Higgs spectrum. The mass of one CP–even Higgs particle is always very close to the $Z'$ boson mass $M_{Z'}$. The masses of another CP–even, the CP–odd and the charged Higgs states are almost degenerate. Furthermore, like in the MSSM and NMSSM, one of the CP–even Higgs bosons is always light irrespective of the SUSY breaking scale. However, in contrast with the MSSM, the lightest Higgs boson in the E$_6$SSM can be heavier than 110 – 120 GeV even at tree level. In the two–loop approximation the lightest Higgs boson mass does not exceed 150 – 155 GeV [9]. Thus the SM–like Higgs boson in the E$_6$SSM can be considerably heavier than in the MSSM and NMSSM, since it contains a similar F-term contribution as the NMSSM but with a larger maximum value for $\lambda(m_t)$ as it is not bounded as strongly by the validity of perturbation theory up to the GUT scale [9]. However in the considered “early discovery” benchmark points in this Letter, it will always be close to the current LEP2 limit.

3. The Constrained E$_6$SSM

The simplified superpotential of the E$_6$SSM involves seven extra couplings ($\mu', \kappa_i$ and $\lambda_i$) as compared with the MSSM with $\mu = 0$. The soft breakdown of SUSY gives rise to many new parameters. The number of fundamental parameters can be reduced drastically though within the constrained version of the E$_6$SSM (cE$_6$SSM). Constrained SUSY models imply that all soft scalar masses are set to be equal to $m_0$ at some high energy scale $M_X$, taken here to be equal to the GUT scale, all gaugino masses $M_i(M_X)$ are equal to $M_{1/2}$ and trilinear scalar couplings are such that $A_i(M_X) = A_0$. Thus the cE$_6$SSM is characterised by the following set of Yukawa couplings, which are allowed to be of the order of unity, and universal soft SUSY breaking terms,

$$\lambda_i(M_X), \kappa_i(M_X), h_t(M_X), h_b(M_X), h_\tau(M_X), m_0, M_{1/2}, A_0,$$

where $h_t(M_X), h_b(M_X)$ and $h_\tau(M_X)$ are the usual $t$–quark, $b$–quark and $\tau$–lepton Yukawa couplings, and $\lambda_i(M_X), \kappa_i(M_X)$ are the extra Yukawa couplings defined in Eq. (2). The universal soft scalar and trilinear masses correspond to an assumed high energy soft SUSY breaking potential of the universal form,

$$V_{soft} = m_0^2 \mathbf{27_i27_i} + A_0 Y_{ijk} \mathbf{27_i27_j27_k} + h.c.,$$

where $Y_{ijk}$ are generic Yukawa couplings from the trilinear terms in Eq. (2) and the $\mathbf{27_i}$ represent generic fields from Eq. (1), and in particular those which appear in Eq. (2).
Since $Z^H_2$ symmetry forbids many terms in the superpotential of the $E_6$ SSM it also forbids similar soft SUSY breaking terms in Eq. (1). To simplify our analysis we assume that all parameters in Eq. (3) are real and $M_{1/2}$ is positive. In order to guarantee correct EWSB $m_0^2$ has to be positive. The set of $cE_6$ SSM parameters in Eq. (3) should in principle be supplemented by $\mu'$ and the associated bilinear scalar coupling $B'$. However, since $\mu'$ is not constrained by the EWSB and the term $\mu' H' \overline{H}'$ in the superpotential is not suppressed by $E_6$, the parameter $\mu'$ will be assumed to be $\sim 10$ TeV so that $H'$ and $\overline{H}'$ decouple from the rest of the particle spectrum. As a consequence the parameters $B'$ and $\mu'$ are irrelevant for our analysis.

To calculate the particle spectrum within the $cE_6$ SSM one must find sets of parameters which are consistent with both the high scale universality constraints and the low scale EWSB constraints. To evolve between these two scales we use two–loop renormalisation group equations (RGEs) for the gauge and Yukawa couplings together with two–loop RGEs for $M_a(Q)$ and $A_i(Q)$ as well as one–loop RGEs for $m_t^2(Q)$. $Q$ is the renormalisation scale. The RGE evolution is performed using a modified version of SOFTSUSY 2.0.5 [15] and the RGEs for the $E_6$ SSM are presented in a longer paper [11]. The details of the procedure we followed are summarized below.

1. The gauge and Yukawa couplings are determined independently of the soft SUSY breaking mass parameters as follows:

   (i) We choose input values for $s = \sqrt{2} \langle S \rangle$ and $\tan \beta = v_2/v_1$ (where $v_2$ and $v_1$ are the usual VEVs of the Higgs fields $H_u$ and $H_d$) as defined by our scenario.

   (ii) We set the gauge couplings $g_1$, $g_2$ and $g_3$ equal to the experimentally measured values at $M_Z$.

