STERILE NEUTRINOS, DARK MATTER, AND RESONANCES in ψ'MSSM

1 Introduction

- The subgroup $SO(10) \times U(1)_\psi$ of E_6 can be decomposed, via $SU(5)$, to the MSSM gauge group times $U(1)_\chi \times U(1)_\psi$.
- One combination of these $U(1)$'s, denoted as $U(1)_{\psi'}$, is assumed here to be broken at a scale at least an order of magnitude greater than the TeV scale of soft SUSY breaking.
- We refer to the MSSM accompanied by $U(1)_{\psi'}$ as $\psi'\mathrm{MSSM}$.
- The RH neutrino in the 16-plet of $SO(10)$ is a $U(1)_{\psi'}$ singlet.
- This enables the three RH neutrinos to acquire large masses, so that the seesaw and leptogenesis scenarios can apply.
- We employ a $U(1)$ R symmetry such that dimension five and higher operators potentially causing proton decay are eliminated.
- The MSSM μ problem is resolved and the usual LSP of MSSM remains a compelling dark matter candidate.
- The three $SO(10)$ singlet sterile neutrino matter fields can only acquire tiny masses $\lesssim 0.1$ eV if $U(1)_{\psi'}$ is broken around 10 TeV.
- The effective number of neutrinos at NS is changed by $\simeq 0.29$.
- The lightest sterile sneutrino and two more particles stabilized by discrete symmetries, can be additional CDM candidates.
- If the breaking scale of $U(1)_{\psi'}$ is increased to 10^3 TeV, the sterile neutrinos become plausible candidates for keV scale warm DM.

- The contribution of the D-term for $U(1)_{\psi'}$ to the mass m_h of the lightest Higgs boson of MSSM can be appreciable.
- So, in the decoupling limit, the observed value of $m_h = 125$ GeV can be obtained with relatively light stop quarks.
- In addition to the Z' gauge boson associated with $U(1)_{\psi'}$, the model predicts diphoton and diquark resonances in the TeV range.
- A high luminosity or energy LHC upgrade may find them.
- The $U(1)_{\psi'}$ breaking produces superconducting strings which may be present in our galaxy.
- If the breaking scale is not too high, a 100 TeV collider may be able to make these strings.

2 The model

- Consider a SUSY model with gauge group $G_{\rm SM} \times U(1)_{\psi'}$, where $G_{\rm SM}$ is SM gauge group.
- The GUT-normalized generator $Q_{\psi'}$ of $U(1)_{\psi'}$ is given by

$$Q_{\psi'} = \frac{1}{4}(Q_\chi + \sqrt{15} Q_\psi).$$

- Here Q_χ and Q_ψ are, respectively, the GUT-normalized generators of the $U(1)_\chi$ in $SO(10)$ which commutes with $SU(5)$ and the $U(1)_\psi$ in E_6 which commutes with $SO(10)$.
- $U(1)_{\psi'}$ is to be spontaneously broken at a scale M.

- The important part of the W is

$$
\begin{aligned}
W &= y_u H_u^1 q u^c + y_d H_d^1 q d^c + y_\nu H_u^1 l \nu^c + y_e H_d^1 l e^c + \frac{1}{2} M_{\nu^c} \nu^c \nu^c \\
&\quad + \lambda_\mu^i N H_u^i H_d^i + \kappa S (N \bar{N} - M^2) + \lambda_D^i N D_i D_i^c + \lambda_q^i D_i q q \\
&\quad + \lambda_{q^c}^i D_i^c u^c d^c + \lambda_L S L \bar{L} + \lambda_{H_d}^\alpha \nu^c \bar{L} H_d^\alpha + \lambda_N^i N_i N_i \frac{\bar{N}^2}{2 m_{\mathrm{P}}}.
\end{aligned}
$$

