A Low Energy Dynamical SUSY Breaking Scenario
Motivated from Superstring Derived Unification

Alon E. Faraggi$^{1,2}$

$^1$ Institute for Fundamental Theory, Department of Physics, University of Florida, Gainesville, FL 32611, USA

$^2$ CERN, Theory Division, 1211 Geneva, Switzerland

Abstract

Recently there has been a resurgence of interest in gauge mediated dynamical supersymmetry breaking scenarios. I investigate how low energy dynamical SUSY breaking may arise from superstring models. In a three generation string derived model I propose that the unbroken hidden non–Abelian gauge group at the string scale is $SU(3)_H$ with matter multiplets. Due to the small gauge content of the hidden gauge group the supersymmetry breaking scale may be consistent with the dynamical SUSY breaking scenarios. The messenger states are obtained in the superstring model from sectors which arise due to the “Wilson–line” breaking of the unifying non–Abelian gauge symmetry. An important property of the string motivated messenger states is the absence of superpotential terms with the Standard Model states. The stringy symmetries therefore forbid the flavor changing processes which may arise due to couplings between the messenger sector states and the Standard Model states. Motivated from the problem of string gauge coupling unification I contemplate a scenario in which the messenger sector consists solely of color triplets. This hypothesis predicts a chargino mass below the $W$–boson mass. Imposing the current limits from the LEP1 and LEP1.5 experiments the lightest supersymmetric particles predicted by this hypothesis are in the mass ranges $M_{\tilde{\chi}^\pm} \approx 55 - 65$ GeV, $M_{\tilde{\chi}^0} \approx 35 - 50$ GeV and $M_{\tilde{\nu}} \approx 45 - 60$ GeV which will be tested in the forthcoming LEP2 experiments.

* E-mail address: faraggi@phys.ufl.edu
†Permanent address.
Recently there has been renewed interest in the dynamical SUSY breaking scenarios [1, 2]. In these scenarios supersymmetry breaking is generated dynamically at a relatively low scale and is transmitted to the observable sector by the gauge interactions of the Standard Model. An important property of this type of gauge mediated supersymmetry breaking is the natural suppression of flavor changing neutral currents. In these scenarios the universality of the Standard Model gauge interactions results in generation blind mass parameters for the supersymmetric scalar spectrum. This attractive property of gauge mediated supersymmetry breaking is an important advantage over some other possible scenarios. A crucial assumption in this regard is the absence of interaction terms between the messenger sector states and the Standard Model states. In this paper I show that the absence of such interaction terms arises naturally in string derived models.

In the dynamical gauge mediated SUSY breaking scenarios, supersymmetry is broken nonperturbatively and the breaking is mediated to the observable sector by a messenger sector. The messenger sector typically consists of vector–like color triplets and electroweak doublets, beyond the spectrum of the Minimal Supersymmetric Standard Model (MSSM). The messenger sector states and their charges under the Standard Model gauge group then determine the sparticle mass spectrum.

The gaugino masses are obtained from one–loop diagrams and are given by

$$M_i(\Lambda) = \frac{\alpha_i(\Lambda)}{4\pi \Lambda}$$

where $\Lambda$ is the SUSY breaking scale and $\alpha_i(\Lambda)$ are the Standard Model coupling constants at the scale $\Lambda$. The scalar masses arise from two–loop diagrams and are given by

$$m^2(\Lambda) = 2\Lambda^2 \left\{ C_3 \left[ \frac{\alpha_3(\Lambda)}{4\pi} \right]^2 + C_2 \left[ \frac{\alpha_2(\Lambda)}{4\pi} \right]^2 + \frac{3}{5} \left( \frac{Y}{2} \right)^2 \left[ \frac{\alpha_1(\Lambda)}{4\pi} \right]^2 \right\}$$

where the weak hypercharge has the standard $SO(10)$ normalization $U(1)_Y = 3/5 U(1)_1$ and $C_3 = 4/3$ for color triplet scalars and zero for sleptons and $C_2 = 3/4$ for electroweak doublets and zero for singlets.

In this paper I examine how a low–energy gauge–mediated dynamical SUSY breaking scenario may arise from superstring derived models. Traditionally it has been assumed that the supersymmetry breaking in superstring models is generated dynamically by a hidden gauge group with a large gauge content, typically $E_8$, $SO(10)$ or $SU(5)$ [3]. I propose that alternative scenarios exist in which the hidden gauge group is broken at the Planck scale to a group with a small gauge content, like $SU(3)$. The appearance of a “small” hidden gauge group may result in the nonperturbative SUSY breaking dynamics at a hierarchically low scale. This illustrates that the SUSY breaking dynamics may indeed be generated at a relatively low scale in accordance with the gauge mediated SUSY breaking scenarios. In this paper the SUSY breaking
sector will not be investigated in detail. Rather, following ref. [4], some heuristic arguments are given which suggest that supersymmetry may indeed be broken in the hidden sector. Instead, the focus of this paper is on a predictive string-motivated scenario with regard to the messenger sector. In this scenario the messenger sector consists solely of color triplets. There are no electroweak doublets in the messenger sector. This scenario differs from previously studied dynamical SUSY breaking scenarios in the context of unified SUSY models and results in specific predictions with regard to the supersymmetric spectrum. Specifically, in the simplest scenario the lightest chargino is predicted to be below the $W$-boson mass, the lightest neutralino is predicted to be of the order $35 - 50$ GeV and the lightest scalar superpartner is the sneutrino with a mass of the order $45 - 60$ GeV.