   (iii) We fix the low energy Yukawa couplings $h_t$, $h_b$, and $h_\tau$ using the relations between the running masses of the fermions of the third generation and VEVs of the Higgs fields, i.e.

   \[ m_t(M_t) = \frac{h_t(M_t) v}{\sqrt{2}} \sin \beta, \quad m_b(M_t) = \frac{h_b(M_t) v}{\sqrt{2}} \cos \beta, \quad m_\tau(M_t) = \frac{h_\tau(M_t) v}{\sqrt{2}} \cos \beta. \quad (5) \]

   (iv) The gauge and Yukawa couplings are then evolved up to the GUT scale $M_X$. Using the beta functions for QED and QCD, the gauge couplings are first evolved up to $m_t$. Since we are employing two–loop RGEs in the SUSY preserving sector, we include one estimated threshold scale for the masses of the superpartners of the SM particles, $T_{MSSM}$, and one for the masses of the new exotic particles, $T_{ESSM}$. Since these are common mass scales we neglect mass splitting within each group of particles. So between $m_t$ and $T_{MSSM}$ we evolve these gauge and Yukawa couplings with SM RGEs and between $T_{MSSM}$ and $T_{ESSM}$ we employ the MSSM RGEs. At $T_{ESSM}$ the values of $E_6$ SSM gauge and Yukawa couplings, $g_1$, $g_2$, $g_3$, $h_t$, $h_b$ and $h_\tau$, form a low energy boundary condition
for what follows. Initial low energy estimates of the new \( E_6 \) SSM Yukawa couplings, \( \lambda_i \) and \( \kappa_i \), are also input here, and all SUSY preserving couplings are evolved up to the high scale using the two–loop \( E_6 \) SSM RGEs.

(v) At the GUT scale \( M_X \) we set \( g_1(M_X) = g_2(M_X) = g_3(M_X) = g_1'(M_X) \equiv g_0 \) and select values for \( \kappa_i(M_X) \) and \( \lambda_i(M_X) \), which are input parameters in our procedure. An iteration is then performed between \( M_X \) and the low energy scale to obtain the values of all the gauge and Yukawa couplings which are consistent with our input values for \( \kappa_i(M_X) \), \( \lambda_i(M_X) \), gauge coupling unification and our low scale boundary conditions, derived from experimental data.

2. Having completely determined the gauge and Yukawa couplings, the low energy soft SUSY breaking parameters are then determined semi-analytically as functions of the GUT scale values of \( A_0, M_{1/2} \) and \( m_0 \). They take the form,

\[
m_i^2(Q) = a_i(Q)M_{1/2}^2 + b_i(Q)A_0^2 + c_i(Q)A_0M_{1/2} + d_i(Q)m_0^2, \quad (6)
\]

\[
A_i(Q) = e_i(Q)A_0 + f_i(Q)M_{1/2}, \quad (7)
\]

\[
M_i(Q) = p_i(Q)A_0 + q_i(Q)M_{1/2}, \quad (8)
\]

where \( Q \) is the renormalisation scale. The coefficients are unknown but may be determined numerically at the low energy scale, as follows:

(i) Set \( A_0 = M_{1/2} = 0 \) at \( M_X \) with \( m_0 \) non-zero, and run the full set of \( E_6 \) SSM parameters down to the low scale to yield the coefficients proportional to \( m_0^2 \) in the expressions for the soft SUSY breaking parameters.

(ii) Repeat for \( A_0 \) and \( M_{1/2} \).

(iii) The coefficient of the \( A_0M_{1/2} \) term is determined using non-zero values of both \( A_0 \) and \( M_{1/2} \) at \( M_X \), using the results in part (ii) to isolate this term.

3. Using the semi-analytic expressions for the soft masses from step 2 above, we then impose conditions for correct EWSB at low energy and determine sets of \( m_0, M_{1/2} \) and \( A_0 \) which are consistent with EWSB, as follows:

(i) Working with the tree–level potential \( V_0 \) (to start with) we impose the minimisation conditions \( \frac{\partial V_0}{\partial s} = \frac{\partial V_0}{\partial v_1} = \frac{\partial V_0}{\partial v_2} = 0 \) leading to a system of quadratic equations in \( m_0, M_{1/2} \) and \( A_0 \). In this approximation, the equations can be reduced to two second order equations with respect to \( A_0 \) and \( M_{1/2} \) which can have up to four solutions for each set of Yukawa couplings.

(ii) For each solution \( m_0, M_{1/2} \) and \( A_0 \), the low energy stop soft mass parameters are determined and the one–loop Coleman-Weinberg Higgs effective potential \( V_1 \) is calculated. The new minimisation conditions for \( V_1 \) are then imposed, and new solutions for \( m_0, M_{1/2} \) and \( A_0 \) are obtained.
(iii) The procedure in (ii) is then iterated until we find stable solutions. Some or all of the obtained solutions can be complex. Here we restrict our consideration to the scenarios with real values of fundamental parameters which do not induce any CP-violating effects. For some values of $\tan \beta$, $s$ and Yukawa couplings the solutions with real $A_0$, $M_{1/2}$ and $m_0$ do not exist. There is a substantial part of the parameter space where there are only two solutions with real values of fundamental parameters. However there are also some regions of the parameters where all four solutions of the non-linear algebraic equations are real.