- y_u, y_d, y_ν, y_e are the Yukawa couplings.
- q, u^c, d^c, l, ν^c, e^c are the usual quark and lepton superfields of MSSM including the right handed neutrinos ν^c.
- H_u^i, H_d^j $(i, j = 1, 2, 3)$ are $SU(2)_{\mathrm{L}}$ doublets with $Y = 1/2$, $-1/2$.
- N, $\bar{N}$ is a conjugate pair of SM singlets and S is a gauge singlet.
- The coupling $\lambda_\mu^{ij} N H_u^i H_d^j$ is diagonalized and a Z_2 symmetry under which H_u^α and H_d^α $(\alpha = 2, 3)$ are odd is imposed.
- So, only H_u^1, H_d^1 couple to quarks and leptons and are the standard electroweak Higgs superfields.
- D_i and D_i^c $(i = 1, 2, 3)$ are color triplets and antitriplets with $Y = -1/3$ and $1/3$ and the coupling $\lambda_D^{ij} N D_i D_j^c$ is diagonalized.
- N_i $(i = 1, 2, 3)$ are SM singlets and $\lambda_N^{ij} N_i N_j \bar{N}^2 / 2 m_{\mathrm{P}}$ is again diagonalized.
- We impose an extra Z_2' under which the N_i's are odd.
- To achieve MSSM gauge couplings unification, we introduced a pair of $SU(2)_{\mathrm{L}}$ doublets L and $\bar{L}$ with $Y = -1/2$ and $1/2$.
- q, u^c, d^c, l, ν^c, e^c, H_u^i, H_d^i, D_i, D_i^c, and N_i form three complete E_6 27-plets, while N, $\bar{N}$ and L, $\bar{L}$ are conjugate pairs from incomplete E_6 multiplets.

Superfields	Represention under G_{SM}	Extra Symmetries Z_2	Z_2'	R	$2\sqrt{10}Q_{\psi'}$
		Matter Superfields			
q	$(\mathbf{3},\mathbf{2},1/6)$	$+$	$+$	$1/2$	1
u^c	$(\bar{\mathbf{3}},\mathbf{1},-2/3)$	$+$	$+$	$1/2$	1
d^c	$(\bar{\mathbf{3}},\mathbf{1},1/3)$	$+$	$+$	$1/2$	2
l	$(\mathbf{1},\mathbf{2},-1/2)$	$+$	$+$	0	2
ν^c	$(\mathbf{1},\mathbf{1},0)$	$+$	$+$	1	0
e^c	$(\mathbf{1},\mathbf{1},1)$	$+$	$+$	1	1
H_u^α	$(\mathbf{1},\mathbf{2},1/2)$	$-$	$+$	1	-2
H_d^α	$(\mathbf{1},\mathbf{2},-1/2)$	$-$	$+$	1	-3
D_i	$(\mathbf{3},\mathbf{1},-1/3)$	$+$	$+$	1	-2
D_i^c	$(\bar{\mathbf{3}},\mathbf{1},1/3)$	$+$	$+$	1	-3
N_i	$(\mathbf{1},\mathbf{1},0)$	$+$	$-$	1	5
		Higgs Superfields			
H_u^1	$(\mathbf{1},\mathbf{2},1/2)$	$+$	$+$	1	-2
H_d^1	$(\mathbf{1},\mathbf{2},-1/2)$	$+$	$+$	1	-3
S	$(\mathbf{1},\mathbf{1},0)$	$+$	$+$	2	0
N	$(\mathbf{1},\mathbf{1},0)$	$+$	$+$	0	5
$\bar{N}$	$(\mathbf{1},\mathbf{1},0)$	$+$	$+$	0	-5
		Extra $SU(2)_{\mathrm{L}}$ Doublet Superfields			
L	$(\mathbf{1},\mathbf{2},-1/2)$	$-$	$+$	0	-3
$\bar{L}$	$(\mathbf{1},\mathbf{2},1/2)$	$-$	$+$	0	3

- Here, we summarize the fields and their transformation properties.
- The symmetries allow also the following higher order terms:

$$\nu^c H_u^\alpha L N,\, e^c H_d^\alpha L \bar{N},\, H_u^1 H_u^1 ll,\, H_u^\alpha H_u^\beta ll,\, H_u^1 H_d^\alpha l\bar{L},\, H_u^\alpha H_d^1 l\bar{L},$$
$$H_d^1 H_d^1 \bar{L}\bar{L},\, H_d^\alpha H_d^\beta \bar{L}\bar{L},\, qu^c qd^c\bar{N},\, qu^c e^c l\bar{N},\, qd^c\nu^c l\bar{N},$$
$$e^c\nu^c LLN,\, H_u^\alpha qd^c lL,\, H_u^1 H_u^\alpha lLN,\, H_u^1 H_u^1 LLNN,$$
$$H_u^\alpha H_u^\beta LLNN,\, H_u^\alpha qu^c l\bar{L}\bar{N},\, H_d^\alpha qd^c l\bar{L}\bar{N},\, \nu^c H_d^1 l\bar{L}\bar{L}\bar{N},$$
$$e^c H_u^1 lLLN,\, qd^c Lqd^c L,\, D_i^c u^c u^c \bar{L}\bar{L}\bar{N},\, D_i^c d^c d^c LLN,$$
$$e^c qd^c lLL,\, H_u^1 qd^c LLN,\, H_d^1 qu^c \bar{L}\bar{L}\bar{N},\, H_d^1 H_d^\alpha l\bar{L}\bar{L}\bar{L}\bar{N},$$
$$H_u^\alpha e^c LLLNN,\, \nu^c qu^c l\bar{L}\bar{L}\bar{N}\bar{N},\, qu^c qu^c \bar{L}\bar{L}\bar{N}\bar{N},$$
$$e^c e^c LLLLNN,\, H_d^\alpha qu^c l\bar{L}\bar{L}\bar{L}\bar{N}\bar{N}.$$

- All the couplings can be multiplied by $N\bar{N}/m_{\mathrm{P}}^2$, $L\bar{L}/m_{\mathrm{P}}^2$, and

$\bar{L}l\bar{N}\bar{L}l\bar{N}/m_{\rm P}^6$ arbitrarily many times.

- We assign baryon number $B=-2/3$ and $2/3$ to D_i and D_i^c.
- We then see that $U(1)_B$ is automatically present to all orders in W and, thus, fast proton decay is avoided.

3 $U(1)_{\psi'}$ breaking

- Assume that the breaking scale of $U(1)_{\psi'}$ is much bigger than the electroweak scale so that this breaking is not affected by it.
- So, the $U(1)_{\psi'}$ breaking can be discussed by considering only

$$\delta W=\kappa S(N\bar{N}-M^2).$$

- This gives the scalar potential

$$\begin{aligned} V &= \kappa^2|N\bar{N}-M^2|^2+\kappa^2|S|^2(|N|^2+|\bar{N}|^2) \\ &\quad+\big(A\kappa SN\bar{N}-(A-2m_{3/2})\kappa M^2S+{\rm H.c.}\big) \\ &\quad+m_0^2(|N|^2+|\bar{N}|^2+|S|^2)+{\rm D-terms}. \end{aligned}$$

- M and κ are made real and positive by field rephasing.
- $m_{3/2}$ is the gravitino mass, $A\sim m_{3/2}$ is the coefficient of the trilinear soft terms taken real and positive, and $m_0\sim m_{3/2}$.
- We assumed minimal SUGRA so that the coefficients of the trilinear and linear soft terms are related as shown.
- Vanishing of the D-terms $\Rightarrow |N|=|\bar{N}|\Rightarrow \bar{N}^*=e^{i\vartheta}N$, while minimization of the potential requires that $\vartheta=0$.
- So, N and $\bar{N}$ can be rotated to the positive real axis by $U(1)_{\psi'}$.

- We find that the scalar potential is minimized at

$$\langle S\rangle = -\frac{m_{3/2}}{\kappa}\left(1+\sum_{n\geq 1} c_n \left(\frac{m_{3/2}}{M}\right)^n\right),$$

$$\langle N\rangle = \langle \bar{N}\rangle \equiv \frac{N_0}{\sqrt{2}} = M\left(1+\sum_{n\geq 1} d_n \left(\frac{m_{3/2}}{M}\right)^n\right),$$

where c_n, d_n are numerical coefficients of order unity.