The messenger sector states are obtained in the superstring model from sectors which arise due to the “Wilson–line” breaking of the unifying non-Abelian gauge symmetry. A consequence of the gauge symmetry breaking in superstring models by “Wilson–line” is the appearance of massless states which do not fit into multiplets of the original unbroken gauge symmetry. I refer to such states generically as exotic “Wilsonian” matter states. This is an important property as it may result in conserved quantum numbers which forbid the couplings of the exotic “Wilsonian” matter states to the Standard Model states. This is precisely what happens in the superstring derived model under consideration. The Standard Model states are obtained from three 16 multiplets of $SO(10)$ and the messenger sector states are “Wilsonian” matter states. In this string derived model it has been shown, to all orders of non-renormalizable terms, that there are no superpotential terms between the “Wilsonian” color triplets and the Standard Model states. This result illustrates that string models can indeed produce the symmetries needed to prevent the interaction of the messenger sector states with the Standard Model states. This demonstrates that, indeed, string motivated gauge mediated SUSY breaking scenarios can resolve the problematic supersymmetric contributions to flavor changing neutral currents.

The motivation to consider a messenger sector with color triplets only arises from the problem of string gauge coupling unification. While in the Minimal Supersymmetric Standard Model (MSSM) the gauge coupling are observed to intersect at a scale of the order of $2 \times 10^{16}$ GeV [7], string theory predicts that the unification scale is of the order of $g_{\text{string}} \times 5 \times 10^{17}$ GeV [8], with $g_{\text{string}} \approx 0.8$ at the unification scale. Thus, approximately a factor of twenty separates the MSSM and string unification scales. It would seem that this discrepancy should have many possible resolutions, keeping in mind that the gauge parameters are extrapolated over fifteen orders of magnitude. Indeed, in string models there are many possible sources that may affect the gauge coupling unification. For instance, heavy string threshold corrections, light SUSY threshold corrections, enhanced gauge group structure at an intermediate mass scale, weak hypercharge normalization which differs from the standard GUT normalization, intermediate matter thresholds and nonperturbative effects. Surprisingly, however, the problem is not easily resolved. In ref. [9] the string gauge coupling
problem was analysed in the context of the realistic free fermionic models. It was shown, in a wide range of realistic free fermionic models, that heavy string threshold correction, non–standard hypercharge normalizations, light SUSY thresholds or intermediate gauge structure do not resolve the problem. Instead, the problem may only be resolved due to the existence of additional intermediate matter thresholds, beyond the MSSM \cite{10, 11, 9}. This additional matter consists of color triplets and electroweak doublets, in vector–like representations. Remarkably, some string models contain in their massless spectrum the additional states, with the specific weak hypercharge assignment needed to achieve string scale unification \cite{10}. In ref. \cite{10, 9} it was shown that there exist many possible scenarios for the mass scales of the additional color triplets and the electroweak doublets. In general, to accommodate the string unification scale with the experimental values for $\alpha_{\text{strong}}(M_Z)$ and $\sin^2 \theta_W(M_Z)$ the additional vector–like color triplets have to be much lighter than the additional electroweak doublets \cite{9}. Furthermore, the mass scale of the color triplets which is required for them to play the role of the messenger sector in the dynamical SUSY breaking scenarios, can also be compatible with the mass scale which is needed to resolve the string gauge coupling unification problem. It was recently also suggested that massive color triplets at the required mass scale, $\Lambda \approx 100$ TeV, are also good dark matter candidates \cite{8}. Thus, the same “Wilsonian” matter states that can provide the missing dark matter, can also play the role of the messenger sector states in the superstring motivated gauge mediated dynamical SUSY breaking scenario. It is important to note that the superstring symmetries which forbid the interaction of the “Wilsonian” matter states with the Standard Model states, insure both their stability as well as the absence of flavor mediating interactions from the messenger sector.

In this paper I propose in a specific superstring derived standard–like model that a $SU(3)_H$ gauge group may be the only non–Abelian part of the hidden gauge group which is left unbroken at the string scale. The hidden $SU(3)_H$ gauge group becomes strongly interacting at a relatively low scale. Gaugino and matter condensation may then drive a non–vanishing F–terms for one of the gauge singlet fields in the massless string spectrum. This singlet couples to the messenger sector. In this paper motivated from the problem of string–scale gauge coupling unification I make the hypothesis that the messenger sector consists only of color triplets while the electroweak doublets are much heavier. The hypothesis that the messenger sector consists solely of color triplets has crucial implications. It predicts the existence of a chargino which is lighter than the $W$–boson mass. Thus, this hypothesis will be confirmed or ruled out in the forthcoming LEP2 experiments.

Examples of semi–realistic superstring models were constructed in the orbifold and free fermionic formulations \cite{12, 14, 15, 16, 10, 17}. The proposed scenario for superstring dynamical SUSY breaking is illustrated in the superstring standard–like model of ref. \cite{14}. The superstring derived standard–like models are constructed in the free fermionic formulation \cite{13}. The realistic free fermionic models \cite{14, 13, 16, 10}. 

