

Diphoton resonances in a $U(1)_{B-L}$ extension of the minimal supersymmetric standard model

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Inspired by the 750 GeV diphoton state recently reported by ATLAS and CMS, we propose a $U(1)_{B-L}$ extension of the MSSM which predicts the existence of four spin zero resonance states that are degenerate in mass in the supersymmetric limit. Vector-like fields, a gauge singlet field, as well as the MSSM Higgsinos are prevented from acquiring arbitrary large masses by a $U(1)$ R-symmetry. Indeed, these masses can be considerably lighter than the Z' gauge boson mass. Depending on kinematics the resonance states could decay into right handed neutrinos and sneutrinos, and/or MSSM Higgs fields and Higgsinos with total decay widths in the multi-GeV range.

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The recently reported [1] 750 GeV diphoton resonance by ATLAS and CMS, if confirmed during Run III at the LHC, would have far reaching ramifications in our quest for new physics beyond the Standard Model (SM). These preliminary results have triggered, not surprisingly perhaps, a flurry of theoretical papers [2] offering a large variety of plausible extensions of the SM in order to explain the reported diphoton excess.

In this paper we propose an extension of the Minimal Supersymmetric Standard Model (MSSM) which naturally yields resonance states in the TeV range. A simple implementation of this framework is realized in a local $U(1)_{B-L}$ extension of the MSSM gauge symmetry. In contrast to radiative electroweak breaking implemented in the MSSM, the additional local symmetry in our case is spontaneously broken at tree level with a superpotential W whose form is uniquely determined by a combination of the underlying gauge symmetry and a $U(1)$ global R-symmetry. The construction of W utilizes an appropriate pair of Higgs superfields ($\Phi, \bar{\Phi}$) as well as a gauge singlet superfield S . The resonance states arise from the scalar components of the $S - \Phi - \bar{\Phi}$ system, and their mass is determined, in the supersymmetric (SUSY) limit, by a dimensionless parameter in W which can be much smaller, if needed, compared to the typical order unity or so gauge coupling constant. Thus, if required, the resonance states can be significantly lighter than the mass of the Z' gauge boson associated with $B - L$. Note that with global $U(1)_{B-L}$ any constraint arising from Z' goes away. This is a plausible alternative to local $U(1)_{B-L}$ considered here.

The spontaneous breaking of $U(1)_{B-L}$ leaves SUSY unbroken. The symmetry breaking scale M may be much larger than the TeV SUSY breaking scale. Superpoten-

tials of this type have previously been employed by Dvali and Shafi [3] in their construction of SUSY trification models based on $SU(3)_c \times SU(3)_L \times SU(3)_R$, and later in the construction of SUSY hybrid inflation models [4, 5].

The scalar component of S acquires a non-zero vacuum expectation value (VEV) proportional to $m_{3/2}$ after SUSY breaking [3, 6], and this has been utilized in the past [6, 7] to resolve the MSSM μ problem. The R-symmetry also protects S from acquiring arbitrarily large masses. In the present scheme we also use this \sim TeV VEV of S to provide masses to suitable vector-like fields including colored fields which play a role in the production and subsequent decay of the scalar resonance(s).

The renormalizable superpotential of the MSSM with R-parity possesses three global symmetries, namely the baryon number $U(1)_B$, lepton number $U(1)_L$ and a R-symmetry $U(1)_R$, where, for simplicity, we ignore the tiny non-perturbative violation of B and L by the $SU(2)_L$ instantons. The new local $U(1)$ symmetry, which we identify as $U(1)_{B-L}$, is to be spontaneously broken at some scale M , and we prefer to implement this breaking by a SUSY generalization of the Higgs mechanism. Motivated by the MSSM example, we require that the new superpotential W respects the global $U(1)_B$ and $U(1)_L$ symmetries as well as a global $U(1)$ R-symmetry.