- For $M \gg m_{3/2}$, these formulas can be approximated as follows:

$$\langle S\rangle \simeq -\frac{m_{3/2}}{\kappa}, \quad \frac{N_0^2}{2} \simeq M^2 + \frac{A m_{3/2} - m_{3/2}^2 - m_0^2}{\kappa^2}.$$

- The trilinear and linear soft terms play an important role.
- Substituting $\langle N\rangle$, $\langle \bar{N}\rangle$, these terms yield a linear term in S which, together with the mass term of S, generates a VEV $\sim$ TeV for S.
- Then substituting $\langle S\rangle$ in $\lambda_L S L \bar{L}$, the superfields L, $\bar{L}$ acquire a mass $m_L = \lambda_L |\langle S\rangle| = \lambda_L m_{3/2}/\kappa$.
- The MSSM μ term is obtained by substituting $\langle N\rangle$ in $\lambda_\mu^1 N H_u^1 H_d^1$ with $\mu = \lambda_\mu^1 N_0/\sqrt{2}$.
- Also H_u^α, H_d^α $(\alpha = 2,3)$ and D_i, D_i^c acquire masses $\sim$ TeV from $\lambda_\mu^\alpha N H_u^\alpha H_d^\alpha$ and $\lambda_D^i N D_i D_i^c$ respectively.
- The mass spectrum of the scalar $S - N - \bar{N}$ system can be constructed by substituting $N = \langle N\rangle + \delta\tilde{N}$ and $\bar{N} = \langle \bar{N}\rangle + \delta\tilde{\bar{N}}$.
- For exact SUSY, we find two complex scalar fields S and $\theta = (\delta\tilde{N} + \delta\tilde{\bar{N}})/\sqrt{2}$ with equal masses $m_S = m_\theta = \sqrt{2}\kappa M$.
- Soft SUSY breaking mixes these fields yielding a mass splitting.
- The $U(1)_{\psi'}$ breaking generates superconducting strings with relatively small tension, which satisfies all the experimental bounds.

4 Electroweak Symmetry Breaking

- The standard V for the radiative electroweak symmetry breaking in MSSM is modified in the present model.
- One modification originates from the D-term for $U(1)_{\psi'}$:

$$V_D = \frac{g_{\psi'}^2}{80}\left[-2|H_u|^2 - 3|H_d|^2 + 5\left(|N|^2 - |\bar{N}|^2\right)\right]^2.$$

- $g_{\psi'}$ is the GUT-normalized gauge coupling for $U(1)_{\psi'}$ and H_u, H_d are the neutral components of the scalar parts of H_u^1, H_d^1.
- To integrate out to one loop N and $\bar{N}$, we express them in terms of the canonically normalized real scalars δN, $\delta\bar{N}$, φ, $\bar{\varphi}$:

$$N = \frac{1}{\sqrt{2}}(N_0 + \delta N)e^{\frac{i\varphi}{N_0}}, \quad \bar{N} = \frac{1}{\sqrt{2}}(N_0 + \delta\bar{N})e^{\frac{i\bar{\varphi}}{N_0}}.$$

- The combination $|N|^2 - |\bar{N}|^2$ in the D-term then becomes

$$|N|^2 - |\bar{N}|^2 = \sqrt{2}N_0\eta + \eta\xi,$$

with

$$\eta = \frac{\delta N - \delta\bar{N}}{\sqrt{2}}, \quad \xi = \frac{\delta N + \delta\bar{N}}{\sqrt{2}}.$$

- The D-term can now be expanded up to second order in η, ξ:

$$V_D = \frac{g_{\psi'}^2}{80}\left[E^2 + 10\sqrt{2}N_0E\eta + 50N_0^2\eta^2 + \cdots\right],$$

where $E \equiv -2|H_u|^2 - 3|H_d|^2$.

- Note that we ignored the mixed quadratic term $\propto \eta\xi$ since its coefficient is much smaller than the coefficient of the η^2 term.

