AN EXTENSION OF THE MSSM WITH DIPHOTON RESONANCES

1 Introduction

- The recently reported 750 GeV diphoton resonance by ATLAS and CMS requires the introduction of new physics beyond the SM.
- There is already a flurry of theoretical papers offering a variety of plausible extensions of the SM to explain this diphoton excess.
- Here, we propose a particular extension of MSSM which naturally yields resonance states in the TeV range.
- Namely, our proposal is based on a local $U(1)_{B-L}$ extension of the MSSM gauge symmetry.
- In contrast to the radiative electroweak symmetry breaking in MSSM, the additional $U(1)$ is spontaneously broken at tree level.
- The relevant W is uniquely determined by the gauge symmetry, global B and L conservation, and a $U(1)$ global R-symmetry.
- It utilizes an appropriate pair of Higgs superfields Φ, $\bar{\Phi}$, as well as a gauge singlet superfield S.
- The resonances arise from the scalar components of S, Φ, $\bar{\Phi}$.
- Their mass is determined, in the SUSY limit, by a dimensionless parameter $\ll$ the gauge coupling.
- Thus, the resonances can be much lighter than the Z' gauge boson associated with $B-L$, whose mass should be at least a few TeV.
- The spontaneous breaking of $U(1)_{B-L}$ leaves SUSY unbroken.

- The symmetry breaking scale M may be much larger than the soft SUSY breaking scale.
- W's of this type have previously been employed in the construction of SUSY hybrid inflation models.
- The scalar component of S acquires a non-zero VEV proportional to $m_{3/2}$ after SUSY breaking.
- This has been utilized in the past to resolve the MSSM μ problem.
- Here we also use this $\langle S \rangle$ to provide masses to vector-like fields which play a role in the production and decay of the resonances.
- The R-symmetry protects S from acquiring arbitrarily large mass.

2 The model

- The new local $U(1)_{B-L}$ symmetry is to be spontaneously broken at some scale M.
- We prefer to implement this breaking by a SUSY generalization of the Higgs mechanism.
- Motivated by MSSM, we require that W respects the global $U(1)_B$ and $U(1)_L$ symmetries and a global $U(1)$ R-symmetry.
- The full renormalizable superpotential is

$$\begin{aligned} W &= y_u H_u q u^c + y_d H_d q d^c + y_\nu H_u l \nu^c + y_e H_d l e^c \\ &\quad + \kappa S(\Phi\bar{\Phi} - M^2) + \lambda_\mu S H_u H_d + \lambda_{\nu^c} \bar{\Phi} \nu^c \nu^c \\ &\quad + \lambda_D S D \bar{D} + \lambda_q D q q + \lambda_{q^c} \bar{D} u^c d^c. \end{aligned}$$

- y_u, y_d, y_ν, y_e are Yukawa couplings with family indices suppressed.

- Also 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, H_d are the standard EW Higgs superfields.
- The gauge singlet S has the same R-charge as W, taken to be 2.
- So, H_u, H_d have opposite R-charges, which are brought to zero by a Y transformation.
- The R-charges of u^c and d^c , as well as of ν^c and e^c are equal.
- Consequently, B and L transformations can make the R-charges of q, u^c, d^c, l, ν^c, e^c all equal to unity.
- To determine the R and $B-L$ charges of the SM singlets Φ, $\bar{\Phi}$, we introduce the coupling $\bar{\Phi}\nu^c\nu^c$.
- This implies that their $B-L$ charge is 2, -2 respectively, and their R-charges are zero.
- Note that $\bar{\Phi}\nu^c\nu^c$ generates masses for the ν^c's after the breaking of $U(1)_{B-L}$ to its Z_2 subgroup by $\langle\Phi\rangle$, $\langle\bar{\Phi}\rangle$.
- We also introduce the coupling $SD\bar{D}$, where D ($\bar{D}$) are color triplet (antitriplet) and $SU(2)_L$ singlet superfields.
- To determine the charges of D, $\bar{D}$, we need an additional coupling.
- Taking Dqq, we find that the R-charges of D, $\bar{D}$ vanish.
- Also, the Y of D is -1/3 and, thus, the Y of $\bar{D}$ is 1/3.
- Finally, the $B-L$ of D is -2/3 and that of $\bar{D}$ is 2/3 with their L vanishing.
- Note that $\bar{D}u^cd^c$ is also present as it respects all the symmetries.
- The Z_2 subgroups of $U(1)_R$, $U(1)_{B-L}$ coincide with matter parity under which the (anti)quark, (anti)lepton superfields are odd.

