GAUGINOS AND HIGGS PARTICLES IN SUPERSTRING MODELS

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ABSTRACT

Superstrings compactified on Calabi-Yau manifolds can yield additional light gauginos, Higgs fermions and bosons beyond those in the supersymmetric Standard Model. We explore the phenomenological constraints on minimal model parameters imposed by the experimental non-observation of charginos and charged Higgs bosons, as well as by cosmological constraints on neutralinos. We find that one or more charged fermions, charged bosons, and neutral bosons could be light enough to be produced at LEP. The couplings and production mechanisms for these particles are discussed.

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1. INTRODUCTION

Even if the heterotic superstring [1] is indeed the unique Theory of Everything (TOE), how it makes contact with four-dimensional reality at energies less than 1 TeV is by no means clear. One of the key uncertainties is how the six surplus dimensions are compactified, and the manner in which this is done determines the quantum numbers of the light particles in the theory. Therefore, at our present stage of understanding there is still considerable ambiguity in the choice of new particles to be predicted by the superstring. There are a few generic possibilities, but no one can yet be sure that any one of them will turn out to be a definite signature of the superstring. For the moment, it is appropriate to investigate all the different possibilities and explore their different phenomenological consequences. The superstring can only be shown to be the true TOE if experiment confirms some distinctive prediction, and we must look for one wherever we can.

One of the favoured possibilities [2] for compactification is some six-dimensional manifold with a complex structure, SU(3) holonomy, and a gauge connection identified with the spin connection. Such Calabi–Yau [3] compactifications predict that each matter generation contains 27 fields with the quantum numbers of the fundamental representation $\underline{27}$ of E_6 , supplemented in general by additional pairs of conjugate fields from $(\underline{27} + \overline{\underline{27}})$ representations as well as by gauge singlet fields. In addition, Calabi–Yau compactification predicts that the four-dimensional gauge group has rank five or more [2, 4]. However, which group of rank 5 or 6 is chosen depends on details of the compactification which are as yet unknown. Other scenarios for compactification have been proposed [5, 6], in which matter particles occupy representations of SO(10) or SU(5). However, none of these has yet been worked out in detail, the vacuum stability of some of them is problematic [7, 8] and they have other phenomenological difficulties [8]. Therefore, we do not consider such models in this paper, but focus on some of the extra matter particles predicted by Calabi–Yau compactification.

Our study is carried out in the phenomenologically attractive [9] context in which the minimal rank-5 gauge group which could be given by Calabi–Yau compactification [4], $G \equiv \text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_Y \times \text{U}(1)_E$, survives down to low energies and there are no light ‘survivors’ coming from $(\underline{27} + \overline{\underline{27}})$ representations of E_6 . However, some qualitative features of our analysis should also be valid for models with different extended low-energy gauge groups [10]. Each chiral supermultiplet transforming according to the $\underline{27}$ representation of E_6 contains five colourless G supermultiplets with an electrically neutral component:

$$H \equiv \begin{pmatrix} H^+ \\ H^0 \end{pmatrix} \sim (1, 2, +1/2, -2/3), \quad (1a)$$

$$\overline{H} \equiv \begin{pmatrix} \overline{H}^0 \\ \overline{H}^- \end{pmatrix}, L \equiv \begin{pmatrix} \nu \\ e \end{pmatrix} \sim (1, 2, -1/2, -1/6), \quad (1b)$$

$$N, \nu^c \sim (1, 1, 0, +5/6). \quad (1c)$$

Treating the superpotential couplings as phenomenological parameters, the model has a discrete ambiguity under the exchange $(\bar{H}, N, D^c) \leftrightarrow (L, \nu^c, d^c)$, where D^c and d^c are colour antitriplets of charge $+1/3$ and (L, ν^c, d^c) are contained in the $\underline{16}$ of $SO(10)$. Assuming that $SU(2)_L$ commutes with lepton number, it is not restrictive to identify L with the lepton doublet ($L = 1$); the requirement that lepton number not be broken by left-handed sneutrino v.e.v.'s then imposes $\langle \tilde{\nu} \rangle = 0$. The doublets giving masses to the ordinary particles through their non-zero vacuum expectation values (v.e.v.'s), as in the supersymmetric Standard Model, must therefore be of type H (1a) and $\bar{H}$ (1b) ($L = 0$). Moreover, to get an acceptable breaking of G at low energies, at least one additional singlet must acquire a v.e.v. of order 400 GeV to 1 TeV [9]. In the framework of radiative breaking [9], which is the only scenario so far available, a non-zero v.e.v. for $\tilde{\nu}^c$ would lead to unacceptable fermion masses and/or lepton number violations. Therefore we assign $L = 0$ to N and $L = -1$ to ν^c , and we stick to the case where only $\langle N \rangle$ can be non-zero. In this paper we examine in detail the $L = 0$ sector, involving the superfields H , $\bar{H}$, and N . The fermionic members of these superfields can mix with the gauginos; we investigate the mass eigenstates of the resulting charged fermion, neutral fermion, charged scalar, and neutral scalar mass matrices, as well as their couplings.

We find that the mass of the lightest charged fermion depends on unknown Yukawa couplings, and could be anywhere above the present experimental limit of 20 GeV. As discussed elsewhere [11], the lightest neutral fermion is expected to be stable, apart from a tiny region of parameter space which gives a left-handed sneutrino as the lightest supersymmetric particle, and must have large gaugino components if its relic cosmological density is not to exceed the critical density. Here we confirm that the other neutral fermions can easily be heavier than this relic. The lightest charged scalar could well be lighter than the $W^\pm$, in contrast with the minimal supersymmetric Standard Model [12]. We identify one neutral scalar which generally weighs between 65 and 150 GeV, and there could well be other neutral scalars weighing less than 100 GeV.