- We see that integrating out the heavy states reduces to the calculation of a path integral over η.
- Substitute N, $\bar{N}$ in terms of δN, $\delta \bar{N}$, φ, $\bar{\varphi}$, keeping only η-dependent terms up to 2nd order and substituting $\langle S \rangle$ and N_0, the potential V becomes

$$\delta V \simeq m_N^2 \eta^2 \quad \text{with} \quad m_N^2 \equiv m_{3/2}^2 + m_0^2.$$

- Adding δV to the D-term potential, we obtain the potential

$$V_\eta = \frac{g_{\psi'}^2 E^2}{80} \left(1 + \frac{5 g_{\psi'}^2 N_0^2}{8 m_N^2}\right)^{-1} + \left(m_N^2 + \frac{5 g_{\psi'}^2 N_0^2}{8}\right) \times \left(\eta + \frac{g_{\psi'}^2 N_0 E}{8\sqrt{2}\left(m_N^2 + \frac{5 g_{\psi'}^2 N_0^2}{8}\right)}\right)^2 + \cdots .$$

- Calculating the path integral

$$\int (d\eta) e^{-i V_\eta \mathcal{V}}$$

($\mathcal{V}$=the spacetime volume), we then find the term

$$\delta V_D \simeq \frac{g_{\psi'}^2}{80} \left[2|H_u|^2 + 3|H_d|^2\right]^2 \left(1 + \frac{m_{Z'}^2}{2 m_N^2}\right)^{-1}$$

to be added to the usual electroweak symmetry breaking potential.
- Here $m_{Z'} = \sqrt{5} g_{\psi'} N_0 / 2$ is the mass of the Z' gauge boson.
- Another modification of the electroweak potential comes from the integration of the heavy field S with mass $\sqrt{2}\kappa M$.

- This gives the extra term in the electroweak potential

$$-\frac{1}{2}\tilde{\lambda}_\mu^2|H_u|^2|H_d|^2, \quad \text{with} \quad \tilde{\lambda}_\mu = \frac{1}{\sqrt{2}}\lambda_\mu^1,$$

which reduces the well-known NMSSM term $\tilde{\lambda}_\mu^2|H_u|^2|H_d|^2$.

- From the modified electroweak V, we find the mass^2 of the lightest neutral CP-even Higgs boson in the decoupling limit ($m_A \gg m_Z$):

$$m_h^2 = m_Z^2\cos^2 2\beta + 4cv^2(2\sin^2\beta + 3\cos^2\beta)^2 + \lambda_\mu^2 v^2 \sin^2 2\beta.$$

- Here $\lambda_\mu \equiv \tilde{\lambda}_\mu/\sqrt{2}$, $v = 246 \text{ GeV}$ and

$$c = \frac{g_{\psi'}^2}{80}\left(1 + \frac{m_{Z'}^2}{2m_N^2}\right)^{-1}.$$

5 Diphoton Resonances

- The real (pseudo)scalar components θ_1 (θ_2) of $\theta = (\theta_1 + i\theta_2)/\sqrt{2}$ with mass $m_\theta = \sqrt{2}\kappa M$ can be produced at the LHC by gluon fusion via a fermionic D_i, D_i^c loop.

- They can then decay into two photons via the same loop diagram as well as a similar fermionic H_u^i, H_d^i loop.

- The cross section of the diphoton excess is

$$\sigma(pp \to \theta_m \to \gamma\gamma) \simeq \frac{C_{gg}}{m_\theta s \Gamma_{\theta_m}}\Gamma(\theta_m \to gg)\Gamma(\theta_m \to \gamma\gamma),$$

where $m = 1, 2$, $C_{gg} \simeq 3163$, $\sqrt{s} \simeq 13 \text{ TeV}$, and Γ_{θ_m} is the total decay width of θ_m.

- The decay widths of θ_m to two gluons or two photons are

$$\begin{aligned}
\Gamma(\theta_m \to gg) &= \frac{m_\theta^3 \alpha_s^2}{512\,\pi^3 \langle N\rangle^2} \left(\sum_{i=1}^{3} A_m(x_i) \right)^2 , \\
\Gamma(\theta_m \to \gamma\gamma) &= \frac{m_\theta^3 \alpha_Y^2 \cos^4\theta_W}{9216\,\pi^3 \langle N\rangle^2} \left[\sum_{i=1}^{3} A_m(x_i) + \right. \\
&\left. \frac{3}{2} \sum_{i=1}^{3} A_m(y_i) \left(1 + \frac{\alpha_2 \tan^2\theta_W}{\alpha_Y} \right) \right]^2 .
\end{aligned}$$