Although correct EWSB is not guaranteed in the cE$\_6$SSM, remarkably, there are always solutions with real $A_0$, $M_{1/2}$ and $m_0$ for sufficiently large values of $\kappa_i$, which drive $m_S^2$ negative. This is easy to understand since the $\kappa_i$ couple the singlet to a large multiplicity of coloured fields, thereby efficiently driving its squared mass negative to trigger the breakdown of the gauge symmetry.

4. Using the obtained solutions we calculate the masses of all exotic and SUSY particles for each set of fundamental parameters in Eq. (3).

Finally at the last stage of our analysis we vary Yukawa couplings, $\tan \beta$ and $s$ to establish the qualitative pattern of the particle spectrum within the cE$\_6$SSM. To avoid any conflict with present and former collider experiments as well as with recent cosmological observations we impose the set of constraints specified in the next section. These bounds restrict the allowed range of the parameter space in the cE$\_6$SSM.

4. Experimental and Theoretical Constraints

The experimental constraints applied in our analysis are: $m_h \geq 114$ GeV, all sleptons and charginos are heavier than 100 GeV, all squarks and gluinos have masses above 300 GeV and the $Z'$ boson has a mass which is larger than 861 GeV [16]. We also impose the most conservative bound on the masses of exotic quarks and squarks that comes from the HERA experiments [17], by requiring them to be heavier than 300 GeV. Finally we require that the inert Higgs and inert Higgsinos are heavier than 100 GeV to evade limits on Higgsinos and charged Higgs bosons from LEP.

In addition to a set of bounds coming from the non-observation of new particles in experiments we impose a few theoretical constraints. We require that the lightest SUSY particle (LSP) should be a neutralino. We also restrict our consideration to values of the Yukawa couplings $\lambda_i(M_X)$, $\kappa_i(M_X)$, $h_t(M_X)$, $h_b(M_X)$ and $h_\tau(M_X)$ less than 3 to ensure the applicability of perturbation theory up to the GUT scale.

In our exploration of the cE$\_6$SSM parameter space we first looked at scenarios with
a universal coupling between exotic coloured superfields and the third generation singlet field $S$, $\kappa(M_X) = \kappa_{1,2,3}(M_X)$, and fixed the inert Higgs couplings $\lambda_{1,2}(M_X) = 0.1$. In fixing $\lambda_{1,2}$ like this we are deliberately pre-selecting for relatively light inert Higgsinos.

The third generation Yukawa $\lambda = \lambda_3$ was allowed to vary along with $\kappa$. Splitting $\lambda_3$ from $\lambda_{1,2}$ seems reasonable since $\lambda_3$ plays a very special role in $E_6$SSM models in forming the effective $\mu$-term when $S$ develops a VEV. Eventually we allowed non–universal $\kappa_i(M_X)$.

For fixed values of $\tan \beta = 3, 10, 30$, we scanned over $s, \kappa, \lambda$. From these input parameters, the sets of soft mass parameters, $A_0, M_{1/2}$ and $m_0$, which are consistent with the correct breakdown of electroweak symmetry, are found. We find that for fixed values of the Yukawas the soft mass parameters scale with $s$, while if $s$ and $\tan \beta$ are fixed, varying the Yukawas, $\lambda$ and $\kappa$, then produces a bounded region of allowed points. The value of $s$ determines the location and extent of the bounded regions. As $s$ is increased the lowest values of $m_0$ and $M_{1/2}$, consistent with experimental searches and EWSB requirements, increase. This is shown in Fig. 1 where the allowed regions for three different values of the singlet VEV, $s = 3, 4$ and $5$ TeV, are compared, with the allowed regions in red, green, magenta respectively and the excluded regions in white. These regions overlap since we are finding soft masses consistent with EWSB conditions that have a non–linear dependence on the VEVs and Yukawas.

5. Predictions of the cE$_6$SSM

5.1 Overview of the spectrum and decay signatures

5.1.1 SUSY spectrum and signatures

From Fig. 1 we see that $m_0 > M_{1/2}$ for each value of $s$ and also that lower $M_{1/2}$ is weakly correlated with lower $s$ and thus lower $Z'$ masses. As is discussed in detail in Ref. [11] this bound is caused, depending on the value of $\tan \beta$, either by the inert Higgs masses being driven below their experimental limit from negative D-term contributions canceling the positive contribution from $m_0$ or the light Higgs mass going below the LEP2 limit.

Another remarkable feature of the cE$_6$SSM is that the low energy gluino mass parameter $M_3$ is driven to be smaller than $M_{1/2}$ by RG running. The reason for this is that the E$_6$SSM has a much larger (super)field content than the MSSM (three 27’s instead of three 16’s), so much so that at one–loop order the QCD beta function (accidentally) vanishes in the E$_6$SSM, and at two loops it loses asymptotic freedom (though the gauge couplings remain perturbative at high energy). This implies that the low energy gaugino masses are all less than $M_{1/2}$ in the cE$_6$SSM, being given by roughly $M_3 \sim 0.7 M_{1/2}$, $M_2 \sim 0.25 M_{1/2}$, $M_1 \sim 0.15 M_{1/2}$. These should be compared to the corresponding low energy values in the
Figure 1: Physical solutions with $\tan \beta = 10$, $\lambda_{1,2} = 0.1$, $s = 3, 4, 5$ TeV fixed and $\lambda \equiv \lambda_3$ and $\kappa \equiv \kappa_{1,2,3}$ varying, which pass experimental constraints from LEP and Tevatron data. On the left-hand side of each allowed region the chargino mass is less than 100 GeV, while underneath the inert Higgses are less than 100 GeV or becoming tachyonic. The region ruled out immediately to the right of the allowed points is due to $m_h < 114$ GeV. The results show that $m_0 > M_{1/2}$ for each value of $s$. They also show that higher $M_{1/2}$ are correlated with higher $s$ (and thus higher $Z'$ masses).