2. GENERAL FRAMEWORK

Each generation of matter particles provided by Calabi-Yau compactification contains $(H^\dagger, \bar{H}^0, N)$ and their supersymmetric partners, and there is at least one additional neutral gauge field beyond those in the Standard Model. Therefore, even in our minimal case with $N_g = 3$ and a rank-5 gauge group, the chargino mass-matrix is 4×4 [for $(\tilde{W}^-, \tilde{\bar{H}}_{1,2,3}^-) \times (\tilde{W}^+, \tilde{H}_{1,2,3}^+)$], the neutralino mass-matrix is 12×12 [for $(\tilde{W}_3, \tilde{B}, \tilde{B}_E, \tilde{H}_{1,2,3}^0, \tilde{\bar{H}}_{1,2,3}^0, \tilde{N}_{1,2,3})^2$], the charged scalar mass-matrix is 6×6 [for $(H_{1,2,3}^+, (\bar{H}^-)_{1,2,3}^*) \times ((H^+)_{1,2,3}^*, \bar{H}_{1,2,3}^-)$], and the neutral scalar mass-matrix is 18×18 [for $(\text{Re } H_{1,2,3}^0, \text{Im } H_{1,2,3}^0, \text{Re } \bar{H}_{1,2,3}^0, \text{Im } \bar{H}_{1,2,3}^0, \text{Re } N_{1,2,3}, \text{Im } N_{1,2,3})^2$]. Short cuts are evidently necessary if we are to extract useful information from these large matrices. We exploit two such short cuts. The most useful is to note [13] that we can always choose to work in a basis for the

matter superfields in which only one of the neutral superfields with each set of quantum numbers has a non-zero v.e.v.^{*)}. These we denote by

$$\langle 0 | H_3^0 | 0 \rangle \equiv v, \quad \langle 0 | \bar{H}_3^0 | 0 \rangle \equiv \bar{v}, \quad \langle 0 | N_3 | 0 \rangle \equiv \infty, \quad (2a)$$

so that

$$\langle 0 | H_{1,2}^0 | 0 \rangle = \langle 0 | \bar{H}_{1,2}^0 | 0 \rangle = \langle 0 | N_{1,2} | 0 \rangle = 0. \quad (2b)$$

In this basis, consideration of the soft supersymmetry-breaking terms in the effective scalar potential shows that the Yukawa coupling constants λ_{ijk} of the H_i , $\bar{H}_j$, and N_k superfields take a particularly simple form. Even if no soft supersymmetry-breaking trilinear interactions of the spin-zero components of the H_i , $\bar{H}_j$, and N_k superfields,

$$V \ni A_{ijk} \lambda_{ijk} H_i \bar{H}_j N_k \quad (3a)$$

are present at the tree level of the effective field theory obtained upon compactification, whilst the gaugino mass $m_{1/2}$ and/or the scalar mass m_0 are non-zero, terms (3a) will be generated from these other soft supersymmetry-breaking terms by radiative corrections. The leading-order renormalization group equations for small λ_{ijk} provide a universal value for the matrix elements A_{ijk} in Eq. (3a), which we write in the form

$$V \ni (\hat{A} m_{1/2}) \lambda_{ijk} H_i \bar{H}_j N_k. \quad (3b)$$

If any two of the three fields in a term of the form (3) have non-zero v.e.v.'s, then the third field will also have a non-zero v.e.v. Recalling that, in our basis, fields with indices 1 and 2 have no v.e.v.'s, we conclude that

$$\lambda_{i33} = \lambda_{3j3} = \lambda_{33k} = 0 \quad \text{for: } i, j, k = 1, 2 \quad (4a)$$

whilst

$$\lambda_{333}, \lambda_{ij3}, \lambda_{i3k}, \lambda_{3jk}, \lambda_{ijk} \neq 0 \quad (4b)$$

in general. Working in this basis greatly simplifies our mixing analysis: the (4×4) chargino matrix becomes $(2 \times 2) + (1 \times 1) + (1 \times 1)$, the (12×12) neutralino matrix becomes $(6 \times 6) + (6 \times 6)$, the (6×6) charged scalar matrix becomes $(2 \times 2) + (2 \times 2) + (2 \times 2)$, and the (18×18) neutral scalar matrix becomes $(6 \times 6) + (12 \times 12)$. Note also that one of the (2×2) submatrices of the

*) We assume that the physical vacuum does not violate charge or colour conservation. A preliminary investigation [9] has shown that this is perfectly consistent with our general theoretical framework. A more detailed investigation might provide interesting constraints on the parameters of our model, but is beyond the scope of this paper.

charged scalar matrix has a massless unphysical Goldstone boson eigenstate, whereas the (6×6) neutral scalar submatrix includes two such Goldstone bosons.

So far we have not made any strong assumptions about the form of the matrix λ_{ijk} of Yukawa couplings. If, as a second short cut, we assume that they are all relatively real, so that CP violation in this sector is negligible, then the (12×12) neutral scalar submatrix further breaks up to become $(6 \times 6) + (6 \times 6)$, whilst the forms of the other matrices are not changed qualitatively. We are now in a position to discuss the different mass matrices.

In doing so, we shall work in the framework in which the only seed of soft supersymmetry breaking in the observable sector at the unification scale is a universal gaugino mass $m_{1/2}$ [14]. Moreover, it is not restrictive to work with a real and positive λ_{333} , so that, since the renormalization group equations drive $\hat{A} < 0$, also the v.e.v.'s v , $\bar{v}$, and x will be all real and positive. When making numerical estimates, we shall restrict the ranges of variation of our parameters according to Table 1, which has been derived [9] taking into account model-independent constraints, including a non-negative mass-squared for the left-handed sneutrinos and an acceptable shift in the Z-boson mass, and specific dynamical calculations of the radiative breaking of the gauge group G around the Fermi scale.