- This Z_2 remains unbroken by all the soft SUSY breaking terms and all the VEVs.
- We summarize below all the superfields of the model together with their transformation properties and charges.

Superfields	Represenţions under G_{SM}	Global Symmetries B	L	R
	Matter Superfields			
q	$(\mathbf{3}, \mathbf{2}, 1/6)$	$1/3$	0	1
u^c	$(\bar{\mathbf{3}}, \mathbf{1}, -2/3)$	$-1/3$	0	1
d^c	$(\bar{\mathbf{3}}, \mathbf{1}, 1/3)$	$-1/3$	0	1
l	$(\mathbf{1}, \mathbf{2}, -1/2)$	0	1	1
ν^c	$(\mathbf{1}, \mathbf{1}, 0)$	0	-1	1
e^c	$(\mathbf{1}, \mathbf{1}, 1)$	0	-1	1
	Higgs Superfields			
H_u	$(\mathbf{1}, \mathbf{2}, 1/2)$	0	0	0
H_d	$(\mathbf{1}, \mathbf{2}, -1/2)$	0	0	0
S	$(\mathbf{1}, \mathbf{1}, 0)$	0	0	2
Φ	$(\mathbf{1}, \mathbf{1}, 0)$	0	-2	0
$\bar{\Phi}$	$(\mathbf{1}, \mathbf{1}, 0)$	0	2	0
	Vector-like Diquark Superfields			
D	$(\mathbf{3}, \mathbf{1}, -1/3)$	$-2/3$	0	0
$\bar{D}$	$(\bar{\mathbf{3}}, \mathbf{1}, 1/3)$	$2/3$	0	0

- Note that W is the most general renormalizable superpotential which obeys the symmetries of the model.
- Had we removed the B and L symmetries and kept only $U(1)_{B-L}$, the terms $\bar{D}ql$, Du^ce^c, and $Dd^c\nu^c$ would be present too.
- They would yield fast proton decay and other B and L violating effects.

- The spontaneous breaking of $U(1)_{B-L}$ to Z_2 will generate a network of local cosmic strings.
- Their string tension, determined by M, satisfies the most stringent relevant upper bound from pulsar timing arrays.
- The 'bare' MSSM μ term is replaced by SH_uH_d.
- So the μ term is generated after S acquires a non-zero VEV $\sim$ TeV from soft SUSY breaking.
- $\langle S \rangle$ is also responsible for generating masses for the diquarks D, $\bar{D}$, which may be found at the LHC.
- These masses are crucial in the production and decay of the diphoton resonances.
- The breaking of $U(1)_{B-L}$ implemented with S, Φ, $\bar{\Phi}$ delivers, for exact SUSY, four scalars all with the same mass $m_S = \sqrt{2}\kappa M$.
- Even for $M \gg 1$ TeV, m_S can be $\simeq 750$ GeV for κ small enough.
- Note that, after SUSY breaking, the four resonances may end up with significantly different masses.

3 Analysis of the Model

- The breaking of $U(1)_{B-L}$ is achieved via the term $\kappa S(\Phi\bar{\Phi} - M^2)$ which, for unbroken SUSY, gives the potential

$$V = \kappa^2|\Phi\bar{\Phi} - M^2| + \kappa^2|S|^2(|\Phi|^2 + |\bar{\Phi}|^2) + \mathrm{D-terms}.$$

- Here we made M, κ real and positive by field rephasing.
- Vanishing of the D-terms yields $|\Phi| = |\bar{\Phi}| \Longrightarrow \bar{\Phi}^* = e^{i\varphi}\Phi$.

- The F-terms vanish for $S = 0$, $\Phi\bar{\Phi} = M^2$, requiring $\varphi = 0$.
- Rotating Φ, $\bar{\Phi}$ to the positive real axis by a $B - L$ transformation, we find the SUSY vacuum

$$S = 0 \quad \text{and} \quad \Phi = \bar{\Phi} = M.$$

- The mass spectrum of the scalar $S - \Phi - \bar{\Phi}$ system is constructed by writing $\Phi = M + \delta\Phi$ and $\bar{\Phi} = M + \delta\bar{\Phi}$.
- For unbroken SUSY, we find two complex scalar fields S and $\theta = (\delta\Phi + \delta\bar{\Phi})/\sqrt{2}$ with equal masses $m_S = m_\theta = \sqrt{2}\kappa M$.
- Soft SUSY breaking can, of course, mix these fields and generate a mass splitting.
- For example, the trilinear soft term $A\kappa S\Phi\bar{\Phi}$ yields a mass^2 splitting $\pm\sqrt{2}\kappa M A$ with mass eigenstates $(S+\theta^*)/\sqrt{2}$, $(S-\theta^*)/\sqrt{2}$.
- We assume, for simplicity, that the mixing is small and ignore it.
- Consider the soft SUSY breaking potential terms