MSSM, $M_3 \sim 2.7M_{1/2}$, $M_2 \sim 0.8M_{1/2}$, $M_1 \sim 0.4M_{1/2}$. Thus, in the cE6SSM, since the low energy gaugino masses $M_i$ are driven by RG running to be small, the lightest SUSY states will generally consist of a light gluino of mass $\sim M_3$, a light wino-like neutralino and chargino pair of mass $\sim M_2$, and a light bino-like neutralino of mass $\sim M_1$, which are typically all much lighter than the Higgsino masses of order $\mu = \lambda s/\sqrt{2}$, where $\lambda$ cannot be too small for correct EWSB. Since $m_0 > M_{1/2}$ the squarks and sleptons are also much heavier than the light gauginos.

Thus, throughout all cE6SSM regions of parameter space there is the striking prediction that the lightest sparticles always include the gluino $\tilde{g}$, the two lightest neutralinos $\chi_1^0, \chi_2^0$, and a light chargino $\chi_1^\pm$. Therefore pair production of $\chi_2^0 \chi_2^0, \chi_2^0 \chi_1^\pm, \chi_1^\pm \chi_1^\mp$ and $\tilde{g} \tilde{g}$ should always be possible at the LHC irrespective of the $Z'$ mass. Due to the hierarchical spectrum, the gluinos can be relatively narrow states with width $\Gamma_\tilde{g} \propto M_\tilde{g}^5/m_\tilde{g}^4$, comparable to that of $W^\pm$ and $Z$ bosons. They will decay through $\tilde{g} \rightarrow q\tilde{q}^* \rightarrow q\bar{q} + E_T^{miss}$, so gluino pair production will result in an appreciable enhancement of the cross section for $pp \rightarrow q\bar{q}q\bar{q} + E_T^{miss} + X$, where $X$ refers to any number of light quark/gluon jets.
The second lightest neutralino decays through $\chi_2^0 \rightarrow \chi_1^0 + l\bar{l}$ and so would produce an excess in $pp \rightarrow l\bar{l}l\bar{l} + E_T^{\text{miss}} + X$, which could be observed at the LHC. Since all squarks and sleptons, as well as new exotic particles, turn out to be rather heavy compared to the low energy wino mass, the calculation of the branching ratio $Br(\chi_2^0 \rightarrow \chi_1^0 + l\bar{l})$ is very similar to that in the MSSM. This branching ratio in the MSSM is known to be very sensitive to the choice of fundamental parameters of the model. For the type of the neutralino spectra presented later, in which the second lightest neutralino is approximately wino, the lightest neutralino is approximately bino, and where the other sparticles are much heavier, $Br(\chi_2^0 \rightarrow \chi_1^0 + l\bar{l})$ is known to vary from $1.5\%$ to $6\%$ [18].

5.1.2 Exotic spectrum and signatures

Other possible manifestations of the $E_6$SSM at the LHC are related to the presence of a $Z'$ and exotic multiplets of matter. The production of a TeV scale $Z'$ will provide an unmistakable and spectacular LHC signal even with first data [9]. At the LHC, the $Z'$ boson that appears in the $E_6$ inspired models can be discovered if it has a mass below $4-4.5$ TeV [19]. The determination of its couplings should be possible if $M_{Z'} \lesssim 2-2.5$ TeV [20].

When the Yukawa couplings $\kappa_i$ of the exotic fermions $D_i$ and $\bar{D}_i$ have a hierarchical structure, some of them can be relatively light so that their production cross section at the LHC can be comparable with the cross section of $tt$ production [9]. In the $E_6$SSM, the $D_i$ and $\bar{D}_i$ fermions are SUSY particles with negative $R$–parity so they must be pair produced and decay into quark–squark (if diquarks) or quark–slepton, squark–lepton (if leptoquarks), leading to final states containing missing energy from the LSP.