The full renormalizable superpotential is

$$\begin{aligned}
 W = & y_u H_u q u^c + y_d H_d q d^c + y_\nu H_\nu l \nu^c + y_e H_d l e^c \\
 & + \kappa S (\Phi \bar{\Phi} - M^2) + \lambda_\mu S H_u H_d + \lambda_{\nu^c} \bar{\Phi} \nu^c \nu^c \\
 & + \lambda_D S D \bar{D} + \lambda_q D q q + \lambda_{q^c} \bar{D} u^c d^c, \quad (1)
 \end{aligned}$$

where y_u, y_d, y_ν, y_e are the Yukawa coupling constants and the family indices are generally suppressed for simplicity. Here $q, u^c, d^c, l, \nu^c, e^c$ are the usual quark and lepton superfields of MSSM including the right handed neutrinos ν^c , and H_u, H_d are the standard electroweak Higgs superfields. The gauge singlet S has necessarily the same R-charge as W , which we take to be 2. Consequently, H_u, H_d have opposite R-charges, which can

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TABLE I: Superfield content of the model.

Superfields	Representations 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

be brought to zero by an appropriate hypercharge (Y) transformation. The R-charges of u^c (ν^c) and d^c (e^c) are equal and, thus, B and L transformations can make the R-charges of q , u^c , d^c , l , ν^c , e^c all equal to unity.

In order to determine the R-charges and $B - L$ quantum numbers of the SM singlets Φ , $\bar{\Phi}$, we introduce the coupling $\bar{\Phi}\nu^c\nu^c$, which implies that their $B - L$ charge is 2, -2 respectively, and their R-charges are zero. This coupling generates masses for the right handed neutrinos after the breaking of $U(1)_{B-L}$ to its Z_2 subgroup by the VEVs of Φ , $\bar{\Phi}$.

We also introduced the coupling $S\bar{D}\bar{D}$, where D (\bar{D}) are color triplet (antitriplet) and $SU(2)_L$ singlet superfields. (Color triplet vector-like superfields D , \bar{D} are perhaps best motivated in the framework of GUT symmetry E_6 .) To determine the charges of these fields we need at least one additional coupling. Taking the coupling Dqq , which is a color and $SU(2)_L$ singlet, the R-charges of D , \bar{D} vanish. Also, the Y charge of D becomes $-1/3$ and consequently, the hypercharge of \bar{D} is $1/3$. Finally, the $B - L$ charge of D is $-2/3$ and that of \bar{D} is $2/3$ with their lepton numbers vanishing. Note that the coupling $\bar{D}u^cd^c$ is also present since it respects all the symmetries of the model. It is also worth mentioning that the Z_2 subgroups of both $U(1)_R$ and $U(1)_{B-L}$ coincide with the Z_2 matter parity under which all the ordinary (anti)quark and (anti)lepton superfields are odd, with the rest of the superfields being even. This symmetry remains unbroken by all the soft SUSY breaking terms and all the VEVs. (See Ref. [8] for a more general discussion of unbroken Z_2 from $SO(10)$.) In Table I, we summarize all the superfields of the model together with their transformation properties under the SM gauge group ($G_{SM} = SU(3)_c \times SU(2)_L \times U(1)_Y$) and their charges under the global symmetries $U(1)_B$, $U(1)_L$, and $U(1)_R$.

The superpotential in Eq. (1) is the most general renormalizable superpotential which obeys the SM gauge sym-

metry and the global symmetries $U(1)_B$, $U(1)_L$, and $U(1)_R$. Had we removed the separate baryon and lepton number symmetries and kept only the gauge $U(1)_{B-L}$ symmetry, the renormalizable superpotential terms $\bar{D}ql$, Du^ce^c , and $Dd^c\nu^c$ would be present leading to fast proton decay and other baryon and lepton number violating effects [9]. The spontaneous breaking of $U(1)_{B-L}$ to its Z_2 subgroup will generate a network of local cosmic strings. Their string tension, which is determined by the scale M , is relatively small and certainly satisfies the most stringent relevant upper bound from pulsar timing arrays [10].

The ‘bare’ MSSM μ term is now replaced by a term SH_uH_d , so that the μ term is generated after S acquires a non-zero VEV of order TeV from soft SUSY breaking [6]. (In the SUSY limit the VEV of S is zero.) The VEV of S also plays an essential role, as we will see, in the generation of masses for the vector-like fields D , \bar{D} that are crucial in the production and decay of the diphoton resonance(s).