4
are defined in terms of a set of boundary condition basis vectors for all the world-sheet fermions, and the one-loop GSO amplitudes. The physical spectrum is obtained by applying the generalised GSO projections. The first five basis vectors in the models that I consider consist of the NAHE set, \( \{1, S, b_1, b_2, b_3\} \). At the level of the NAHE set the observable gauge group is \( SO(10) \times SO(6)^3 \) and the hidden gauge group is \( E_8 \).

The sectors \( b_1, b_2 \) and \( b_3 \) produce 48 generations in the 16 representation of \( SO(10) \). Adding to the NAHE set three additional boundary condition basis vectors, \( \{\alpha, \beta, \gamma\} \), reduces the number of generations to three, one from each of the sectors \( b_1, b_2 \) and \( b_3 \). At the same time the observable \( SO(10) \) gauge group is broken to one of its subgroups, the flavor \( SO(6)^3 \) symmetries are broken to product of \( U(1)'s \), and the hidden \( E_8 \) is broken to one of its subgroups. It is important to note the correspondence between free fermionic models and orbifold models. The free fermionic models correspond to \( Z_2 \times Z_2 \) orbifold models with nontrivial background fields \([\mathbb{S}]\). The Neveu–Schwarz sector corresponds the untwisted sector, and the sectors \( b_1, b_2 \) and \( b_3 \) correspond to the three twisted sectors of the \( Z_2 \times Z_2 \) orbifold models. The three sectors which break the \( SO(10) \) symmetry correspond to Wilson lines in the orbifold terminology.

In the superstring derived standard-like models the Neveu–Schwarz sector gives rise to the generators of the \( SU(3)_C \times SU(2)_L \times U(1)_B \times U(1)_{T_{3R}} \times U(1)^6 \) gauge group in the observable sector. In the hidden sector the Neveu–Schwarz sector produces the generators of the \( SU(5) \times SU(3) \times U(1)^2 \) hidden gauge group. The three sectors \( b_1, b_2 \) and \( b_3 \) produce three chiral 16 of \( SO(10) \) decomposed under \( SU(3) \times SU(2) \times U(1)^2 \). In addition to the spin one and two multiplets, the Neveu–Schwarz (NS) sector produces three pairs of electroweak doublets, \( \{h_1, h_2, h_3, \bar{h}_1, \bar{h}_2, \bar{h}_3\} \), three pairs of \( SO(10) \) singlets with \( U(1) \) charges, \( \{\Phi_{12}, \Phi_{23}, \Phi_{13}, \bar{\Phi}_{12}, \bar{\Phi}_{23}, \bar{\Phi}_{13}\} \), and three singlets of the entire four dimensional gauge group, \( \{\xi_1, \xi_2, \xi_3\} \). The sector \( b_1 + b_2 + \alpha + \beta \) produces one or two additional Higgs pairs, \( \{h_{45}, \bar{h}_{45}, h'_{45}, \bar{h}'_{45}\} \), and several \( SO(10) \) singlet fields with horizontal \( U(1) \) charges, \( \{\Phi_{45}, \bar{\Phi}_{45}, \Phi'_{45}, \bar{\Phi}'_{45}, \Phi_{12}, \bar{\Phi}_{12}\} \). The three sectors \( b_j + 2\gamma \) produce massless states in the vector 16 representation of the \( SO(16) \) subgroup of the hidden \( E_8 \), decomposed under the final hidden gauge group, \( \{T_{1,2,3}, \bar{T}_{1,2,3}, V_{1,2,3}, \bar{V}_{1,2,3}\} \). The \( T_i (\bar{T}_i) \) are 5 (5) and the \( V_i (\bar{V}_i) \) are 3 (3) of the hidden \( SU(5) \) and \( SU(3) \) gauge groups, respectively. In addition, vectors that are combinations of the NAHE set basis vectors and of the basis vectors \( \{\alpha, \beta, \gamma\} \), produce additional massless sectors which break the \( SO(10) \) symmetry explicitly. These are the sectors which arise due to the “Wilson–line” breaking of the \( SO(10) \) gauge symmetry. In the model of ref. \([10]\) there are two pairs of additional vector–like \( SU(3)_C \) color triplets \( \{D_1, D_2, \bar{D}_1, \bar{D}_2\} \) from the sectors \( b_{1,2} + b_3 + \beta + \gamma \). These two color triplet pairs have the standard down–type weak hypercharge assignment \([\mathbb{I}]\). However, they carry non–standard \( SO(10) \) charges under the \( U(1)_{Z'} \) which is embedded in \( SO(10) \) and is orthogonal to the weak hypercharge \([\mathbb{I}]\). In the model of ref. \([11]\), because of the exotic charges of these color triplets under the \( U(1)_{Z'} \) symmetry, there is a residual discrete symmetry which forbids the couplings of these color triplets to the Standard Model states \([\mathbb{I}]\). This is a crucial observation from which follows that the interaction of these color triplets
with the Standard Model states is indeed only through the gauge interactions. The stringy local discrete symmetries therefore forbid the couplings which may result in flavor changing processes. An additional pair of color triplets, \( \{D_3, \bar{D}_3\} \) with a non-standard weak hypercharge assignment is obtained from the sector \( 1 + \alpha + 2 \gamma \). Three pairs of additional electroweak doublets are obtained from the sectors \( \{1 + b_1 + \alpha + 2 \gamma\} \). The sectors \( b_{1,2} + b_3 + \beta \pm \gamma \) also produce two additional pairs of triplets of the hidden \( SU(3)_H \) gauge group. The total number of triplets of the hidden \( SU(3)_H \) gauge group in the model of ref. \[10\] is ten. Further details on the construction of the superstring standard-like models and their spectrum are given in ref. \[13, 16, 6\].