3. CHARGINOS

We start with the *chargino* mass term, which in the basis introduced above becomes

$$\mathcal{L}_c = -(\tilde{W}^- \tilde{H}_3^- \tilde{H}_1^- \tilde{H}_2^-) \begin{pmatrix} M_2 & -g v & 0 & 0 \\ -g \bar{v} & -\lambda_{333} x & 0 & 0 \\ 0 & 0 & -\lambda_{113} x & 0 \\ 0 & 0 & 0 & -\lambda_{223} x \end{pmatrix} \begin{pmatrix} \tilde{W}^+ \\ \tilde{H}_3^+ \\ \tilde{H}_1^+ \\ \tilde{H}_2^+ \end{pmatrix} + h.c. \dots (5)$$

Two eigenvalues associated with (5) are trivial, namely

$$m_{1,2}^\pm = |\lambda_{113,223}| x = |\lambda_{113,223}| \left(\frac{x}{v} \right) \sqrt{\frac{\sqrt{2}}{4 G_F [1 + (\bar{v}/v)^2]}} \quad (6)$$

The condition $m_{1,2}^\pm \geq 20$ GeV from PEP and PETRA [15] corresponds to a region of the $(|\lambda_{113,223}|, x/v)$ plane which depends only slightly on $\bar{v}/v$ in the dynamically preferred range $0.2 \leq \bar{v}/v \leq 0.6$, as seen in Fig. 1a. These lower bounds on $|\lambda_{113,223}|$ are much smaller than typical values of λ_{333} found in previous dynamical calculations [9]. The eigenvalues of the non-trivial (2×2) submatrix in (5) are given by

$$m_{3,4}^\pm = \frac{(m_{LL}^2 + m_{RR}^2) \mp \sqrt{(m_{LL}^2 - m_{RR}^2)^2 + 4 m_{LR}^4}}{2}, \quad (7a)$$

where

$$m_{LL}^2 = M_2^2 + (g\bar{v})^2, \quad m_{RR}^2 = (\lambda_{333}x)^2 + (gv)^2, \quad (7b)$$

$$m_{LR}^2 = -M_2(gv) + (\lambda_{333}x)(g\bar{v}), \quad M_2 = \frac{\alpha_{e.m.} m_{1/2}}{\alpha_3 \sin^2 \theta_W}.$$

We have studied (7), varying λ_{333} , $m_{1/2}$, x/v , and $\bar{v}/v$ over the dynamically preferred ranges shown in Table 1 and taking into account the correlation between $m_{1/2}$ and x/v . The lowest values of $m_{\frac{1}{2}}$ that we found are shown in Fig. 1b. We see that an absolute lower bound is

$$m_{\frac{1}{2}}^{\pm} \gtrsim 43 \text{ GeV}, \quad (8)$$

that there could be as many as three charginos lighter than $m_Z/2$, and that two of them could have masses close to the present e^+e^- experimental lower limit [15].

4. NEUTRALINOS

We now turn to the *neutralino* mass-matrix. As in the case of the charginos (5), the $(\tilde{W}^3, \tilde{B}, \tilde{B}_E, \tilde{H}_3^0, \tilde{\tilde{H}}_3^0, \tilde{N}_3)$ sector decouples from the $(\tilde{H}_{1,2}^0, \tilde{\tilde{H}}_{1,2}^0, \tilde{N}_{1,2})$ sector. The former sector was discussed in Refs. 9 and 11, where it was found that the lightest neutralino χ in that sector most likely had large gaugino components, and a mass

$$m_{\chi} \approx \left(\frac{1}{5} \text{ to } \frac{1}{7} \right) m_{1/2}. \quad (9)$$

Indeed, if the lightest neutralino were predominantly a Higgsino, its annihilation rate in the early Universe would have been so small that its present relic density would be unacceptably high [11]. The second neutralino sector has the mass matrix

$$\mathcal{L}_N = -\frac{1}{2} \begin{pmatrix} \tilde{H}_2^0 & \tilde{\tilde{H}}_2^0 & \tilde{N}_2 & \tilde{H}_1^0 & \tilde{\tilde{H}}_1^0 & \tilde{N}_1 \end{pmatrix} \begin{pmatrix} 0 & \lambda_{223}x & \lambda_{232}\bar{v} & 0 & 0 & \lambda_{231}\bar{v} \\ 0 & \lambda_{322}v & 0 & 0 & 0 & \lambda_{321}v \\ 0 & \lambda_{132}\bar{v} & \lambda_{312}v & 0 & 0 & 0 \\ \hline \text{symmetric} & 0 & \lambda_{113}x & \lambda_{131}\bar{v} & 0 & 0 \\ 0 & \lambda_{311}v & 0 & 0 & 0 & 0 \end{pmatrix} \begin{pmatrix} \tilde{H}_2^0 \\ \tilde{\tilde{H}}_2^0 \\ \tilde{N}_2 \\ \tilde{H}_1^0 \\ \tilde{\tilde{H}}_1^0 \\ \tilde{N}_1 \end{pmatrix} \quad (10)$$

+ hermitian conjugate.

In order to understand its eigenvalues, a convenient limit of (10) to study is when $x/v \rightarrow \infty$ with $\bar{v}/v$ fixed. In this case there are $2 + 2$ heavy Majorana neutralinos with masses $\sim |\lambda_{113,223}|x$ which almost make two neutral Dirac fermions, as well as two light Majorana particles with masses

$$m_{\chi_L} \sim \frac{(\lambda_{i3j}, \lambda_{3ij})^2}{|\lambda_{i33}, \lambda_{j33}|} \frac{v^2}{x} \quad (i, j = 1, 2) \quad (11)$$

As already mentioned, consistency with cosmology requires $m_{\chi_L} < m_\chi$. This is not difficult to arrange, given all the uncertainties in the couplings λ_{ijk} . We do not explore in detail the bounds on the λ_{i3j} , λ_{3ij} , and λ_{i33} that come from requiring $m_\chi < m_{\chi_L}$ (11), as they would be complicated and difficult to interpret. As an example, if we take $\lambda_{i3j} \approx \lambda_{3ij} \approx 0.15$, $\lambda_{i33} \approx 0.05$, and $x/v = 3$, $\bar{v}/v = 0.5$, the estimate (11) yields $m_{\chi_L} \approx 23$ GeV, which is comfortably above the comparable minimum value of m_χ .