- $A_1(x) = 2x + (1-x)A_2(x)$, $A_2(x) = 2x \arcsin^2(1/\sqrt{x})$, $x_i = 4m_{D_i}^2/m_\theta^2$, $y_i = 4m_{H_i}^2/m_\theta^2$, $m_{D_i} = \lambda_D^i \langle N\rangle$, $m_{H_i} = \lambda_\mu^i \langle N\rangle$.
- The cross section simplifies if θ_m decay predominantly into gluons, i.e. $\Gamma_{\theta_m} \simeq \Gamma(\theta_m \to gg)$:

$$\sigma(pp \to \theta_m \to \gamma\gamma) \simeq 7.3 \times 10^6 \, \frac{\Gamma(\theta_m \to \gamma\gamma)}{m_\theta} \text{ fb}.$$

- Assume that x_i, y_i are just above unity, which maximizes $A_1(x_i)$, $A_2(y_i)$ while still blocks the decay of θ_m to D_i, D_i^c and H_u^i, H_d^i.
- We also consider the decay of θ_2 since $A_2(x) > A_1(x)$.
- In this case, the cross section becomes

$$\sigma(pp \to \theta_2 \to \gamma\gamma) \simeq 5.5 \left(\frac{m_\theta}{\langle N\rangle} \right)^2 \text{fb} \simeq 11\,\kappa^2 \text{ fb}.$$

- θ could also decay into a bosonic L, $\bar{L}$ pair.
- Our estimate of the cross section holds if that the direct decay of θ into a D_i, D_i^c, or H_u^i, H_d^i, or L, $\bar{L}$ is kinematically blocked.

- This is achieved for

$$\kappa \lesssim \sqrt{2}\lambda_D^i,\ \sqrt{2}\lambda_\mu^i,\ 2\lambda_L \frac{m_{3/2}}{m_\theta}.$$

- Note that our estimate of the maximal diphoton excess corresponds to saturating the first two of these inequalities.
- For simplicity and for not disturbing the MSSM gauge coupling unification, we choose to saturate the third inequality too.

6 A Numerical Example

- $g_{\psi'}$ unifies with the MSSM gauge couplings provided that its value at low energies is equal to about 0.45.
- Demanding that the Z' gauge boson mass $m_{Z'} \simeq \sqrt{5} g_{\psi'} M/\sqrt{2} > 3.8\ \mathrm{TeV}$, say, we then find $M \gtrsim 5.34\ \mathrm{TeV}$.
- As an example, we will set $M = 10\ \mathrm{TeV}$.
- We can show that κ, $\tilde{\lambda}_\mu$ remain perturbative up to the GUT scale provided that they are not much bigger than about 0.7.
- Requiring that the diphoton resonance mass $m_\theta = \sqrt{2}\kappa M \gtrsim 4.5\ \mathrm{TeV}$ as indicated by CMS, implies that $\kappa \gtrsim 0.32$.
- If the first two inequalities above are saturated, we have $0.5 \gtrsim \lambda_D^i, \lambda_\mu^i \gtrsim 0.22$.
- We set $\lambda_D^i \simeq \lambda_\mu^i \simeq 0.3 \Rightarrow \tilde{\lambda}_\mu \simeq 0.3$, $\kappa \simeq 0.42$, $m_{D_i} \simeq m_{H_i} \simeq 3\ \mathrm{TeV}$ ($\mu \simeq 3\ \mathrm{TeV}$), $m_\theta \simeq 6\ \mathrm{TeV}$, $m_{Z'} \simeq 7.1\ \mathrm{TeV}$.
- Saturating the third inequality too, we obtain $m_L \simeq 3\ \mathrm{TeV}$.
- For $\kappa \lesssim 0.7$, the resonance mass remains below $9.9\ \mathrm{TeV}$.