The cubic level and higher order nonrenormalizable terms in the superpotential are obtained by calculating correlators between vertex operators, \( A_N \sim \langle V_1^f V_2^f V_3^b \cdots V_N^b \rangle \), where \( V_i^f (V_i^b) \) are the fermionic (scalar) components of the vertex operators. The non-vanishing terms must be invariant under all the symmetries of the string models and must satisfy all the string selection rules \[19\]. A detailed analysis of texture of fermion mass matrices was done in refs. \[20, 21\] for the model of ref. \[15\]. The analysis was done up to nonrenormalizable terms of order \( N = 8 \). From this analysis it was found that the two sectors \( b_1 \) and \( b_2 \) produce the two heavy generations while the sector \( b_3 \) produces the lightest generation. The mixing terms between the generations are obtained by exchanging states which transform under the 16 vector representation of the hidden \( SO(16) \) gauge group. For example, the lowest order mixing terms in the model of ref. \[13\] arise at order \( N = 6 \),

\[
\begin{align*}
d_3Q_2 h_{45} \Phi_{45} V_3 \bar{V}_2 & \quad d_2Q_3 h_{45} \Phi_{45} V_2 \bar{V}_3, \\
d_3Q_1 h_{45} \Phi_{45} V_1 V_3 & \quad d_1Q_3 h_{45} \Phi_{45} V_1 \bar{V}_3.
\end{align*}
\]

where \( V_i, \bar{V}_i \) transform as 3 and \( \bar{3} \) of the hidden \( SU(3)_H \) gauge group. Thus, in this model in order to obtain a phenomenologically acceptable Cabbibo angle the hidden \( SU(3)_H \) gauge group has to be broken \[21\]. In that case the nonperturbative dynamics in the hidden \( SU(5)_H \) gauge group generate gaugino and matter condensation that may break supersymmetry with \( m_{3/2} \sim 1 \text{ TeV} \) \[4\]. Thus, the supersymmetry breaking scenario in this model is similar to supersymmetry breaking in the traditional supergravity models.

The structure of the model of ref. \[10\] is similar to the structure of the model of ref. \[15\]. In particular, the observable spectrum from the NS sector and the sectors \( b_1 \), \( b_2 \) and \( b_3 \) is identical in the two models with some differences in the charges under the horizontal \( U(1) \) symmetries, \( U(1)_{L,R_{4,5,6}} \). The model of ref. \[10\] produces similarly potential mass terms for the two heavy generations from quartic and quintic order terms. A detailed analysis of the texture of fermion mass matrices in not the purpose of the present paper. For our purposes here, it is sufficient to note that in the model or ref. \[10\] the generation mixing terms are obtained at order \( N = 6 \)

\[
\begin{align*}
d_3Q_2 h_{45}' \Phi_{45} T_2 \bar{T}_3 & \quad d_3Q_1 h_{45}' \Phi_{45} T_1 \bar{T}_3, \\
d_2Q_3 h_{45}' \Phi_{45} \bar{T}_2 T_3 & \quad d_1Q_3 h_{45}' \Phi_{45} \bar{T}_1 T_3.
\end{align*}
\]
where $T_i$ and $\bar{T}_i$ transform as 5 and $\bar{5}$ of the hidden $SU(5)_H$ gauge group. Thus, in this model in order to generate a Cabbibo angle of a phenomenologically acceptable order of magnitude, the hidden $SU(5)$ gauge group has to be broken while the hidden $SU(3)_H$ gauge group is left unbroken. The hidden $SU(5)_H$ gauge group can of course be broken in several possible patterns. I will assume that there exist a solution in which it is completely broken. In this case the SUSY breaking dynamics will be driven by the hidden $SU(3)_H$ gauge group.

Thus, in the model of ref. [10], to obtain a sizable generation mixing the hidden gauge group needs to be broken, and only $SU(5)_H$ remains unbroken. In this case because of the small gauge content of the $SU(3)$ gauge group, the scale where the hidden $SU(3)_H$ gauge group becomes strongly interacting can be much smaller, in accordance with the gauge mediated dynamical SUSY breaking scenarios. The scale at which the hidden $SU(3)_H$ gauge group becomes strongly interacting is given by

$$\Lambda_3 = M_{\text{string}} \exp\left(\frac{2\pi}{b} \frac{(1 - \alpha_0)}{\alpha_0}\right), \quad (5)$$

where $M_{\text{string}}$ is the string unification scale, $b = 1/2 \ n_3 - 9$ is the $\beta$-function coefficient of the hidden $SU(3)_H$ gauge group and $\alpha_0$ is the gauge coupling constant at the string unification scale. The scale at which the hidden $SU(3)$ gauge group becomes strongly interacting depends on the $M_{\text{string}}$, $\alpha_0$ and on the number of hidden $SU(3)_H$ triplets which are massless at the string scale. For example, with $M_{\text{string}} = 4 \times 10^{17}$ GeV, $\alpha_0 = 1/24$ and $n_3 = 8$, gives $\Lambda_3 \approx 100$ TeV. This is roughly the scale required in the dynamical SUSY breaking scenarios to obtain phenomenologically viable gaugino masses. This illustrates that dynamical low energy SUSY breaking may indeed be generated from superstring derived models. Detailed scenarios for the mass scales of the hidden $SU(3)_H$ matter states can be studied from an analysis of nonrenormalizable terms in the superpotential.