The lifetime and decay modes of the exotic coloured fermions are determined by the $Z_2^H$ violating couplings. If $Z_2^H$ is broken significantly the presence of the light exotic quarks gives rise to a remarkable signature. Assuming that $D_i$ and $\bar{D}_i$ fermions couple most strongly to the third family (s)quarks and (s)leptons, the lightest exotic $D_i$ and $\bar{D}_i$ fermions decay into $\bar{t}b$, $\bar{t}b$, $\bar{t}\bar{t}$, $\bar{t}\bar{b}$ (if they are diquarks) or $\bar{t}\tau$, $t\bar{\tau}$, $b\nu$, $b\bar{\nu}$ (if they are leptoquarks). This can lead to a substantial enhancement of the cross section of either $pp \rightarrow t\bar{t}b\bar{b} + E_T^{\text{miss}} + X$ (if diquarks) or $pp \rightarrow t\bar{t}\tau\bar{\tau} + E_T^{\text{miss}} + X$ or $pp \rightarrow \bar{b}\bar{b} + E_T^{\text{miss}} + X$ (if leptoquarks). Notice that SM production of $t\bar{t}\tau^+\tau^-$ is $(a_W/\pi)^2$ suppressed in comparison to the leptoquark decays. Therefore light leptoquarks should produce a strong signal with low SM background at the LHC. In principle the detailed LHC analyses is required to establish the feasibility of extracting the excess of $t\bar{t}b\bar{b}$ or $t\bar{t}\tau^+\tau^-$ production induced by the light exotic quarks predicted by our model.

We have already remarked that the lifetime and decay modes of the exotic coloured fermions are determined by the $Z_2^H$ violating couplings. If $Z_2^H$ is only very slightly broken
exotic quarks may be very long lived, with lifetimes up to about 1 s. This is the case, for example, in some minimal versions of the model [13]. In this case the exotic $D_i$ and $\bar{D}_i$ fermions could hadronize before decaying, leading to spectacular signatures consisting of two low multiplicity jets, each containing a single quasi-stable heavy D-hadron, which could be stopped for example in the muon chambers, before decaying much later.

In Tab. 2 we estimate the total production cross section of exotic quarks at the LHC for a few different values of their masses assuming that all masses of exotic quarks are equal (i.e. $\mu_{D_i} = \mu_D$) and all sparticles as well as other new exotic particles are very heavy. The results in Tab. 2 suggest that the observation of the $D$ fermions might be possible if they have masses below about 1.5-2 TeV [9].

| $\mu_D$ [GeV] | 300 | 400 | 500 | 700 | 1000 | 1500 | 2000 | 3000 |
|---------------|-----|-----|-----|-----|------|------|------|------|
| $\sigma(pp \rightarrow D\bar{D})$ [pb] | 76.4 | 17.4 | 5.30 | 0.797 | 0.0889 | 4.94 $\cdot$ 10^{-3} | 4.09 $\cdot$ 10^{-4} | 3.51 $\cdot$ 10^{-6} |

Table 2: The cross section of $D\bar{D}$ production at the LHC as a function of the masses of exotic quarks. For simplicity we assume that three families of exotic quarks have the same masses.

Similar considerations apply to the case of exotic $\tilde{D}_i$ and $\tilde{D}_i$ scalars except that they are non–SUSY particles so they may be produced singly and decay into quark–quark (if diquarks) or quark–lepton (if leptoquarks) without missing energy from the LSP. It is possible to have relatively light exotic coloured scalars due to mixing effects. The RGEs for the soft breaking masses, $m_{\tilde{D}_i}^2$ and $m_{\tilde{D}_i}^2$, are very similar, with $\frac{d}{dt}(m_{\tilde{D}_i}^2 - m_{\tilde{D}_i}^2) = g_1^2 M_1^2$, resulting in comparatively small splitting between these soft masses. Consequently, mixing can be large even for moderate values of the $A_0$, leading to a large mass splitting between the two scalar partners of the exotic coloured fermions. Recent, as yet unpublished, results from Tevatron searches for dijet resonances [21] rule out scalar diquarks with mass less than 630 GeV. However, scalar leptoquarks may be as light as 300 GeV since at hadron colliders they are pair produced through gluon fusion. Scalar leptoquarks decay into quark–lepton final states through small $ZH_2$ violating terms, for example $\tilde{D} \rightarrow t\tau\bar{\tau}$, and pair production leads to an enhancement of $pp \rightarrow t\bar{t}\tau\bar{\tau}$ (without missing energy) at the LHC.

In addition, the inert Higgs bosons and Higgsinos (i.e. the first and second families of Higgs doublets predicted by the $E_6$SSM which couple weakly to quarks and leptons and do not get VEVs) can be light or heavy depending on their free parameters. The light inert Higgs bosons decay via $Z_2^H$ violating terms which are analogous to the Yukawa

Note that the diagonal entries of the exotic squark mass matrices have substantial negative contributions from the $U(1)_N$ D–term quartic interactions in the scalar potential. These contributions reduce the masses of exotic squarks and also contribute to their mass splitting since the $U(1)_N$ charges of $D_i$ and $\bar{D}_i$ are different.
interactions of the Higgs superfields, $H_u$ and $H_d$. One can expect that the couplings of the inert Higgs fields would have a similar hierarchical structure as the couplings of the normal Higgs multiplets, therefore we assume the $Z_2^H$ breaking interactions predominantly couple the inert Higgs bosons to the third generation. So the neutral inert Higgs bosons decay predominantly into 3rd generation fermion–anti-fermion pairs like $H^0_{1,i} \rightarrow b \bar{b}$. The charged inert Higgs bosons also decay into fermion–anti-fermion pairs, but in this case it is the antiparticle of the fermions’ EW partner, e.g. $H^-_{1,i} \rightarrow \tau \bar{\nu}_\tau$. The inert Higgs bosons may also be quite heavy, so that the only light exotic particles are the inert Higgsinos. Similar couplings govern the decays of the inert Higgsinos; the electromagnetically neutral Higgsinos predominantly decay into fermion-anti-sfermion pairs (e.g. $\tilde{H}_0^i \rightarrow t \tilde{t}^*, \tilde{H}_0^i \rightarrow \tau \tilde{\nu}_\tau^*$). The charged Higgsinos decays similarly but in this case the sfermion is the SUSY partner of the EW partner of the fermion (e.g. $\tilde{H}_i^+ \rightarrow t \tilde{b}^*, \tilde{H}_i^- \rightarrow \tau \tilde{\nu}_\tau^*$).