The spontaneous breaking of $U(1)_{B-L}$ implemented with the fields S , Φ , $\bar{\Phi}$ delivers, in the exact SUSY limit, four spin zero particles all with the same mass given by $\sqrt{2}\kappa M$. This mass, even for $M \gg 1$ TeV, can be of order TeV (more precisely $\simeq 750$ GeV in the present case) by selecting an appropriate value for κ . We should point out though that depending on the SUSY breaking mechanism, the four resonance states may end up with significantly different masses. The VEV of S , with suitable choice of the gauge and R-charges, yields masses for the appropriate fields that are vector-like under the MSSM gauge symmetry. This is in addition to possibly additional such fields that acquire masses from their coupling to the Higgs fields with VEV M that break the $U(1)_{B-L}$ gauge symmetry. The diquarks associated with the vector-like fields may be found [11] at the LHC.

The spontaneous breaking of $U(1)_{B-L}$ is achieved via the first term in the second line of Eq. (1), which gives the following potential for unbroken SUSY

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

Here we assumed that the mass parameter M and the dimensionless coupling constant κ are made real and positive by field rephasing, and the scalar components of the superfields are denoted by the same symbol. Vanishing of the D-terms implies that $|\Phi| = |\bar{\Phi}|$, which yields $\bar{\Phi}^* = e^{i\varphi}\Phi$, while the F-terms vanish for $S = 0$ and $\Phi\bar{\Phi} = M^2$, which requires that $\varphi = 0$. So, after rotating Φ and $\bar{\Phi}$ to the positive real axis by a $B - L$ transformation, we find that the SUSY vacuum lies at

$$S = 0 \quad \text{and} \quad \Phi = \bar{\Phi} = M. \quad (3)$$

The mass spectrum of the scalar $S - \Phi - \bar{\Phi}$ system is constructed by writing $\Phi = M + \delta\Phi$ and $\bar{\Phi} = M + \delta\bar{\Phi}$. In the unbroken SUSY limit, 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 squared splitting $\pm\sqrt{2}\kappa MA$ with the mass eigenstates now being $(S+\theta^*)/\sqrt{2}$ and $(S-\theta^*)/\sqrt{2}$. This splitting is small for $A \ll \sqrt{2}\kappa M$. Let us assume that the mixing is in general sub-dominant and ignore it. This simplifies our analysis.

The soft SUSY breaking terms

$$V_1 = A\kappa S\Phi\bar{\Phi} - (A - 2m_{3/2})\kappa M^2 S \quad (4)$$

in the potential with $m_{3/2}$ being the gravitino mass and $A \sim m_{3/2}$, which arise from the first term in the second line of Eq. (1), play an important role in our scheme. Here we assume minimal supergravity so that the coefficients of the trilinear and linear soft terms are related as shown. Substituting $\Phi = \bar{\Phi} = M$ in Eq. (4), we obtain a linear term in S which, together with the mass term $2\kappa^2 M^2 |S|^2$ of S , generates [3] a VEV for S :

$$\langle S \rangle = -\frac{m_{3/2}}{\kappa}. \quad (5)$$

Substituting this VEV of S in the superpotential term $\lambda_\mu SH_u H_d$, we obtain [6, 7] the MSSM μ term with $\mu = -\lambda_\mu m_{3/2}/\kappa$. The crucial point here is that the same VEV generates mass terms $-\lambda_D m_{3/2} D\bar{D}/\kappa$ for the vector-like superfields D, \bar{D} via the superpotential terms $\lambda_D S D\bar{D}$. The trilinear terms corresponding to these superpotential terms will produce mixing between the scalar components of D, \bar{D} . However, we will assume here that this mixing is sub-dominant.