The superstring model under consideration contains in its massless spectrum two pairs of color triplets $\{D_1, D_2, \bar{D}_1, \bar{D}_2\}$ from the sectors $h_{1,2} + b_3 + \beta \pm \gamma$. with the charge assignment $(3,1)_{1/3}$, $(3,1)_{-1/3}$, and one pair, $\{D_3, \bar{D}_3\}$ from the sector $1 + \alpha + 2\gamma$ with charges $(3,1)_{1/6}$, $(3,1)_{-1/6}$, under $SU(3)_{3} \times SU(2)_{L} \times U(1)_{Y}$. The first two pairs transform as regular down-type quarks under the Standard Model gauge group. The cubic level superpotential in this model is given by

$$W_3 = \{(u_{L_1}^c Q_1 h_1 + N_{L_1}^c L_1 h_1 + u_{L_2}^c Q_2 h_2 + N_{L_2}^c L_2 h_2 + u_{L_3}^c Q_3 h_3 + N_{L_3}^c L_3 h_3)
+ h_1 \bar{h}_2 \Phi_{12} + h_1 \bar{h}_3 \Phi_{13} + h_2 \bar{h}_3 \Phi_{23} + \bar{h}_1 h_2 \Phi_{12} + \bar{h}_1 h_3 \Phi_{13} + \bar{h}_2 h_3 \Phi_{23}
+ \Phi_{23} \Phi_{13} \Phi_{12} + \Phi_{23} \Phi_{13} \bar{\Phi}_{12} + \Phi_{12} (\Phi_1 \Phi_1 + \Phi_2 \Phi_2) + \Phi_{12} (\Phi_1 \Phi_1 + \Phi_2 \Phi_2)
+ \frac{1}{2} \xi_3 (\Phi_{45} \bar{\Phi}_{45} + h_{45} \bar{h}_{45} + \Phi_{45} \bar{\Phi}_{45} + h_{45} \bar{h}_{45} + \bar{\Phi}_{45} \Phi_{45} + h_{45} h_{45} + \bar{\Phi}_{45} \Phi_{45})
+ h_{45} \bar{h}_{45} \Phi_{45} + \bar{h}_{45} h_{45} \Phi_{45} + h_{45} \bar{h}_{45} \Phi_{45} + \bar{h}_{45} h_{45} \Phi_{45})
+ \frac{1}{2} (\xi_1 D_1 \bar{D}_1 + \xi_2 D_2 \bar{D}_2) + \frac{1}{\sqrt{2}} (D_1 \bar{D}_2 \phi_2 + \bar{D}_1 D_2 \bar{\phi}_1)\} \quad (6)$$
The color triplet pairs \{D_1, D_2, \bar{D}_1, \bar{D}_2\} can serve as the messenger sector states in the superstring model and the terms in the cubic level superpotential \(\frac{1}{2}(\xi_1 D_1 \bar{D}_1 + \xi_2 D_2 \bar{D}_2) + \sqrt{2}(D_1 \bar{D}_2 \phi_2 + \bar{D}_1 D_2 \phi_1)\) can serve as the coupling between the messenger sector and the SUSY breaking sector, provided that a non–vanishing \(F\)–term is generated in the \(\xi_{1,2}\) or \(\phi_{1,2}\) directions. A detailed analysis of the matter and gaugino condensates was performed in ref. [4] for the model of ref. [15]. There indeed it was argued that a non–vanishing \(F\)–term in the \(\xi_i\) direction is generated due to the hidden matter condensates. There, [4], it was argued that the flat \(F\) directions of the cubic level superpotential are lifted once the nonrenormalizable terms which include the hidden sector matter condensates are included in the analysis. Again it should be emphasised that the analysis was carried out in detail for the model of ref. [15] where the hidden \(SU(5)_H\) is left unbroken at the string scale. However, due to the similar structure of the spectrum of the two models and in particular the almost identical spectrum from the NS sector and the sectors \(b_j + 2\gamma\), similar results are also expected to hold in the case of the model of ref. [10]. Here I am assuming that indeed such a non–vanishing \(F\)–term is generated by nonperturbative effects in the \(\xi\) direction and a more detailed analysis is left for future work. With this assumption the color triplet pairs \{\(D_1, D_2, \bar{D}_1, \bar{D}_2\)\} serve as the messenger sector. The messenger sector therefore consists solely of color triplets. This implies specific predictions for the supersymmetric mass spectrum which are studied below.