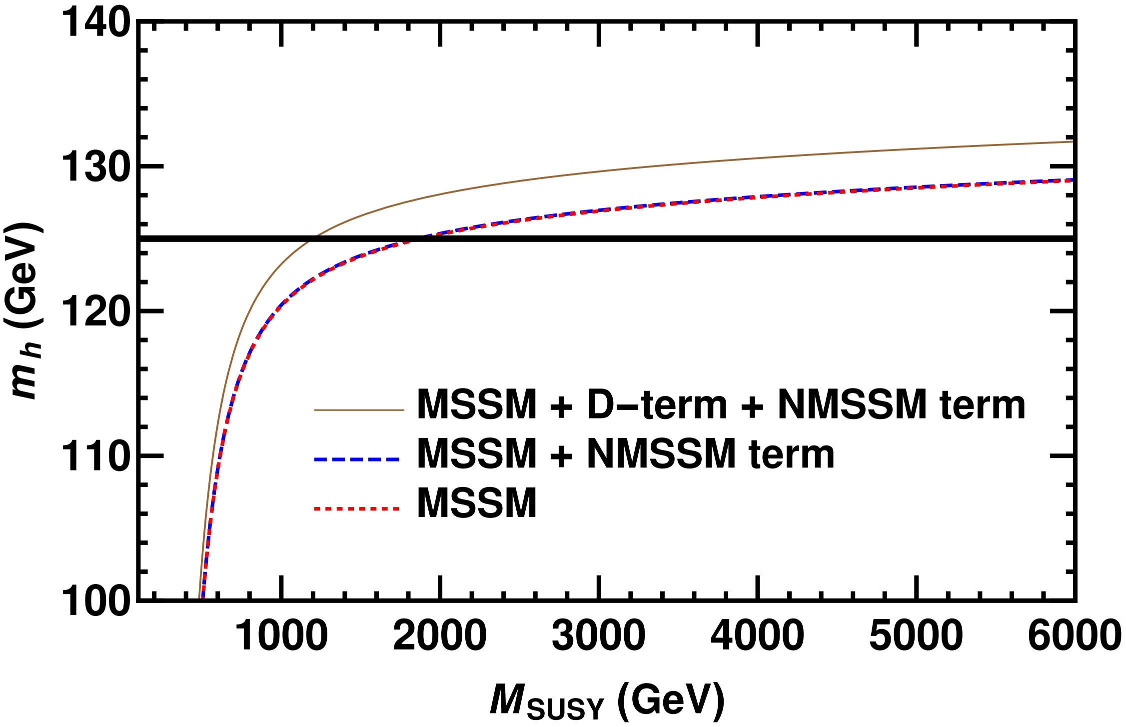

Figure 1: Higgs boson mass m_h in the decoupling limit and for maximal stop quark mixing versus $M_{\rm SUSY}$ for $M = 10$ TeV, $\tilde{\lambda}_\mu = 0.3$, $\tan\beta = 20$, and $m_{3/2} = 4$ TeV. The bold horizontal line corresponds to $m_h = 125$ GeV.

- We plot the Higgs mass m_h in the decoupling limit versus $M_{\rm SUSY}$, which is the geometric mean of the stop quark mass eigenvalues.
- We assume maximal stop quark mixing, which maximizes m_h, and include the two-loop radiative corrections to m_h in MSSM.
- The NMSSM and D-term contributions to m_h are also included.
- In this figure, $\tan\beta = 20$ and $m_{3/2} = 4$ TeV.
- The NMSSM correction is very small since $\tilde{\lambda}_\mu$ is relatively small.
- The D-term correction, however, is sizable and allows us to obtain the observed m_h with much smaller stop quark masses than the ones required in MSSM or NMSSM.
- Indeed, the inclusion of the D-term from $U(1)_{\psi'}$ reduces $M_{\rm SUSY}$ from about 1900 GeV to about 1200 GeV.

7 Sterile Neutrinos

- The sterile neutrinos, which are the fermionic parts of N_i, acquire masses $\sim 10^{-1}$ eV for $M \sim 10$ TeV via $\lambda_N^i N_i N_i \bar{N}^2/2m_{\rm P}$.
- These fermionic fields, which are stable on account of the Z_2' symmetry, can act as sterile neutrinos.
- In the early universe, sterile neutrinos are in equilibrium through reactions like $N_i\bar{N}_i \leftrightarrow$ a pair of SM particles via a Z' exchange.
- The interaction rate per sterile neutrino is $\Gamma_{N_i} \sim T^5/M^4$.
- The decoupling temperature $T_{\rm D}$ is then found from the condition $\Gamma_{N_i} \sim H =$the Hubble parameter, which implies that