To preserve gauge coupling unification one should introduce vector-like color singlet, $SU(2)_L$ doublet superfields L, \bar{L} equal in number to the color triplets D, \bar{D} . (With D, \bar{D} and L, \bar{L} masses \sim TeV, the gauge couplings stay in the perturbative domain for up to four such pairs.) These fields with a superpotential coupling $\lambda_L S L\bar{L}$ can enhance the branching ratio of the decay of the spin zero fields S and θ to photons and, in addition, allow the decay into W^\pm . Introducing the superpotential coupling Lle^c , the hypercharge of L (\bar{L}) is $-1/2$ ($1/2$). Their baryon, lepton numbers, and R-charges are all zero. These quantum numbers allow the superpotential couplings $SLH_u, SH_d\bar{L}, Lqd^c, \bar{L}qu^c$, and $\bar{L}l\nu^c$. Substituting $\langle S \rangle$ in $\lambda_L S L\bar{L}$, the superfields L, \bar{L} acquire a mass $m_L = -\lambda_L m_{3/2}/\kappa$.

The real scalar S_1 and real pseudoscalar S_2 components of S ($= (S_1 + iS_2)/\sqrt{2}$) with mass $m_S = \sqrt{2}\kappa M$ in the exact SUSY limit can be produced at the LHC by gluon fusion via a fermionic D, \bar{D} loop as indicated in Fig. 1. In the absence of the vector-like L, \bar{L} superfields, they can decay into gluons, photons, or Z gauge bosons via the same loop diagram, but not to W^\pm bosons since the D, \bar{D} are $SU(2)_L$ singlets. The most promising decay channel to search for these resonances is into two photons with the relevant diagram also shown in Fig. 1.

Applying the results of S.F. King and R. Nevzorov in

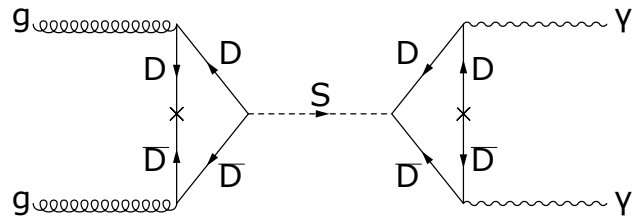


FIG. 1: Production of the bosonic component of S at the LHC by gluon (g) fusion and its subsequent decay into photons (γ). Solid (dashed) lines represent the fermionic (bosonic) component of the indicated superfields. The arrows depict the chirality of the superfields and the crosses are mass insertions which must be inserted in each of the lines in the loops.

Ref. [2], the cross section of the diphoton excess is

$$\sigma(pp \rightarrow S_i \rightarrow \gamma\gamma) \simeq \frac{C_{gg}}{m_S s \Gamma_{S_i}} \Gamma(S_i \rightarrow gg) \Gamma(S_i \rightarrow \gamma\gamma), \quad (6)$$

where $i = 1, 2$, $C_{gg} \simeq 3163$, $\sqrt{s} \simeq 13$ TeV, Γ_{S_i} is the total decay width of S_i , and the decay widths of S_i to two gluons (g) or two photons (γ) are given by

$$\Gamma(S_i \rightarrow gg) = \frac{n^2 \alpha_s^2 m_S^3}{256 \pi^3 \langle S \rangle^2} A_i^2(x), \quad (7)$$

$$\Gamma(S_i \rightarrow \gamma\gamma) = \frac{n^2 \alpha_Y^2 m_S^3 \cos^4 \theta_W}{4608 \pi^3 \langle S \rangle^2} A_i^2(x). \quad (8)$$

Here n is the number of D, \bar{D} pairs, which are taken, for simplicity, to have a common coupling constant λ_D to S , $A_1(x) = 2x[1 + (1-x)\arcsin^2(1/\sqrt{x})]$, $A_2(x) = 2x\arcsin^2(1/\sqrt{x})$, $x = 4m_D^2/m_S^2 > 1$, and α_s, α_Y are the strong and hypercharge fine-structure constants. If the L, \bar{L} superfields are present, they will also contribute to the decay width of S to photons via loop diagrams similar to the ones in the right part of Fig. 1, and Eq. (8) will be replaced by

$$\Gamma(S_i \rightarrow \gamma\gamma) = \frac{n^2 m_S^3 \alpha_Y^2 \cos^4 \theta_W}{4608 \pi^3 \langle S \rangle^2} A_i^2(x) \left[1 + \frac{3A_i(y)}{2A_i(x)} \left(1 + \frac{\alpha_2 \tan^2 \theta_W}{\alpha_Y} \right) \right]^2, \quad (9)$$

where α_2 is the $SU(2)_L$ fine-structure constant and $y = 4m_L^2/m_S^2 > 1$.