5.2 Early discovery benchmarks

5.2.1 The benchmark input parameters

In Tab. 3 we present a set of “early discovery” benchmark points, each associated with a $Z'$ gauge boson mass close to 1 TeV which should be discovered using first LHC data. The first block of Tab. 3 shows the input parameters which define the benchmark points. We have selected $s = 2.7 - 3.3$ TeV corresponding to $M_{Z'} = 1020 - 1250$ GeV, where $M_{Z'} \approx g'_1 s Q_S$ with $Q_S = 5/\sqrt{40}$ and $g'_1 \approx g_1$. We have also restricted ourselves to $(m_0, M_{1/2}) < (700, 400)$ GeV leading to very light gauginos, associated with the three low gaugino masses $M_i$, and in addition a light stop and Higgs mass. Note that for all the benchmark points the trilinear soft mass is fixed to lie in the range $A_0 = 650 - 1150$ GeV in order to achieve EWSB.

The benchmark points cover three different values of $\tan \beta = 3, 10, 30$. In each case we have taken $|\lambda_3|$ to be larger than $\lambda_{1,2} = 0.1$ (fixed) at the GUT scale. In benchmark points A, B, E (corresponding to $\tan \beta = 3, 10, 30$) we have taken the $\kappa_i$ to be universal at the GUT scale and large enough to trigger EWSB. Since the $\kappa_i$’s control the exotic coloured fermion masses, this implies that all the $D_i$ and $\overline{D}_i$ fermions are all very heavy in these cases. However it is not necessary for the $\kappa_i$’s to be universal and these Yukawa couplings may be hierarchical as for the quark and lepton couplings. To illustrate this possibility we have considered two benchmark points, C and D, both for $\tan \beta = 10$, in which $\kappa_3 \gg \kappa_{1,2}$ at the GUT scale. In these points C, D we have taken $\kappa_3$ to be large enough to trigger EWSB, while allowing $\kappa_{1,2}$ to be low enough to yield light $D_{1,2}$ and $\overline{D}_{1,2}$ fermion masses.
5.2.2 The benchmark spectra

The full spectrum for each of the benchmark points is given in Tab. 3 and illustrated in Fig. 2. The benchmark points all exhibit the characteristic SUSY spectrum described above of a light gluino $\tilde{g}$, two light neutralinos $\chi^0_1, \chi^0_2$, and a light chargino $\chi^\pm_1$. The lightest neutralino $\chi^0_1$ is essentially pure bino, while $\chi^0_2$ and $\chi^\pm_1$ are the degenerate components of the wino. Since $M_{1/2} < 400$ GeV for all the points the (two-loop corrected) gluino mass is below 350 GeV, and the wino mass just above the LEP2 limit of 100 GeV, while the bino is around 60 GeV in each case. The question of the resulting cosmological dark matter relic abundance is not considered in this Letter but one should not regard such points with a light bino as being excluded by cosmology for reasons that will be discussed later.

The Higgsino states are much heavier with the degenerate Higgsinos $\chi^0_{3,4}$ and $\chi^\pm_2$ having masses given by $\mu = \lambda s/\sqrt{2}$ in the range 675–830 GeV for all the benchmark points. The remaining neutralinos are dominantly singlet Higgsinos with masses approximately given by $M_{Z'}$.

The Higgs spectrum for all the benchmark points contains a very light SM–like CP–even Higgs boson $h_1$ with a mass close to the LEP limit of 115 GeV, making it accessible to LHC or even Tevatron. The heavier CP–even Higgs $h_2$, the CP–odd Higgs $A_0$, and the charged Higgs $H^\pm$ are all closely degenerate with masses in the range 600–1000 GeV making them difficult to discover. The remaining mainly singlet CP–even Higgs $h_3$ is closely degenerate with the $Z'$. For benchmarks A, B, E (corresponding to $\tan \beta = 3, 10, 30$) we have taken the $\kappa_i$ to be universal and the exotic coloured fermions have masses in the range 1–1.5 TeV. However, due to the mixing effect mentioned previously, we find a light exotic coloured scalar with a mass of 393 GeV for point E and one at 628 GeV for B. For benchmark points C and D, with $\kappa_3 \gg \kappa_{1,2}$ at the GUT scale, there are light exotic coloured fermions in the range 300–400 GeV, together with a light exotic coloured scalar as before.