5. CHARGED SCALARS

Next we examine the *charged Higgs* mass-matrix, which breaks up into three independent (2×2) submatrices, as already noted. The (2×2) submatrix $\mathcal{M}_3^2$ mixing the ‘true’ charged Higgses $H_3^\pm$ whose neutral partners have v.e.v.’s is

$$\mathcal{L}_+ = - (H_3^{+*}, \bar{H}_3^-) \mathcal{M}_3^2 \begin{pmatrix} H_3^+ \\ \bar{H}_3^{-*} \end{pmatrix}, \quad (12a)$$

where we use the minimization condition on the potential to eliminate the soft supersymmetry-breaking mass terms in favour of the model parameters λ_{333} , A_{333} :

$$\mathcal{M}_3^2 = \begin{pmatrix} -\lambda_{333} A_{333} \frac{x\bar{v}}{v} + \left(\frac{g^2}{2} - \lambda_{333}^2\right) \bar{v}^2 & -\lambda_{333} A_{333} x + \left(\frac{g^2}{2} - \lambda_{333}^2\right) v\bar{v} \\ -\lambda_{333} A_{333} x + \left(\frac{g^2}{2} - \lambda_{333}^2\right) v\bar{v} & -\lambda_{333} A_{333} \frac{xv}{\bar{v}} + \left(\frac{g^2}{2} - \lambda_{333}^2\right) v^2 \end{pmatrix}. \quad (12b)$$

Solving the eigenvalue problem, one finds a zero-mass Goldstone boson which is eaten by the $W^\pm$, and a physical field of mass $m_{H^\pm}$:

$$m_{H^\pm}^2 = m_W^2 + \left[-\lambda_{333} A_{333} x \left(\frac{\bar{v}}{v} + \frac{v}{\bar{v}} \right) - \lambda_{333}^2 (v^2 + \bar{v}^2) \right]. \quad (13)$$

Using the dynamically preferred values of Table 1 we find that the square bracket in (13) is always positive, and indeed

$$\left[\right]_{(13)} \gtrsim 1700 \text{ GeV}^2 \quad (14a)$$

implying

$$m_{H^\pm} \gtrsim 90 \text{ GeV} \quad (14b)$$

for $m_{W^\pm} = 82 \text{ GeV}$.

It is more difficult to analyse the eigenvalues of the *charged unHiggs* mass matrices $\mathcal{M}_{1,2}^2$, i.e. those whose neutral partners do not have v.e.v's. They take the forms

$$\mathcal{M}_i^2 = \begin{pmatrix} \tilde{m}_{H_i}^2 + m_{H_i^\pm}^{D^2} & -\lambda_{ii3} \lambda_{333} v \bar{v} \\ + \lambda_{ii3}^2 x^2 & -\lambda_{ii3} A_{ii3} x \\ \hline -\lambda_{ii3} \lambda_{333} v \bar{v} & \tilde{m}_{\bar{H}_i}^2 + m_{\bar{H}_i^\pm}^{D^2} \\ -\lambda_{ii3} A_{ii3} x & + \lambda_{ii3}^2 x^2 \end{pmatrix}, \quad i = 1, 2 \quad (15)$$

where $\tilde{m}_{H_i, \bar{H}_i}^2$ are the soft supersymmetry-breaking masses that are calculable using the renormalization group equations, and the D-term contributions are

$$m_{H_i^\pm}^{D^2} = -\frac{g^2 - g'^2}{4} (v^2 - \bar{v}^2) - \frac{g'^2}{9} (5x^2 - 4v^2 - \bar{v}^2), \quad (16a)$$

$$m_{\bar{H}_i^\pm}^{D^2} = \frac{g^2 - g'^2}{4} (v^2 - \bar{v}^2) - \frac{g'^2}{36} (5x^2 - 4v^2 - \bar{v}^2). \quad (16b)$$

Since the D-terms (16) are negative for $x > v > \bar{v}$ as expected, it is possible for charged unHiggses to be relatively light, even close to the e^+e^- experimental lower limit of about 20 GeV [15]. A full analysis of this possibility is complicated by the large number of unknown parameters in $\mathcal{M}_i^2$ (15): $m_{1/2}$, x/v , $\bar{v}/v$, λ_{ii3} , λ_{333} , A_{ii3} . We use the parameter values in Table 1, which are valid if $|\lambda_{ii3}| < \lambda_{333}$. We present in Fig. 2 results for the lighter charged unHiggs mass $m_{h_i^\pm}$ over the $(m_{1/2}, x/v)$ plane in the limit $\lambda_{ii3} \rightarrow 0$: larger values of λ_{ii3} give larger values of $m_{h_i^\pm}$. On the right of the diagonal solid line in Fig. 2 which corresponds to

$$m_{1/2} \sim (57 \text{ GeV}) \cdot \frac{x}{v} \quad (17)$$

for the representative choice $A_{ii3} = -0.6m_{1/2}$, $\lambda_{333} = 0.2$, $\bar{v}/v = 0.5$, we find $m_{h_1}^\pm < 20$ GeV in the limit $\lambda_{ii3} \rightarrow 0$. In this domain, whilst $m_{h_1}^\pm$ is generically $O(x)$ for larger values of λ_{ii3} , there is for any value of $(m_{1/2}, x/v)$ a restricted range of λ_{ii3} values which give $20 \text{ GeV} < m_{h_1}^\pm < 100 \text{ GeV}$. On the left of the diagonal solid line in Fig. 2 there is a somewhat larger domain of $(m_{1/2}, x/v)$ space for which $20 \text{ GeV} < m_{h_1}^\pm < 100 \text{ GeV}$ in the limit as $\lambda_{ii3} \rightarrow 0$, and values of λ_{ii3} up to 0.1 give $m_{h_1}^\pm$ only slightly heavier.