$$T_{\rm D} \sim M \left(\frac{M}{m_{\rm P}}\right)^{\frac{1}{3}}.$$

- The strategy is the same as the one used for the SM neutrino decoupling via processes involving weak gauge boson exchange.
- For SM neutrinos, M should be the electroweak scale ~ 100 GeV, and the decoupling temperature turns out to be ~ 1 MeV.
- So, for $M \simeq 10$ TeV, $T_{\rm D}$ is expected to be $\simeq 460$ MeV, which is well above the critical temperature for the QCD transition.
- The effective number of massless degrees of freedom in equilibrium right after the decoupling of sterile neutrinos is 61.75.
- At decoupling of the SM neutrinos, this number becomes 10.75.
- Due to entropy conservation, the T of SM neutrinos is raised relative to that of the sterile neutrinos by a factor $(61.75/10.75)^{1/3}$.

- Consequently, the contribution of the three sterile neutrinos to the effective number of neutrinos at big bang nucleosynthesis is

$$\Delta N_\nu = 3 \times \left(\frac{10.75}{61.75}\right)^{\frac{4}{3}} \simeq 0.29.$$

- This is perfectly compatible with the Planck satellite bound

$$N_\nu = 3.15 \pm 0.23.$$

8 Dark Matter

- The bosonic N_i with mass $\sim m_{3/2}$ can decay into a fermionic N_i and a particle-sparticle pair via a Z' gaugino exchange.
- A necessary condition for this is that there exist sparticles which are lighter than the scalar N_i.
- If the decay of the lightest scalar N_i (denoted as $\hat{N}$) is kinematically blocked, this particle can contribute to the CDM.
- We estimate the freeze-out temperature $T_{\rm f}$ of $\hat{N}$ and its relic abundance $\Omega_{\hat{N}} h^2$ for the lowest $M \simeq 5.34$ TeV.
- The requirement that $\Omega_{\hat{N}} h^2$ equals the $\Omega_{\rm CDM} h^2 \simeq 0.12$ implies that $m_{\hat{N}} \simeq 1.25$ TeV and $T_{\rm f} \simeq 51$ GeV.
- The model possesses an accidental lepton parity symmetry $Z_2^{\rm lp}$ under which l, e^c, ν^c, L, $\bar{L}$ are odd.
- Combining $Z_2^{\rm lp}$ with the $Z_2^{\rm bp} \subset U(1)_B$, we obtain a matter parity symmetry $Z_2^{\rm mp}$ under which q, u^c, d^c, l, e^c, ν^c, L, $\bar{L}$ are odd.
- A R-parity is then generated combining $Z_2^{\rm mp}$ with fermion parity.

- Particles with negative R-parity except the bosonic L, $\bar{L}$ and the fermionic H_u^α, H_d^α, N_i decay to the LSP which is CDM candidate.
- Z_2 and R-parity $\Rightarrow$ the lightest state in the bosonic (fermionic) L, $\bar{L}$ and fermionic (bosonic) H_u^α, H_d^α is stable.
- We thus have two more candidates for CDM with their relic abundances depending on details.
- Finally, if $\langle N \rangle$ is increased to $\sim 10^3$ TeV, the sterile neutrinos become plausible candidates for keV scale warm dark matter.

9 Summary

- We appended $U(1)_{\psi'}$ to the MSSM gauge group.
- This $U(1)_{\psi'}$ is a linear combination of $U(1)_\chi$, $U(1)_\psi \subset E_6$.
- The three matter 27-plets in E_6 give rise to three $SO(10)$ singlet fermions N_i, called sterile neutrinos.
- For a relatively low (~ 10 TeV) breaking scale of $U(1)_{\psi'}$, the sterile neutrinos acquire masses $\lesssim 0.1$ eV.
- Their contribution as fractional cosmic neutrinos is acceptable.
- The model possesses many possible candidates for DM.
- The D-term for $U(1)_{\psi'}$ contributes appreciably to m_h and, thus $m_h = 125$ GeV can be obtained with relatively light stop quarks.
- The model predicts superconducting cosmic strings as well as diquark and diphoton resonances.
- The μ problem is naturally solved and the RH neutrinos masses are large allowing the seesaw and leptogenesis scenarios to apply.
- Baryon number is conserved to all orders in perturbation theory.