The inert Higgs masses may be very light depending on the particular parameters chosen. For example, for benchmarks B and E the lightest inert Higgs bosons of the first and second generation have relatively small masses ($m_{H^0_{1,i}} = 154$ GeV and $m_{H^0_{1,i}} = 220$ GeV respectively). For all the benchmarks the inert Higgsinos are light, as $\mu_{\tilde{H}^i} = 230–300$ GeV.

The lightest stop mass is in the range 430 – 550 GeV for all the benchmark points, with the remaining squark and slepton masses being all significantly heavier than the stop mass but below 1 TeV. Note that the gluino mass, being below 350 GeV, is always lighter than all the squark masses for all the benchmark points.
6. Conclusions

We have discussed the predictions of a constrained version of the exceptional supersymmetric standard model (cE₆SSM), based on a universal high energy soft scalar mass $m₀$, soft trilinear mass $A₀$ and soft gaugino mass $M₁/₂$. We have seen that the cE₆SSM predicts a characteristic SUSY spectrum containing a light gluino, a light wino-like neutralino and chargino pair, and a light bino-like neutralino, with other sparticle masses except the lighter stop being much heavier. In addition, the cE₆SSM allows the possibility of light exotic colour triplet charge $1/3$ fermions and scalars, leading to early exotic physics signals at the LHC.

We have focussed on the possibility of low values of $(m₀, M₁/₂) < (700, 400) \text{ GeV}$, and a $Z'$ gauge boson with mass close to 1 TeV, which would correspond to an early LHC discovery using “first data”, and have proposed a set of benchmark points to illustrate this in Tab. 3 and Fig. 2. For some of the benchmarks (C and D) there are exotic colour triplet charge $±1/3 \ D$ fermions and scalars below 500 GeV, with distinctive final states as discussed in Section 5.1.2. All the benchmark points have a SM–like Higgs close to the LEP2 limit of 115 GeV with the rest of the Higgs spectrum significantly heavier. The inert Higgs bosons may be relatively light, but will be difficult to produce, having zero VEVs and small couplings to quarks and leptons. The lightest stop mass is in the range 430 – 550 GeV for all the benchmark points, with the remaining squark and slepton masses being all significantly heavier than the stop mass but below 1 TeV. The gluino mass is very light, being below 350 GeV in all cases, and in particular is lighter than the stop squark for all the benchmark points. The chargino and second neutralino masses are just above the LEP2 limit of 100 GeV, while the lightest neutralino is around 60 GeV.

We have not considered the question of cosmological cold dark matter (CDM) relic abundance due to the neutralino LSP and so one may be concerned that a bino-like lightest neutralino mass of around 60 GeV might give too large a contribution to $Ω_{CDM}$. Indeed a recent calculation of $Ω_{CDM}$ in the USSM [22], which includes the effect of the MSSM states plus the extra $Z'$ and the active singlet $S$, together with their superpartners, indicates that for the benchmarks considered here that $Ω_{CDM}$ would be too large. However the USSM does not include the effect of the extra inert Higgs and Higgsinos that are present in the E₆SSM. While we have considered the inert Higgsino masses given by $µ_{Hα} = λₐs/√2$, we have not considered the mass of the inert singlinos. These are generated by mixing with the Higgs and inert Higgsinos, and are thus of order $fv²/s$, where $f$ are additional Yukawa couplings that we have not specified in our analysis. Since $s ≫ v$ it is quite likely that the LSP neutralino in the cE₆SSM will be an inert singlino with a mass lighter than 60 GeV. This would imply that the state $χ⁰₁$ considered here is not cosmologically stable but
would decay into lighter (essentially inert) singlinos. Such inert singlinos can annihilate via an s-channel Z-boson, due to their doublet component, yielding an acceptable CDM relic abundance, as has been recently been demonstrated in the $E_6$SSM \cite{23}. The question of the calculation of the relic abundance of such an LSP singlino within the framework of the $cE_6$SSM is beyond the scope of this Letter and will be considered elsewhere. In summary, it is clear that one should not regard the benchmark points with $|m_{\chi^0_1}| \approx 60$ GeV as being excluded by $\Omega_{CDM}$.

To conclude, in this Letter we have argued that the $cE_6$SSM is a very well motivated SUSY model and leads to distinctive predictions at the LHC. We have presented sample benchmark points for which not only the Higgs boson, but also SUSY particles such as gauginos and stop, and even more exotic states such as a light $Z'$ and colour triplet charge $\pm 1/3$ $D$ fermions and scalars, could be just around the corner in early LHC data. If such states are discovered, this would not only represent a revolution in particle physics, but would also point towards an underlying high energy $E_6$ gauge structure, providing a window into string theory.