The cross section in Eq. (6) simplifies under the assumption that the spin zero fields S_i decay predominantly into gluons, namely, $\Gamma_{S_i} \simeq \Gamma(S_i \rightarrow gg)$. In this case, as pointed by R. Franceschini et al. in Ref. [2], one obtains $\sigma(pp \rightarrow S_i \rightarrow \gamma\gamma) \simeq 8$ fb if

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

For x and y just above unity, which guarantees that the decay of S_i to D, \bar{D} and L, \bar{L} pairs is kinematically

blocked, $A_1(x)$ and $A_2(y)$ are maximized with values $A_1 \simeq 2$ and $A_2 \simeq \pi^2/2$. Note that x close to unity means $m_D \simeq 375$ GeV. However, one should work with somewhat larger m_D as indicated by ATLAS and CMS. So we take $m_D = 700$ GeV. It is more beneficial to consider the decay of the pseudoscalar S_2 since $A_2(x) > A_1(x)$ for all $x > 1$. Using Eq. (9), we then find that the condition in Eq. (10) is satisfied for $nm_S/|\langle S \rangle| \simeq 2.97$. Therefore, for $n = 3$, we require that $m_S/|\langle S \rangle| \simeq 0.99$, which, for $m_S \simeq 750$ GeV, implies that $|\langle S \rangle| \simeq 758$ GeV. In this case, $\lambda_D \simeq 0.92$ and λ_L is just above 0.49. Comparing Eqs. (8) and (9), we find that the inclusion of the vector-like L, \bar{L} superfields enhances the decay width of S_2 to photons by about a factor 58.5.

In the exact SUSY limit, the complex scalar field S could decay into MSSM Higgsinos (potential dark matter candidate) via the superpotential term $\lambda_\mu SH_u H_d$ if this is kinematically allowed. S also could decay into right handed sneutrinos via the F-term $F_{\bar{\Phi}}$ between the superpotential terms $\kappa S \Phi \bar{\Phi}$ and $\bar{\Phi} \nu^c \nu^c$ after substituting the VEV of Φ . The decay widths in the two cases are

$$\Gamma_H^S = \frac{\lambda_\mu^2}{8\pi} m_S, \quad \Gamma_{\nu^c}^S = \frac{\lambda_{\nu^c}^2}{8\pi} m_S, \quad (11)$$

respectively, where we assumed that the masses of the Higgsinos and the relevant right handed sneutrinos are much smaller than m_S . Depending on the kinematics the total decay width of the resonance could easily lie in the multi-GeV range. The diphoton, dijet, and diboson decay modes in this case would be sub-dominant.

Our estimate of $m_S/|\langle S \rangle|$ after Eq. (10) requires that the decay widths of S into MSSM Higgsinos and right handed sneutrinos are sub-dominant or kinematically blocked. The latter 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 [12, 13], say, we find that $g_{B-L} M \gtrsim 1225$ GeV (g_{B-L} is the GUT normalized $B-L$ gauge coupling constant). From Eq. (5) and setting, say, $m_{3/2} = 50$ 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 – for a recent discussion and references, see Ref. [14]. Finally, we have checked that, in this example, $g_{B-L} \lesssim 0.25$, λ_D , and λ_L remain perturbative up to the GUT scale (M_{GUT}). In particular, if $g_{B-L} \simeq 0.24$, it unifies with the MSSM gauge coupling constants at M_{GUT} . So the requirements for a viable diphoton resonance are met.