$$V_1 = A\kappa S\Phi\bar{\Phi} - (A - 2m_{3/2})\kappa M^2 S, \quad A \sim m_{3/2}$$

  arising from the W term $\kappa S(\Phi\bar{\Phi} - M^2)$.
- In minimal SURGA, the coefficients of the trilinear and linear soft terms are related as shown.
- Substituting $\Phi = \bar{\Phi} = M$, we obtain a linear term in S which, together with the mass term $2\kappa^2 M^2 |S|^2$, generates a VEV:

$$\langle S \rangle = -\frac{m_{3/2}}{\kappa}.$$

- From $\lambda_\mu S H_u H_d$, we obtain the μ term with $\mu = -\lambda_\mu m_{3/2}/\kappa$.

- The same VEV generates mass terms $m_D D\bar{D}$ for the vector-like superfields D, $\bar{D}$ via $\lambda_D S D\bar{D}$ with $m_D = -\lambda_D m_{3/2}/\kappa$.
- To preserve gauge coupling unification, we introduce vector-like colorless, $SU(2)_L$ doublets L, $\bar{L}$ equal in number to D, $\bar{D}$.
- Note that with up to four D, $\bar{D}$ and L, $\bar{L}$ pairs with masses $\sim \mathrm{TeV}$, the gauge couplings stay perturbative up to $M_{\rm GUT}$.
- The L, $\bar{L}$'s with a W coupling $\lambda_L S L\bar{L}$ can enhance the branching ratio of the decay of the scalars S and θ to photons.
- They also allow the decay into $W^\pm$.
- Introducing the W coupling Lle^c, the Y of L ($\bar{L}$) is -1/2 (1/2).
- Their B, L, and R-charges are all zero.
- These quantum numbers allow also the couplings SLH_u, $SH_d\bar{L}$, Lqd^c, $\bar{L}qu^c$, and $\bar{L}l\nu^c$.
- Substituting $\langle S\rangle$ in $\lambda_L S L\bar{L}$, L, $\bar{L}$ get a mass $m_L = -\lambda_L m_{3/2}/\kappa$.

4 Diphoton Resonance

- The real (pseudo)scalar S_1 (S_2) in $S = (S_1 + iS_2)/\sqrt{2}$ with $m_S = \sqrt{2}\kappa M$ for exact SUSY can be produced by g fusion.
- The relevant graphs via a fermionic D, $\bar{D}$ loop are shown in Fig. 1.
- In the absence of L, $\bar{L}$, they can decay into g, γ, or Z via the same diagram, but not to $W^\pm$ since D, $\bar{D}$ are $SU(2)_L$ singlets.
- The most promising decay channel to search for these resonances is into two γ with the relevant diagram also shown in Fig. 1.

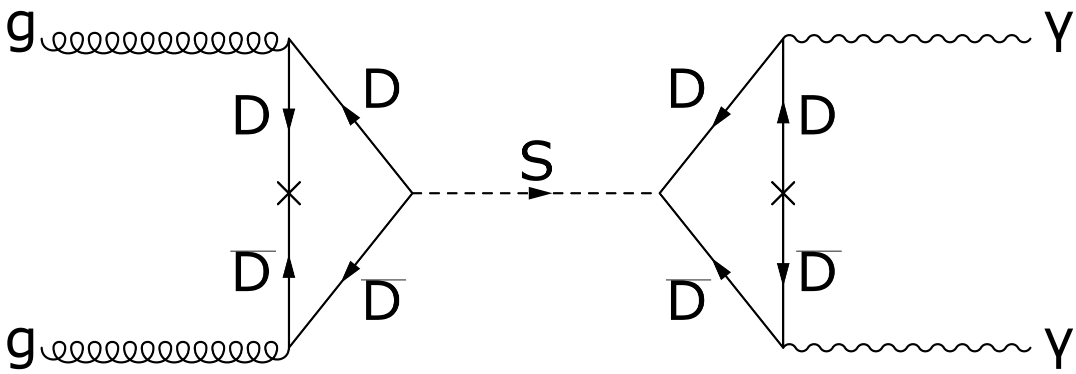

Figure 1: Production of the bosonic S by g fusion and its decay into γ. Solid (dashed) lines represent fermions (bosons). The arrows depict the chirality of the superfields and the crosses mass insertions.