The gaugino and Higgsino mass spectrum is obtained by diagonalizing the chargino and neutralino mass matrices. The chargino mass matrix is given by

\[
M_\tilde{C} = \begin{pmatrix}
\tilde{M}_2 & M_W \sqrt{2} \sin \beta \\
M_W \sqrt{2} \cos \beta & \mu
\end{pmatrix},
\]

and the neutralino mass matrix is given by

\[
M_\tilde{N} = \begin{pmatrix}
\tilde{M}_1 & 0 & -M_Z \sin \theta_W \cos \beta & M_Z \sin \theta_W \sin \beta \\
0 & \tilde{M}_2 & M_Z \cos \theta_W \sin \beta & -M_Z \cos \theta_W \cos \beta \\
-M_Z \sin \theta_W \cos \beta & M_Z \cos \theta_W \sin \beta & 0 & \mu \\
M_Z \sin \theta_W \sin \beta & -M_Z \cos \theta_W \cos \beta & \mu & 0
\end{pmatrix},
\]

With the assumption that the messenger sector consists only of color triplets, it follows from Eq. [4] that

\[
\tilde{M}_2 = 0,
\]

in the chargino and neutralino mass matrices, Eqs. [7] and [8], respectively. It follows from this hypothesis that the lightest chargino is lighter than the \(W\)–boson mass. This is therefore a precise prediction of the hypothesis that the messenger sector consists solely of color triplets, which is motivated from the string gauge coupling unification. This should be compared with the gauge mediated SUSY breaking scenarios in the context of the MSSM. There one introduces a messenger sector which consists of color triplets and electroweak doublets in order not to spoil the intersection of the gauge couplings at the MSSM unification scale.
To study the possible spectrum of the lightest superparticles the two parameters in Eqs. 8 and 9, \( \tan \beta \) and \( \mu \) are varied in the ranges \( 1 \leq \tan \beta \leq 3 \) and \(-50 \text{ GeV} \leq \mu \leq 50 \text{ GeV} \) and with \( \Lambda = 100 \text{ TeV} \). The current limits on the chargino and neutralino masses from the LEP1 and LEP1.5 experiments are imposed in the analysis.

Using the LEP1 and LEP1.5 data the four LEP experiments have imposed strong constraints on the existence of charginos and neutralinos below the \( W \)–boson mass \[22, 23\]. These limits are often obtained by making various assumptions on GUT boundary conditions for the gaugino masses and on the scalar mass spectrum. In particular if the sneutrino mass is assumed to be larger than 200 GeV then the chargino mass is constrained to be larger than \( \approx 65 \text{ GeV} \). However, if the assumption on the sneutrino mass is relaxed then destructive interference between the \( s \)–channel exchange diagram of the \( Z \) and \( \gamma \) gauge bosons and the \( t \)–channel exchange diagram of the sneutrino allows smaller chargino masses. A recent analysis was done by the DELPHI collaboration \[23\] from which one can infer the conservative limits of \( m_{\chi^{\pm}} \geq 56 \text{ GeV} \) and \( m_{\chi^0} \geq 35 \text{ GeV} \). The light sneutrino region is precisely the scenario predicted by the superstring motivated dynamical SUSY breaking scenario with a messenger sector which consists only of color triplets. As can be seen from Eq. 9 in this case the second term in Eq. 9 is equal to zero and the only contribution to the sneutrino mass arises from the term due to the weak hypercharge. In the analysis of the sparticle masses I have included the \( D \)–term contribution given by

\[
d_{\tilde{\rho}} = 2 \left( T^3_{\tilde{\rho}L} - \frac{3}{5} Y_{\tilde{\rho}} \tan^2 \theta_W \right) \cos 2\beta M_W^2,
\]

where \( T^3_{\tilde{\rho}L} \) and \( Y_{\tilde{\rho}} \) are the \( SU(2)_L \) and \( U(1)_Y \) charges of a sparticle. In fig. 1 the predicted sneutrino mass is plotted versus the electroweak VEVs ratio \( \tan \beta \) and the scale \( \Lambda \) is taken at 100 TeV.

In fig. 2, the lightest chargino mass is plotted versus the lightest neutralino mass, where the constraints \( M_{\chi^0} \geq 35 \text{ GeV} \) and \( M_{\chi^\pm} \geq 56 \text{ GeV} \) have been imposed. The predicted lightest neutralino mass is approximately in the range \( 35 \text{–} 50 \text{ GeV} \) and the lightest chargino mass is in the approximate range \( 56 \text{–} 65 \text{ GeV} \). The next lightest neutralino, \( \chi^1 \), is found to be heavier than 53 GeV with \( M_{\chi^0} + M_{\chi^1} > 95 \text{ GeV} \) over the entire parameter space. The next lightest and heaviest neutralinos, \( \chi^2 \) and \( \chi^3 \) are found with \( M_{\chi^2} > 84 \text{ GeV} \) and \( M_{\chi^3} > 165 \text{ GeV} \) and the heavy chargino is found with \( M_{h^\pm} > 100 \text{ GeV} \) over the parameter space.