6. NEUTRAL SCALARS

We now turn to the neutral Higgs bosons. As mentioned earlier, their mixing matrix breaks into two (3×3) and two (6×6) submatrices in our basis in which only H_3 , $\bar{H}_3$, and N_3 are ‘true’ Higgses with v.e.v.’s, and if one assumes that the λ_{ijk} are real. The mass terms for the real and imaginary components of the neutral true Higgses are:

$$\mathcal{L}_{R,I} = -\frac{1}{2} \begin{pmatrix} H_{R,I}^0 & \bar{H}_{R,I}^0 & N_{R,I} \end{pmatrix} \mathcal{M}_{R,I}^2 \begin{pmatrix} H_{R,I}^0 \\ \bar{H}_{R,I}^0 \\ N_{R,I} \end{pmatrix}, \quad (18)$$

with

$$\mathcal{M}_R^2 = \begin{pmatrix} \begin{array}{c} (g^2 + \frac{25}{9}g'^2)v^2 \\ -2\lambda_{333}A_{333}\frac{\bar{v}x}{v} \end{array} & \begin{array}{c} (4\lambda_{333}^2 - g^2 - \frac{5}{9}g'^2)v\bar{v} \\ +2\lambda_{333}A_{333}x \end{array} & \begin{array}{c} (4\lambda_{333}^2 - \frac{20}{9}g'^2)v x \\ +2\lambda_{333}A_{333}\bar{v} \end{array} \\ \text{SYM} & \begin{array}{c} (g^2 + \frac{10}{9}g'^2)\bar{v}^2 \\ -2\lambda_{333}A_{333}\frac{vx}{\bar{v}} \end{array} & \begin{array}{c} (4\lambda_{333}^2 - \frac{5}{9}g'^2)\bar{v}x \\ +2\lambda_{333}A_{333}v \end{array} \\ \text{MET} & \text{RIC} & \begin{array}{c} \frac{25}{9}g'^2x^2 \\ -2\lambda_{333}A_{333}\frac{v\bar{v}}{x} \end{array} \end{pmatrix} \quad (19)$$

In the limit of large x/v and $m_{1/2}$ the lightest mass eigenstate of $\mathcal{M}_R^2$ (19) is $H_L^0 = (vH_R^0 + \bar{v}\bar{H}_R^0)/\sqrt{v^2 + \bar{v}^2}$. It is easy to check that $m_{H_L^0} = O(v, \bar{v})$, not $O(x)$. This is to be expected since for large $m_{1/2}$ and x the low-energy effective theory reduces to the Standard Model, which must contain at least one real neutral true Higgs boson with mass $O(m_W)$. We find in this limit that

$$m_{H_L^0} \simeq \frac{1}{v^2 + \bar{v}^2} \left\{ (g^2 + g'^2)(v^2 - \bar{v}^2)^2 + \frac{g'^2}{9}(4v^2 + \bar{v}^2)^2 + 8\lambda_{333}^2 v^2 \bar{v}^2 \right\} - \frac{\epsilon^2}{C}, \quad (20)$$

where

$$\epsilon = \frac{1}{\sqrt{v^2 + \bar{v}^2}} \left[\left(4\lambda_{333}^2 - \frac{20}{9}g'^2 \right) v^2 + 4\lambda_{333}A_{333} \frac{v\bar{v}}{x} + \left(4\lambda_{333}^2 - \frac{5}{9}g'^2 \right) \bar{v}^2 \right], \quad (21)$$

$$G = \frac{25}{9} g'^2$$

The dominant term in (20) is the first one, whilst ϵ^2/c makes less than a 10% change in $m_{H_L^0}$ for small $\bar{v}/v$, and a larger one for $\bar{v}/v = 0.6$. Varying only the parameters λ_{333} and A_{333} , one obtains an interval of values of $m_{H_L^0}^2$ which are possible for a given set of values of x/v , $m_{1/2}$, and $\bar{v}/v$. For $x/v \approx 3$ and $m_{1/2} \approx 100$ GeV the approximate expression (20) gives a lower bound for $m_{H_L^0}^2$ which is larger than the more precise numerical calculations by 10% for $\bar{v}/v = 0.2$ and by 30% for $\bar{v}/v = 0.6$, whilst the corresponding upper bound is reproduced with a few percent accuracy. We have verified numerically that (20) and (21) give a good approximation to $m_{H_L^0}$ (i.e. are accurate to within a few percent) when $x/v \geq 4$ and $m_{1/2} \geq 200$ GeV.

In general, we find that

$$65 \text{ GeV} \lesssim m_{H_L^0} \lesssim 150 \text{ GeV} \quad (22)$$

for values of the parameters in the usual ranges. Figure 3a shows contours of $m_{H_L^0}$ in the $(x/v, \bar{v}/v)$ plane for typical values of λ_{333} , A_{333} and $m_{1/2}$. The other two real neutral true Higgses with masses $O(x)$ are too heavy to be detectable at LEP.

Next we turn to the *imaginary neutral true Higgses*, whose mixing matrix is $\mathcal{M}_I^2$:

$$\mathcal{M}_I^2 = (-2\lambda_{333}A_{333}x) \begin{pmatrix} \bar{v}/v & 1 & \bar{v}/x \\ 1 & v/\bar{v} & v/x \\ \bar{v}/x & v/x & v\bar{v}/x^2 \end{pmatrix}. \quad (23)$$

This has two massless eigenstates which are Goldstone bosons, eaten by the massive neutral gauge bosons Z and Z' , and one massive state

$$a = \frac{1}{\sqrt{\frac{\bar{v}}{v} + \frac{v}{\bar{v}} + \frac{v\bar{v}}{x^2}}} \left[\sqrt{\frac{\bar{v}}{v}} H_1^0 + \sqrt{\frac{v}{\bar{v}}} \bar{H}_1^0 + \sqrt{\frac{v\bar{v}}{x^2}} N_1 \right],$$

with

$$m_a^2 = (-2\lambda_{333}A_{333}x) \left(\frac{\bar{v}}{v} + \frac{v}{\bar{v}} + \frac{v\bar{v}}{x^2} \right). \quad (24)$$

This state is also heavy in the limit $x \rightarrow \infty$, but can be interestingly light. Taking the minimum values of λ_{333} , A_{333} , and x/v discussed above, we find

$$m_a \gtrsim 100 \text{ GeV}, \quad (25)$$

which could even be lighter than the real neutral true Higgs (22). Figure 3b shows contours of m_a in the $(x/v, m_{1/2})$ plane for typical values of λ_{333} , A_{333} , and $\bar{v}/v$.