- The cross section of the diphoton excess is

$$\sigma(pp \to S_i \to \gamma\gamma) \simeq \frac{C_{gg}}{m_S s \Gamma_{S_i}} \Gamma(S_i \to gg)\Gamma(S_i \to \gamma\gamma).$$

- Here $C_{gg} \simeq 3163$, $\sqrt{s} \simeq 13$ TeV, Γ_{S_i} = total decay width of S_i.
- The decay widths of S_i to two g or two γ are

$$\Gamma(S_i \to gg) = \frac{n^2 \alpha_s^2 m_S^3}{256\, \pi^3 \langle S \rangle^2} A_i^2(x),$$
$$\Gamma(S_i \to \gamma\gamma) = \frac{n^2 \alpha_Y^2 m_S^3 \cos^4 \theta_W}{4608\, \pi^3 \langle S \rangle^2} A_i^2(x).$$

- n is the number of D, $\bar{D}$ pairs with a common coupling λ_D to S.
- $A_1(x) = 2x[1 + (1 - x) \arcsin^2(\frac{1}{\sqrt{x}})]$, $A_2(x) = 2x \arcsin^2(\frac{1}{\sqrt{x}})$ with $x = 4m_D^2/m_S^2 > 1$.
- α_s, α_Y are the strong and Y fine-structure constants.
- If L, $\bar{L}$ are present, they also contribute to the Γ of S to γ via loop diagrams similar to the ones in the right part of Fig. 1.

- In this case, the equation for $\Gamma(S_i \to \gamma\gamma)$ is replaced by

$$\Gamma(S_i \to \gamma\gamma) = \frac{n^2 m_S^3 \alpha_Y^2 \cos^4\theta_W}{4608\,\pi^3 \langle S\rangle^2} A_i^2(x) \left[1 + \frac{3A_i(y)}{2A_i(x)}\left(1 + \frac{\alpha_2 \tan^2\theta_W}{\alpha_Y}\right)\right]^2.$$

- α_2 is the $SU(2)_L$ fine-structure constant, $y = 4m_L^2/m_S^2 > 1$.
- The cross section simplifies if S_i decay predominantly into g, namely, if $\Gamma_{S_i} \simeq \Gamma(S_i \to gg)$.
- In this case, one obtains $\sigma(pp \to S_i \to \gamma\gamma) \simeq 8$ fb if

$$\frac{\Gamma(S_i \to \gamma\gamma)}{m_S} \simeq 1.1 \times 10^{-6}.$$

- For x and y just above 1, the S_i decay to D, $\bar{D}$ and L, $\bar{L}$ is blocked and $A_1(x)$, $A_2(y)$ are maximized with $A_1 \simeq 2$, $A_2 \simeq \pi^2/2$.
- However, x close to 1 means $m_D \simeq 375$ GeV, which is excluded by ATLAS and CMS. So we take $m_D = 700$ GeV.
- Moreover, it is more beneficial to consider the decay of the S_2 since $A_2(x) > A_1(x)$ for all $x > 1$.
- We then find that the condition above is satisfied for $|\langle S\rangle| \simeq 758$ GeV, for $n = 3$, $m_S \simeq 750$ GeV.
- In this case, $\lambda_D \simeq 0.92$ and λ_L is just above 0.49.
- Note that the inclusion of L, $\bar{L}$ enhances the decay width of S_2 to γ by about a factor 58.5.
- For exact SUSY, the complex scalar S could decay into Higgsinos via the term $\lambda_\mu S H_u H_d$ if this is kinematically allowed – Fig. 2(a).