In figs. 3 and 4 the predicted lightest neutralino mass is shown versus the electroweak VEV ratio \( \tan \beta \) and the Higgs mixing parameter \( \mu \), and in figs. 5 and 6 the predicted lightest chargino mass is plotted versus the same parameters. It is important to note that the hypothesis made in Eq. 9 severely restricts the allowed parameter space for \( \tan \beta \) and \( \mu \). It is seen from Eq. 9, with \( M_2 = 0 \), that the absolute value of the \( \mu \) parameter is constrained. With \( M_2 = 0 \) in Eq. 7 \( \mu \) acts as a seesaw scale and as \( \mu \) increases, the lightest chargino mass is pushed down. Thus, in this scenario there is an upper limit on the allowed \( \mu \) value. Similarly, as can be seen...
from figs. 3 and 5 the small values for $\tan \beta$, with $\tan \beta \approx 1.1 - 1.4$ are preferred.

In this paper I examined how low energy dynamical SUSY breaking may arise from superstring derived models. The gauge mediated SUSY breaking scenarios have the important property of naturally suppressing supersymmetric contributions to neutral flavor changing transitions. In a specific superstring derived standard–like model I have shown that requiring potentially realistic fermion mass spectrum may result in a small hidden gauge group, like $SU(3)$. In the case that the hidden gauge group is broken entirely to $SU(3)_H$ the nonperturbative hidden gauge dynamics may indeed result in supersymmetry breaking at a low scale, in accordance with the gauge mediated dynamical SUSY breaking scenarios. The messenger sector states are obtained in the superstring model from sectors which arise due to the “Wilson–line” breaking of the $SO(10)$ gauge symmetry. As a result the interaction terms with the Standard Model states are suppressed. This illustrates that superstring models can provide the symmetries which are needed to suppress the supersymmetric contributions to flavor changing neutral currents. In this case the SUSY breaking is indeed transmitted to the observable sector only by the Standard Model gauge interaction which are generation blind. Thus, in the string inspired gauge mediated dynamical SUSY breaking scenario the suppression of flavor changing neutral currents arises naturally. It is also noted that the same color triplet fields arise as stable states in the superstring models and can be good dark matter candidate. Interestingly the mass scale needed for this type of color triplets to be good dark matter candidates is roughly the same as the mass scale for these states to serve as the messenger sector in the dynamical SUSY breaking scenarios. Motivated from the problem of string gauge coupling unification I have made the hypothesis that the messenger sector in the string motivated dynamical supersymmetry breaking scenario consists solely of color triplets. This hypothesis results in specific predictions for the superparticle spectrum which will be tested in the forthcoming LEP2 experiments. In particular, the lightest chargino mass is predicted to be below the $W$–boson mass. More detailed analysis of the composition of the lightest chargino and neutralino, which results from this hypothesis, and their experimental signature is of further interest.

I would like to thank I. Antoniadis, S. Dimopoulos, G. Farrar S. Thomas, F. Zwirner and especially J. Wells for important discussions, and the CERN theory division for its hospitality while this work was conducted. This work is supported in part by DOE Grant No. DE-FG-0586ER40272.
References

[1] M. Dine, W. Fischler and M. Srednicki, *Nucl. Phys.* **B189** (1981) 575;
S. Dimopoulos and S. Raby, *Nucl. Phys.* **B192** (1981) 353;
L. Alvarez–Gaume, M. Claudson and M. Wise, *Nucl. Phys.* **B207** (1982) 96;
C.R. Nappi and B.A. Ovrut, *Phys. Lett.* **B113** (1982) 175.

[2] M. Dine and A. Nelson, *Phys. Rev.* **D48** (1993) 1277;
M. Dine, A. Nelson and Y. Shirman, *Phys. Rev.* **D51** (1995) 1362;
M. Dine, A. Nelson, Y. Nir and Y. Shirman, *Phys. Rev.* **D53** (1996) 2658;
S. Dimopoulos, M. Dine, S. Raby and S. Thomas, *Phys. Rev. Lett.* **76** (1996) 3494;
S. Ambrosanio *et al.*, *Phys. Rev. Lett.* **76** (1996) 3498;
S. Dimopoulos, S. Thomas and J. Wells, [hep-ph/9604452];
G. Dvali, G. Giudice and A. Pomarol, [hep-ph/9603238];
K.S. Babu, C. Kolda and F. Wilczek, [hep-ph/9603238].

[3] J.P. Derendinger, L.E. Ibanez and H.P. Nilles, *Phys. Lett.* **B155** (1985) 65;
M. Dine, R. Rhom, N. Seiberg and E. Witten, *Phys. Lett.* **B156** (1985) 55;
S. Ferrara, N. Magnoli, T. Taylor and G. Veneziano, *Phys. Lett.* **B245** (1990) 409;
J.A. Casas, Z. Lalak, C. Munoz, G.G. Ross, *Nucl. Phys.* **B347** (1990) 243;
M. Cvetic *et al.*, *Nucl. Phys.* **B361** (1991) 194.

[4] A.E. Faraggi and E. Halyo, *Int. J. Mod. Phys.* **A11** (1996) 2357.

[5] A.E. Faraggi, *Nucl. Phys.* **B428** (1994) 111, [hep-ph/9403312].

[6] S. Chang, C. Corianò and A.E. Faraggi, [hep-ph/9603272] and [hep-ph/9605325].