The final bosons to be analysed are the neutral unHiggses, i.e. those in multiplets without v.e.v.'s, starting with the *real neutral unHiggses* ($H_{R,1,2}^0, \bar{H}_{R,1,2}^0, N_{R,1,2}$) whose (6×6) mixing matrix is

$$\mathcal{M}_R^u = \begin{array}{|c|c|c|c|c|c|} \hline \begin{array}{l} \tilde{m}_{H_1}^2 + m_{H^0}^2 \\ + \lambda_{113}^2 x^2 + \\ (\lambda_{131}^2 + \lambda_{132}^2) \bar{v}^2 \end{array} & & & S & Y & M \\ \hline \begin{array}{l} (\lambda_{131} \lambda_{231} + \\ \lambda_{132} \lambda_{232}) \bar{v}^2 \end{array} & \begin{array}{l} \tilde{m}_{H_2}^2 + m_{H^0}^2 \\ + \lambda_{223}^2 x^2 + \\ (\lambda_{231}^2 + \lambda_{232}^2) \bar{v}^2 \end{array} & & M & E & T \\ \hline \begin{array}{l} \lambda_{113} A_{113} x + \\ (\lambda_{113} \lambda_{333} + \lambda_{131} \lambda_{311} \\ + \lambda_{132} \lambda_{322}) v \bar{v} \end{array} & \begin{array}{l} (\lambda_{231} \lambda_{311} + \\ \lambda_{232} \lambda_{312}) v \bar{v} \end{array} & \begin{array}{l} \tilde{m}_{H_1}^2 + m_{H^0}^2 \\ + \lambda_{113}^2 x^2 + \\ (\lambda_{311}^2 + \lambda_{312}^2) v^2 \end{array} & R & I & C \\ \hline \begin{array}{l} (\lambda_{131} \lambda_{321} + \\ \lambda_{132} \lambda_{322}) v \bar{v} \end{array} & \begin{array}{l} \lambda_{223} A_{223} x + \\ (\lambda_{223} \lambda_{333} + \lambda_{231} \lambda_{321} \\ + \lambda_{232} \lambda_{322}) v \bar{v} \end{array} & \begin{array}{l} (\lambda_{311} \lambda_{321} + \\ \lambda_{312} \lambda_{322}) v^2 \end{array} & \begin{array}{l} \tilde{m}_{H_2}^2 + m_{H^0}^2 \\ + \lambda_{223}^2 x^2 + \\ (\lambda_{321}^2 + \lambda_{322}^2) v^2 \end{array} & & \\ \hline \begin{array}{l} \lambda_{131} A_{131} \bar{v} + \\ (\lambda_{131} \lambda_{333} + \\ \lambda_{113} \lambda_{311}) x v \end{array} & \begin{array}{l} \lambda_{231} A_{231} \bar{v} + \\ (\lambda_{231} \lambda_{333} + \\ \lambda_{223} \lambda_{321}) x v \end{array} & \begin{array}{l} \lambda_{311} A_{311} v + \\ (\lambda_{311} \lambda_{333} + \\ \lambda_{113} \lambda_{131}) x \bar{v} \end{array} & \begin{array}{l} \lambda_{321} A_{321} v + \\ (\lambda_{321} \lambda_{333} + \\ \lambda_{223} \lambda_{231}) x \bar{v} \end{array} & \begin{array}{l} \tilde{m}_{N_1}^2 + m_N^2 + \\ (\lambda_{131}^2 + \lambda_{231}^2) \bar{v}^2 + \\ (\lambda_{311}^2 + \lambda_{321}^2) v^2 \end{array} & \\ \hline \begin{array}{l} \lambda_{132} A_{132} \bar{v} + \\ (\lambda_{132} \lambda_{333} + \\ \lambda_{113} \lambda_{312}) x v \end{array} & \begin{array}{l} \lambda_{232} A_{232} \bar{v} + \\ (\lambda_{232} \lambda_{333} + \\ \lambda_{223} \lambda_{322}) x v \end{array} & \begin{array}{l} \lambda_{312} A_{312} v + \\ (\lambda_{312} \lambda_{333} + \\ \lambda_{113} \lambda_{132}) x \bar{v} \end{array} & \begin{array}{l} \lambda_{322} A_{322} v + \\ (\lambda_{322} \lambda_{333} + \\ \lambda_{223} \lambda_{132}) x \bar{v} \end{array} & \begin{array}{l} (\lambda_{131} \lambda_{132} + \\ \lambda_{231} \lambda_{232}) \bar{v}^2 + \\ (\lambda_{311} \lambda_{312} + \\ \lambda_{321} \lambda_{322}) v^2 \end{array} & \begin{array}{l} \tilde{m}_{N_2}^2 + m_N^2 + \\ (\lambda_{132}^2 + \lambda_{232}^2) \bar{v}^2 + \\ (\lambda_{312}^2 + \lambda_{322}^2) v^2 \end{array} \\ \hline \end{array} \quad (26)$$

where

$$\begin{aligned} m_{H^0}^{D^2} &= \frac{g^2 + g'^2}{4} (v^2 - \bar{v}^2) - \frac{g'^2}{9} (5x^2 - 4v^2 - \bar{v}^2) \\ m_{H^0}^{D^2} &= -\frac{g^2 + g'^2}{4} (v^2 - \bar{v}^2) - \frac{g'^2}{36} (5x^2 - 4v^2 - \bar{v}^2) \\ m_N^{D^2} &= \frac{5g'^2}{36} (5x^2 - 4v^2 - \bar{v}^2) \end{aligned} \quad (27)$$

and $\tilde{m}_{H_i, \bar{H}_i, N_i}$ are the soft supersymmetry-breaking masses $O(m_{1/2})$ which are calculable using the renormalization group equations. If we take the limit of (26) as $x, m_{1/2} \rightarrow \infty$, we see that all six eigenstates have masses $O(x)$, as one might expect since none of them have any reason to weigh $O(v, \bar{v})$. It is of course possible to obtain a light-mass eigenstate by playing the positive and negative $O(x^2)$ contributions to (26) off against each other, a game which is easier for smaller values of x and $m_{1/2}$, but this is not a generic feature.