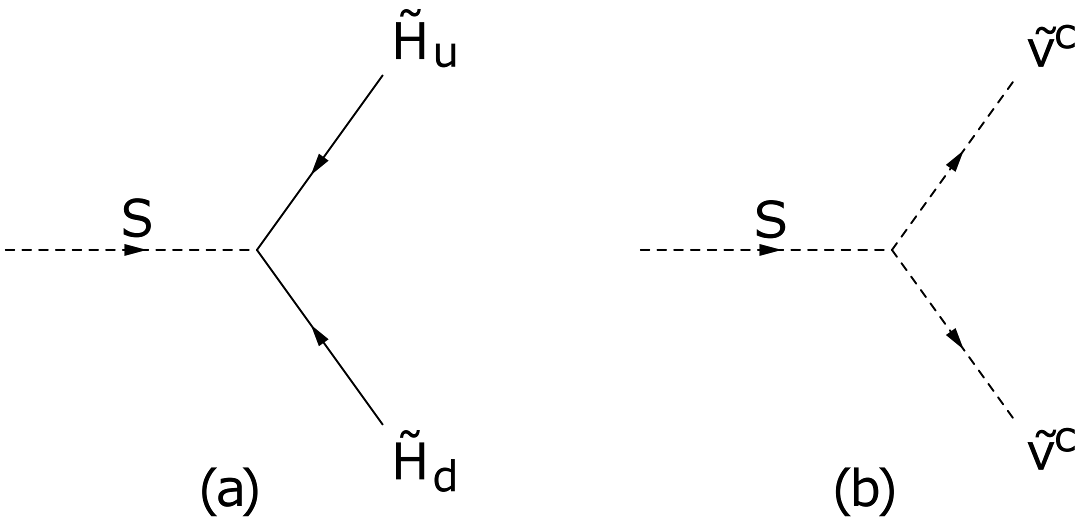

Figure 2: Decay of the bosonic component of S into MSSM Higgsinos ($\tilde{H}_u$, $\tilde{H}_d$) and right handed sneutrinos ($\tilde{\nu}^c$). The notation is as in Fig. 1.

- S could also decay into right handed sneutrinos via the F-term $F_{\bar{\Phi}}$ between $\kappa S\Phi\bar{\Phi}$ and $\bar{\Phi}\nu^c\nu^c$ after substituting $\langle\Phi\rangle$ – Fig. 2(b).
- The decay widths in the two cases are, respectively,

$$\Gamma_H^S = \frac{\lambda_\mu^2}{8\pi} m_S, \quad \Gamma_{\nu^c}^S = \frac{\lambda_{\nu^c}^2}{8\pi} m_S,$$

- Depending on the kinematics the total decay width of the resonance could easily lie in the multi-GeV range.
- The diphoton, dijet, diboson events in this case are sub-dominant.
- Our estimate of $|\langle S\rangle|$ holds if the decay of S into Higgsinos and right handed sneutrinos is kinematically blocked.
- This is achieved for $|\mu| = \lambda_\mu|\langle S\rangle| > m_S/2 \simeq 375$ GeV (or $\lambda_\mu \gtrsim 0.49$) and $\lambda_{\nu^c}M > m_S/2$.
- Demanding that the mass of the $B-L$ gauge boson $m_{Z'} = \sqrt{6}g_{B-L}M > 3$ TeV, we find that $g_{B-L}M \gtrsim 1225$ GeV.

- Setting $m_{3/2} = 50$ GeV, we obtain $\kappa \simeq 0.066$, $M \simeq 8040$ GeV, $\lambda_{\nu^c} \gtrsim 0.047$, and $g_{B-L} \gtrsim 0.15$.
- A gravitino in this mass range is a plausible cold matter candidate.
- $g_{B-L} \lesssim 0.25$, λ_D, and λ_L remain perturbative up to $M_{\rm GUT}$.
- If $g_{B-L} \simeq 0.24$, it unifies with the MSSM gauge couplings.
- So the requirements for a viable diphoton resonance are met.
- The spin zero field $\theta = (\theta_1 + i\theta_2)/\sqrt{2}$ consists of a (pseudo)scalar θ_1 (θ_2) field with mass $m_\theta = \sqrt{2}\kappa M$ in the SUSY limit.
- It couples to the scalar vector-like fields D, $\bar{D}$ via the F-term F_S between $\kappa S \Phi \bar{\Phi}$ and $\lambda_D S D \bar{D}$ with coupling constant $\lambda_D m_\theta$.