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|                  | A    | B   | C   | D   | E   |
|------------------|------|-----|-----|-----|-----|
| $\tan \beta$     | 3    | 10  | 10  | 10  | 30  |
| $\lambda_3(M_X)$ | -0.465 | -0.37 | -0.378 | -0.395 | -0.38 |
| $\lambda_{1,2}(M_X)$ | 0.1  | 0.1  | 0.1  | 0.1  |     |
| $\kappa_3(M_X)$  | 0.3  | 0.2  | 0.42 | 0.43 | 0.17 |
| $\kappa_{1,2}(M_X)$ | 0.3  | 0.2  | 0.06 | 0.08 | 0.17 |
| $s[\text{TeV}]$  | 3.3  | 2.7  | 2.7  | 2.7  | 3.1  |
| $M_{1/2}[\text{GeV}]$ | 365  | 363  | 388  | 358  | 365  |
| $m_0[\text{GeV}]$ | 640  | 537  | 681  | 623  | 702  |
| $A_0[\text{GeV}]$ | 798  | 711  | 645  | 757  | 1148 |

|                  | A    | B   | C   | D   | E   |
|------------------|------|-----|-----|-----|-----|
| $m_{\tilde{B}_1}(3)[\text{GeV}]$ | 1797 | 628  | 1465 | 1445 | 393  |
| $m_{\tilde{B}_2}(3)[\text{GeV}]$ | 1156 | 1439 | 2086 | 2059 | 1617 |
| $\mu_D(3)[\text{GeV}]$  | 1466 | 1028 | 1747 | 1747 | 1055 |
| $m_{\tilde{D}_1}(1,2)[\text{GeV}]$ | 1797 | 628  | 520  | 370  | 393  |
| $m_{\tilde{D}_2}(1,2)[\text{GeV}]$ | 1156 | 1439 | 906  | 916  | 1617 |
| $\mu_D(1,2)[\text{GeV}]$  | 1466 | 1028 | 300  | 391  | 1055 |

|                  | A    | B   | C   | D   | E   |
|------------------|------|-----|-----|-----|-----|
| $|m_{\chi_0^3}|[\text{GeV}]$ | 1278 | 1052 | 1054 | 1051 | 1203 |
| $m_{h_3} \simeq M_Z[\text{GeV}]$ | 1248 | 1020 | 1021 | 1127 |     |
| $|m_{\chi_2^3}|[\text{GeV}]$ | 1220 | 993  | 994  | 994  | 1143 |

|                  | A    | B   | C   | D   | E   |
|------------------|------|-----|-----|-----|-----|
| $m_S(1,2)[\text{GeV}]$ | 1097 | 908  | 1001 | 961  | 1093 |
| $m_{H_2}(1,2)[\text{GeV}]$ | 468  | 479  | 627  | 561  | 704  |
| $m_{H_1}(1,2)[\text{GeV}]$ | 165  | 154  | 459  | 345  | 220  |
| $\mu_{\tilde{D}}(1,2)[\text{GeV}]$ | 249  | 244  | 233  | 229  | 298  |

|                  | A    | B   | C   | D   | E   |
|------------------|------|-----|-----|-----|-----|
| $m_{\tilde{u}_1}(1,2)[\text{GeV}]$ | 893  | 788  | 911  | 845  | 929  |
| $m_{\tilde{d}_1}(1,2)[\text{GeV}]$ | 910  | 807  | 929  | 862  | 945  |
| $m_{\tilde{u}_2}(1,2)[\text{GeV}]$ | 910  | 807  | 929  | 862  | 945  |
| $m_{\tilde{d}_2}(1,2)[\text{GeV}]$ | 975  | 850  | 964  | 903  | 998  |
| $m_{\tilde{u}_3}(1,2,3)[\text{GeV}]$ | 874  | 733  | 849  | 796  | 900  |
| $m_{\tilde{d}_3}(1,2,3)[\text{GeV}]$ | 762  | 631  | 765  | 708  | 804  |
| $m_{\tilde{b}_1}[\text{GeV}]$ | 974  | 841  | 955  | 894  | 890  |
| $m_{\tilde{b}_2}[\text{GeV}]$ | 758  | 668  | 777  | 712  | 694  |
| $m_{\tilde{t}_1}[\text{GeV}]$ | 821  | 734  | 829  | 772  | 773  |
| $m_{\tilde{t}_2}[\text{GeV}]$ | 493  | 433  | 546  | 474  | 463  |

|                  | A    | B   | C   | D   | E   |
|------------------|------|-----|-----|-----|-----|
| $|m_{\chi_1^3}| \simeq |m_{\chi_4^3}|[\text{GeV}]$ | 832  | 684  | 674  | 685  | 803  |
| $m_{b_2} \simeq m_A \simeq m_{H^2}[\text{GeV}]$ | 615  | 664  | 963  | 720  | 593  |
| $m_{b_1}[\text{GeV}]$ | 114  | 115  | 115  | 114  | 119  |

|                  | A    | B   | C   | D   | E   |
|------------------|------|-----|-----|-----|-----|
| $m_{\tilde{g}}[\text{GeV}]$ | 336  | 330  | 353  | 327  | 338  |
| $|m_{\chi_2^3}| \simeq |m_{\chi_4^3}|[\text{GeV}]$ | 107  | 103  | 109  | 101  | 103  |
| $|m_{\chi_1^3}|[\text{GeV}]$ | 59   | 58   | 61   | 57   | 58   |

Table 3: The “early discovery” eE6 SSM benchmark points.
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Figure 2: Spectra for the “early discovery” benchmark points A (top left), B (top right), C (middle left), D (middle right) and E (bottom centre).