[7] P. Langacker and M. Luo, *Phys. Rev.* **D44** (1991) 817;
J. Ellis, S. Kelley, and D. V. Nanopoulos, *Phys. Lett.* **B260** (1991) 131;
U. Amaldi, W. de Boer, and H. Füstenau, *Phys. Lett.* **B260** (1991) 447.

[8] P. Ginsparg, *Phys. Lett.* **B197** (1987) 139;
M. Dine and N. Seiberg, *Phys. Rev. Lett.* **55** (1985) 366;
V.S. Kaplunovsky, *Nucl. Phys.* **B307** (1988) 145; Erratum: *ibid.* **B382** (1992) 436, [hep-th/9205070].

[9] K.R. Dienes and A.E. Faraggi, *Phys. Rev. Lett.* **75** (1995) 2646, [hep-th/9505018];
*Nucl. Phys.* **B457** (1995) 409, [hep-th/9505048].

[10] A.E. Faraggi, *Phys. Lett.* **B302** (1993) 202, [hep-ph/9301268].
[11] I. Antoniadis, J. Ellis, S. Kelley and D.V. Nanopoulos, *Phys. Lett.* **B272** (1991) 31;
S. Kelley, J. Lopez, and D.V. Nanopoulos, *Phys. Lett.* **B278** (1992) 140;
I. Antoniadis, G.K. Leontaris, and N.D. Tracas, *Phys. Lett.* **B279** (1992) 58;
M.K. Gaillard and R. Xiu, *Phys. Lett.* **B296** (1992) 71, hep-ph/9206206;
S.P. Martin and P. Ramond, *Phys. Rev.* **D51** (1995) 6515, hep-ph/9501244.

[12] A. Font, L.E. Ibanez, F. Quevedo and A. Sierra, *Nucl.Phys.* **B331** (1990) 421;
D. Bailin, A. Love and S. Thomas, *Nucl. Phys.* **B298** (1988) 75;
J.A. Casas, E.K. Katehou and C. Muñoz, *Nucl. Phys.* **B317** (1989) 171.

[13] H. Kawai, D.C. Lewellen, and S.-H.H. Tye, *Nucl. Phys.* **B288** (1987) 1;
I. Antoniadis, C. Bachas, and C. Kounnas, *Nucl. Phys.* **B289** (1987) 87.

[14] I. Antoniadis, J. Ellis, J. Hagelin, and D.V. Nanopoulos, *Phys. Lett.* **B231** (1989) 65;
A.E. Faraggi, D.V. Nanopoulos and K. Yuan, *Nucl. Phys.* **B335** (1990) 347;
I. Antoniadis, G.K. Leontaris, and J. Rizos, *Phys. Lett.* **B245** (1990) 161;
J. Lopez, D.V. Nanopoulos, and K. Yuan, *Nucl. Phys.* **B399** (1993) 654;
S. Chaudhuri, G. Hockney, and J. Lykken, hep-th/9510241;
G.K. Leontaris, hep-ph/9601337.

[15] A.E. Faraggi, *Phys. Lett.* **B278** (1992) 131.

[16] A.E. Faraggi, *Nucl. Phys.* **B387** (1992) 239, hep-th/9208024.

[17] A.E. Faraggi, *Phys. Lett.* **B339** (1994) 223, hep-ph/9408333.

[18] A.E. Faraggi, *Phys. Lett.* **B326** (1994) 62, hep-ph/9311312.

[19] S. Kalara, J. Lopez, and D.V. Nanopoulos, *Nucl. Phys.* **B353** (1991) 650.

[20] A.E. Faraggi, *Nucl. Phys.* **B403** (1993) 101, hep-th/9208023; *Nucl. Phys.* **B407** (1993) 57, hep-ph/9210256.

[21] A.E. Faraggi and E. Halyo, *Phys. Lett.* **B307** (1993) 305, hep-ph/9301261; *Nucl. Phys.* **B416** (1994) 63, hep-ph/9306233.

[22] The ALEPH collaboration, *Phys. Lett.* **B373** (1996) 246;
The DELPHI collaboration, CERN–PPE/96–75;
The L3 collaboration, CERN–PPE/96–29;
The OPAL collaboration, CERN–PPE/96–20.

[23] The second reference in [22] provides an analysis of the constraints on the lightest chargino mass as a function of the sneutrino mass.

12
Figure 1: The predicted value of the lightest scalar superparticle versus tan $\beta$, with $1 \leq \tan \beta \leq 3$, $-50 \text{ GeV} \leq \mu \leq 50 \text{ GeV}$ and $\Lambda_3 = 100 \text{ TeV}$. 
Figure 2: The predicted values of the lightest neutralino and the lightest chargino for the same parameters values as in figure 1.

Figure 3: The predicted value of the lightest neutralino mass versus $\tan \beta$ the electroweak VEVs ratio, for the same parameters values as in figure 1.
Figure 4: The predicted value of the lightest neutralino mass versus $\mu$ the Higgs mixing parameter, for the same parameters values as in figure 1.

Figure 5: The predicted value of the lightest chargino mass versus $\tan \beta$ the electroweak VEVs ratio, for the same parameters values as in figure 1.
Figure 6: The predicted value of the lightest chargino mass versus $\mu$ the Higgs mixing parameter, for the same parameters values as in figure 1.