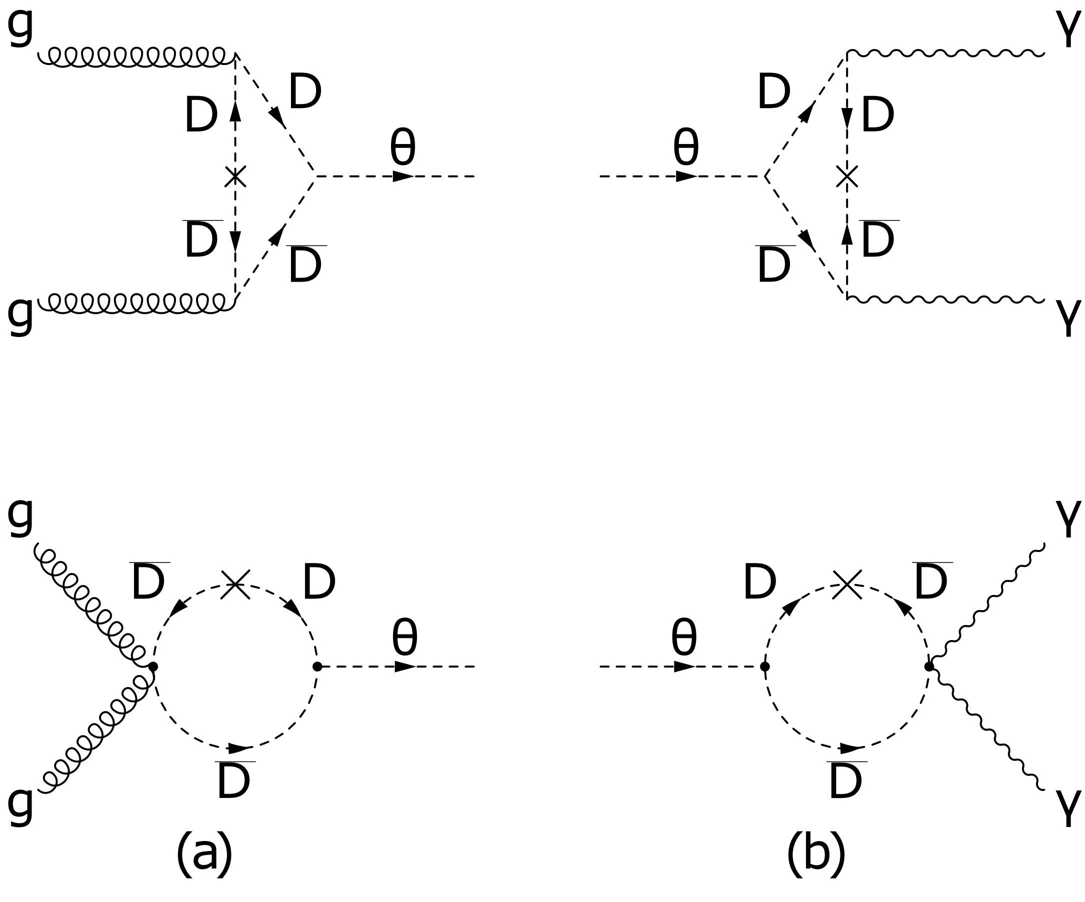

Figure 3: Production of the bosonic component of θ by g fusion and its subsequent decay into γ. The notation is as in Fig. 1 with the crosses indicating mass squared insertions.

- It also can be produced by g fusion via scalar D, $\bar{D}$ loops – Fig. 3(a), and decay into two γ via the graphs in Fig. 3(b).

- Note that, in the presence of L, $\bar{L}$, similar graphs with scalar L, $\bar{L}$ loops also contribute to the decay of θ into γ.

- The mass squared insertions in the graphs now arise from the soft SUSY breaking trilinear term $A'\lambda_D SD\bar{D}$ and are equal to $A'm_D$.

- Thus, for $A' \ll m_D$, the cross sections for the diphoton excess are suppressed by a factor $(A'/m_D)^4$ relative to the ones for S.

- Larger soft SUSY breaking trilinear terms will enhance the diagrams in Fig. 3 and also cause larger mixing between S and θ.

- In this case all four states can contribute to the diphoton excess.

- θ can decay, for exact SUSY, into H_u, H_d and ν^c's if this is kinematically allowed with decay widths equal to Γ_H^S, $\Gamma_{\nu^c}^S$ respectively.

5 Summary

- We presented a realistic $U(1)_{B-L}$ extension of the MSSM with resonances observable at the LHC and/or future colliders.

- The symmetries prevent the μ parameter and the masses of vector-like fields and a gauge singlet field from being arbitrarily large.

- Four spin zero resonances arise from a gauge singlet scalar and a pair of conjugate Higgs superfields responsible for $B-L$ breaking.

- One or more of them could explain the 750 GeV diphoton excess.

- Their total decay widths can lie in the multi-GeV range with the diphoton, diboson, dijet events being sub-dominant.