The complex spin zero field $\theta = (\delta\Phi + \delta\bar{\Phi})/\sqrt{2} = (\theta_1 + i\theta_2)/\sqrt{2}$, which consists of a real scalar (θ_1) and a real pseudoscalar (θ_2) field and has mass $m_\theta = \sqrt{2}\kappa M$ in the SUSY limit, couples to the scalar vector-like fields D, \bar{D} via the F-term F_S between the superpotential terms $\kappa S \Phi \bar{\Phi}$ and $\lambda_D S D \bar{D}$. The coupling constant is $\lambda_D m_\theta$. It also can be produced at the LHC by gluon fusion via scalar D, \bar{D} loops as shown in Fig. 2(a), and decay into two photons via the diagrams in Fig. 2(b). (In the presence of L, \bar{L} superfields similar diagrams with scalar $L,$

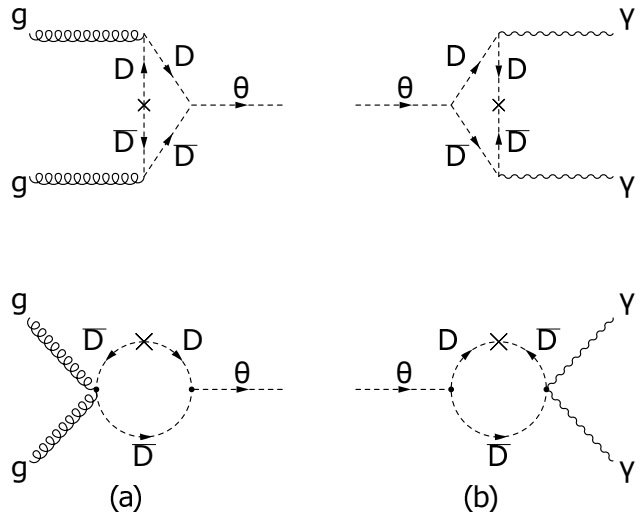


FIG. 2: Production of the bosonic component of θ at the LHC by gluon (g) fusion and its subsequent decay into photons (γ). The notation is the same as in Fig. 1 with the crosses indicating mass squared insertions.

\bar{L} loops also contribute to the decay of θ into photons.) The important point here is that the mass squared insertions in all the diagrams of Fig. 2 arise from the soft SUSY breaking trilinear term $A' \lambda_D S D \bar{D}$ and are thus equal to $A' m_D$, where $m_D = -\lambda_D m_{3/2}/\kappa$ is the mass of the D, \bar{D} superfields generated by the VEV of S in Eq. (5). Consequently, 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 the spin zero field S . Larger soft SUSY breaking trilinear terms will enhance the diagrams in Fig. 2 and also cause larger mixing between S and θ . In this case all four spin zero states can contribute to the diphoton excess.

The field θ can decay, in the exact SUSY limit, into MSSM Higgs fields via the F-term F_S between the superpotential terms $\kappa S \Phi \bar{\Phi}$ and $\lambda_\mu S H_u H_d$ if this is kinematically allowed. The relevant coupling constant is $\lambda_\mu m_\theta$, and thus the decay width is the same as Γ_H^S in Eq. (11) provided that the masses of the Higgs fields are much smaller than m_θ (for $m_S = m_\theta$). θ also could decay into right handed neutrinos via the superpotential term $\lambda_{\nu^c} \bar{\Phi} \nu^c \nu^c$ with a decay width equal to $\Gamma_{\nu^c}^S$ in Eq. (11) under the same assumption.

In conclusion, we have presented a realistic SUSY model based on a $U(1)_{B-L}$ extension of the MSSM that contains resonances observable at the LHC and/or future colliders. The underlying gauge and R-symmetries are such that the MSSM μ parameter and the masses of vector-like superfields and a gauge singlet superfield cannot be arbitrarily large. A resonance system consisting of four spin zero states arises from a gauge singlet scalar and a pair of conjugate Higgs superfields responsible for the $B-L$ breaking. These states are degenerate in mass in the SUSY limit, and depending on the details of SUSY

breaking, one or more of these states could explain the observed 750 GeV diphoton excess. Their total decay widths can lie in the multi-GeV range, depending on the kinematics, in which case the diphoton, diboson and dijet events will be sub-dominant.

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