Similar comments apply to the *imaginary neutral unHiggs* mixing matrix $\mathcal{M}_I^{u^2}$:

$$\mathcal{M}_I^{u^2} = \begin{array}{c|c|c|c|c|c} \begin{array}{l} \tilde{m}_{H_1}^2 + m_{H^0}^2 \\ + \lambda_{113}^2 x^2 + \\ (\lambda_{131}^2 + \lambda_{132}^2) \bar{v}^2 \end{array} & & & S & Y & M \\ \hline \begin{array}{l} (\lambda_{131} \lambda_{231} + \\ \lambda_{132} \lambda_{232}) \bar{v}^2 \end{array} & \begin{array}{l} \tilde{m}_{H_2}^2 + m_{H^0}^2 \\ + \lambda_{223}^2 x^2 + \\ (\lambda_{231}^2 + \lambda_{232}^2) \bar{v}^2 \end{array} & & M & E & T \\ \hline \begin{array}{l} -\lambda_{113} A_{113} x + \\ (-\lambda_{113} \lambda_{333} + \lambda_{131} \lambda_{311} \\ + \lambda_{132} \lambda_{312}) v \bar{v} \end{array} & \begin{array}{l} (\lambda_{231} \lambda_{311} + \\ \lambda_{232} \lambda_{312}) v \bar{v} \end{array} & \begin{array}{l} \tilde{m}_{H_4}^2 + m_{H^0}^2 \\ + \lambda_{113}^2 x^2 + \\ (\lambda_{311}^2 + \lambda_{312}^2) v^2 \end{array} & R & I & C \\ \hline \begin{array}{l} (\lambda_{131} \lambda_{321} + \\ \lambda_{132} \lambda_{322}) v \bar{v} \end{array} & \begin{array}{l} -\lambda_{223} A_{223} x + \\ (-\lambda_{223} \lambda_{333} + \lambda_{231} \lambda_{321} \\ + \lambda_{232} \lambda_{322}) v \bar{v} \end{array} & \begin{array}{l} (\lambda_{311} \lambda_{321} + \\ \lambda_{312} \lambda_{322}) v^2 \end{array} & \begin{array}{l} \tilde{m}_{H_2}^2 + m_{H^0}^2 \\ + \lambda_{223}^2 x^2 + \\ (\lambda_{321}^2 + \lambda_{322}^2) v^2 \end{array} & & \\ \hline \begin{array}{l} -\lambda_{131} A_{131} \bar{v} + \\ (-\lambda_{131} \lambda_{333} + \\ \lambda_{113} \lambda_{311}) x v \end{array} & \begin{array}{l} -\lambda_{231} A_{231} \bar{v} + \\ (-\lambda_{231} \lambda_{333} + \\ \lambda_{223} \lambda_{321}) x v \end{array} & \begin{array}{l} -\lambda_{311} A_{311} v + \\ (-\lambda_{311} \lambda_{333} + \\ \lambda_{113} \lambda_{131}) x \bar{v} \end{array} & \begin{array}{l} -\lambda_{321} A_{321} v + \\ (-\lambda_{321} \lambda_{333} + \\ \lambda_{223} \lambda_{231}) x \bar{v} \end{array} & \begin{array}{l} \tilde{m}_{N_1}^2 + m_N^2 + \\ (\lambda_{131}^2 + \lambda_{231}^2) \bar{v}^2 + \\ (\lambda_{311}^2 + \lambda_{321}^2) v^2 \end{array} & \\ \hline \begin{array}{l} -\lambda_{132} A_{132} \bar{v} + \\ (-\lambda_{132} \lambda_{333} + \\ \lambda_{113} \lambda_{312}) x v \end{array} & \begin{array}{l} -\lambda_{232} A_{232} \bar{v} + \\ (-\lambda_{232} \lambda_{333} + \\ \lambda_{223} \lambda_{322}) x v \end{array} & \begin{array}{l} -\lambda_{312} A_{312} v + \\ (-\lambda_{312} \lambda_{333} + \\ \lambda_{113} \lambda_{132}) x \bar{v} \end{array} & \begin{array}{l} -\lambda_{322} A_{322} v + \\ (-\lambda_{322} \lambda_{333} + \\ \lambda_{223} \lambda_{232}) x \bar{v} \end{array} & \begin{array}{l} (\lambda_{131} \lambda_{132} + \\ \lambda_{231} \lambda_{232}) \bar{v}^2 + \\ (\lambda_{311} \lambda_{312} + \\ \lambda_{321} \lambda_{322}) v^2 \end{array} & \begin{array}{l} \tilde{m}_{N_2}^2 + m_N^2 + \\ (\lambda_{132}^2 + \lambda_{232}^2) \bar{v}^2 + \\ (\lambda_{312}^2 + \lambda_{322}^2) v^2 \end{array} \end{array}$$

(28)

which we will not discuss in detail.

7. PRODUCTION IN e^+e^- COLLISIONS

Reviewing our results, we see that there could be as many as three charginos with masses below $m_Z/2$, additional neutralinos not much heavier than the lightest one discussed in Ref. 11, a charged Higgs boson with $m_{H^\pm} \geq 90$ GeV, two charged unHiggses as light as 20 GeV, a real neutral Higgs boson H_L^0 as light as 65 GeV, and an imaginary neutral Higgs boson a as light as 100 GeV. Searches for light charginos and charged Higgs bosons are sufficiently straightforward [16] that we do not discuss them here.

The neutral Higgs boson H_L^0 could be produced at LEP via the reaction $e^+e^- \rightarrow Z^0 + H_L^0$ [16] with a cross-section

$$\frac{\sigma(e^+e^- \rightarrow Z^0 + H_L^0)}{\sigma(e^+e^- \rightarrow Z^0 + H^0)|_{\text{STANDARD MODEL}}} = \left(\frac{g_{H_L^0 ZZ}}{g_{H^0 ZZ}} \right)^2, \quad (29)$$

where

$$\frac{g_{H_L^0 ZZ}}{g_{H^0 ZZ}} = \frac{1}{\sqrt{v^2 + \bar{v}^2}} \left\{ (\alpha v + \beta \bar{v}) \cos^2 \theta_E + \frac{2}{3} \sin \theta_W \sin \theta_E \cos \theta_E \right. \\ \left. (4v\alpha - \bar{v}\beta) + \frac{1}{9} \sin^2 \theta_W \sin^2 \theta_E (25x\gamma + 16v\alpha + \bar{v}\beta) \right\} \quad (30)$$

is the ratio of the $H_L^0 - Z^0 - Z^0$ and the Standard Model $H^0 - Z^0 - Z^0$ coupling constants. In Eq. (30), α, β, γ are the coefficients in the decomposition $H_L^0 = \alpha H^0 + \beta \bar{H}^0 + \gamma N$, and θ_E is the mixing angle of the two neutral weak bosons present in the class of models we are considering. Since the neutral-current phenomenology and the data on the Z^0 mass constrain the mixing angle θ_E to be small ($|\sin \theta_E| \leq 0.1$), one can approximate expression (30) with

$$\frac{g_{H_L^0 ZZ}}{g_{H^0 ZZ}} \simeq \frac{1}{\sqrt{v^2 + \bar{v}^2}} \left\{ \alpha v + \beta \bar{v} + 0.6 \sin^2 \theta_E x\gamma \right\}. \quad (31)$$

For values of $m_{H_L^0}$ close to the lower bound of 65 GeV, the H^0 and $\bar{H}^0$ components in H_L^0 are dominant ($|\alpha|, |\beta| > |\gamma|$), and the ratio of the cross-sections (29) is of the order of unity. For instance, in the case $m_{H_L^0} = 65$ GeV, one has $\bar{v}/v = 0.6$, $x/v = 3$, $\alpha = -0.61$, $\beta = -0.78$, $\gamma = -0.11$, and, correspondingly, $(g_{H_L^0 ZZ}/g_{H^0 ZZ}) \simeq (\alpha v + \beta \bar{v})^2/(v^2 + \bar{v}^2) \simeq 0.8$, giving a large production rate at LEP. On the other hand, the imaginary neutral Higgs boson α has no $Z^0 - Z^0 - \alpha$ coupling at the tree level and so cannot be produced with an appreciable rate in association with the Z^0 boson in e^+e^- annihilation.

If toponium θ is heavy enough, it can decay [17] into both H_L^0 and α with the branching ratios

$$\frac{B(\theta \rightarrow H_L^0 + \gamma)}{B(\theta \rightarrow H^0 + \gamma)|_{\text{STANDARD MODEL}}} = \beta^2 \left(1 + \frac{\bar{v}^2}{v^2} \right) \quad (32)$$

and

$$\frac{B(\theta \rightarrow a + \gamma)}{B(\theta \rightarrow H^0 + \gamma) \Big|_{\text{STANDARD MODEL}}} = \frac{1 + \frac{\bar{v}^2}{v^2}}{1 + \frac{\bar{v}^2}{v^2} + \frac{\bar{v}^2}{x^2}} . \quad (33)$$

Both ratios (32) and (33) can be of order unity when $m_{H_L^0}$ and m_a are close to their lower bounds. In particular, in the cases $m_{H_L^0} = 65$ GeV and $m_a = 100$ GeV, which correspond to $\bar{v}/v = 0.6$, $\bar{v}/x = 0.2$, and $\beta = -0.78$, we have 0.8 and 1.0 for the ratios (32) and (33), respectively.

The variety of superstring phenomenology we have explored here certainly does not exclude plenty of gaugino and Higgs excitement at LEP!

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Table 1

Typical values of the model parameters^{*)}

$0.2 \leq \bar{v}/v \leq 0.6$
$2.8 \leq x/v \leq 10$
$100 \leq m_{1/2} \leq 500$
$0.15 \leq \lambda_{333} \leq 0.35$
$-0.65 \leq A_{333}/m_{1/2} \leq -0.35$

*) The renormalized parameters λ_{333} and A_{333} are specified at the electroweak symmetry-breaking scale and have been derived in the approximation $|\lambda_{113}|, |\lambda_{223}| < \lambda_{333}$.

Figure captions

- Fig. 1** a) Contours in the $(|\lambda_{113,223}|, x/v)$ -plane corresponding to the constraint $m_{\tilde{H}_{1,2}}^{\pm} \geq 20$ GeV, for the cases $\bar{v}/v = 0.2$ and $\bar{v}/v = 0.6$;
- b) The lightest attainable value of the chargino mass $m_3^{\pm}$ as a function of $m_{1/2}$, for some discrete values of $\bar{v}/v$, with x/v and λ_{333} both free to vary within the ranges given in Table 1. The input values $m_W = 82$ GeV, $\sin^2 \theta_W = 0.222$, $\alpha_{em}(m_W) = 1/128$ and $\alpha_3(m_W) = 0.11$ are used in calculating this and the following figures.
- Fig. 2** Dependence of the lighter charged unhiggs mass $m_{\tilde{H}_1}^{\pm}$ on $m_{1/2}$ and x/v , for the representative values $\bar{v}/v = 0.5$, $\lambda_{333} = 0.2$ and $A_{333}/m_{1/2} = -0.6$. The region to the right of the diagonal lines corresponds to $m_{\tilde{H}_1}^{\pm} < 20$ GeV in the limit $|\lambda_{ii3}| \rightarrow 0$. Lines corresponding to $m_{\tilde{H}_1}^{\pm} = 20, 60$ and 100 GeV in the limit $|\lambda_{ii3}| \rightarrow 0$ are shown. The region to the left of these lines corresponds to $m_{\tilde{H}_1}^{\pm} > 100$ GeV in the limit $|\lambda_{ii3}| \rightarrow 0$.
- Fig. 3** a) Contours of $m_{H_L^0}$ in the $(x/v, \bar{v}/v)$ -plane corresponding to the representative values $m_{1/2} = 250$ GeV, $\lambda_{333} = 0.15$ and $A_{333}/m_{1/2} = -0.35$; the dependence on $m_{1/2}$ is very slight.
- b) Contours of m_a in the $(x/v, m_{1/2})$ -plane corresponding to the representative values $\bar{v}/v = 0.6$, $\lambda_{333} = 0.15$ and $A_{333}/m_{1/2} = -0.35$; m_a increases by ~ 1.5 if $\bar{v}/v$ is decreased to 0.2.

Fig. 1

Fig. 2

Fig. 3