Diphoton resonance at 750 GeV in the broken R-symmetric MSSM

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Nonobservation of superpartners of the Standard Model particles at the early runs of the LHC provide strong motivation for introducing an R-symmetric minimal supersymmetric Standard Model. This model also comes with a pair of extra scalars which couple only to superpartners at the tree level. We demonstrate that in the case when the $U(1)_R$ symmetry is mildly broken one of these scalars develops all the properties necessary to explain the 750 GeV diphoton resonance recently observed at the LHC, as well as the nonobservation of associated signals in other channels. Some confirmatory tests in the upcoming LHC runs are proposed.

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I. INTRODUCTION

In many ways, supersymmetric models remain the best option for new physics beyond the Standard Model (SM) of electroweak and strong interactions. The discovery of a light, probably elementary, scalar in 2012 [1] has made this motivation, if anything, stronger than ever. However, the decay modes of the 125 GeV scalar found in 2012 appear increasingly to resemble those of the SM Higgs boson [2]. Moreover, this must be coupled with the somewhat disappointing fact that the early runs of the LHC at CERN, Geneva, have not found any of the promised signals for supersymmetry (SUSY) [3,4]. Though all this does not invalidate the *idea* of SUSY *per se*, it has made it increasingly difficult to fit the observed results with popular models of SUSY, such as the so-called minimal supersymmetric SM (MSSM).

To add to this tension, we have the recent announcement that both the ATLAS and CMS collaborations seem to have observed [5,6] an excess of diphoton events in the 13 TeV run, which commenced last year. The excess events appear to arise from a resonant production of an intermediate particle of mass around 750 GeV and a width which is best fit as 45 GeV. At the same time, both the experimental collaborations have announced that searches for deviations from the SM prediction in all other channels have produced null results. Their principal results on the diphoton excess are summarized below:

(i) The ATLAS Collaboration has observed [5] an excess of 14 events, with a peak at 750 GeV and a best-fit width of 45 GeV, in 3.2 fb^{-1} of data at $\sqrt{s} = 13$ TeV. The local significance of this excess is 3.9σ, but it reduces to about 2.6σ if the look-elsewhere effect is included. Taking into account the experimental acceptance value of about 0.4, this corresponds to an excess signal of 10 ± 3 fb.

(ii) The CMS Collaboration has observed [6] an excess of 13 events, with a peak at 760 GeV, in 2.6 fb^{-1} of data at $\sqrt{s} = 13$ TeV. The local significance of this excess is 2.3σ, but it reduces to about 2.0σ in a global fit. Taking into account the experimental acceptance value of about 40%, this corresponds to an excess signal of 6 ± 3 fb.

Subsequently, both ATLAS and CMS have updated their results [7]; however, these only lead to very minor changes in the theoretical cross sections and are not considered further in our discussion.

While there is a strong probability that this excess is only a statistical fluctuation in the data, there is always the more exciting possibility that this may be the first observed manifestation of new physics at the LHC—or, for that matter, any other collider experiment. Undoubtedly, this announcement has stirred the theoretical mind, for several new physics interpretations of this excess have already appeared in the literature. For example, models with vectorlike fermions and extended scalar sectors [8–15], SUSY [8,9,16], extra dimensions [12,17], axions and composite scalars [18,19], vector resonances [20], leptoquarks [21], dark matter candidates [14,22], and minimal gauge extensions of the SM and MSSM [23] have been studied. Some have proposed model-independent tests of the signal [19,24], and others have constructed scenarios in which the presence of a diphoton excess and the absence of any other signals arise in a natural way [25]. In addition, electroweak vacuum stability and inflation in the presence of this new resonance has been analyzed in Ref. [26]. However, it is probably a fair statement to say that an explanation of the current results is rather difficult to obtain in any of the popular "minimal" models which have hitherto been the mainstay of phenomenological studies of physics beyond the SM. Quite naturally, therefore, many of the proposed scenarios invoke exotic options, which are

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barely permitted by the experimental data and do not conform to the choices commonly seen earlier in the literature [27]. It is interesting, therefore, to ask if a well-motivated model, where a specific scenario in the parameter choices could explain the observed facts in this regard, can be found.

It is now common knowledge [8,9,16] that general SUSY models in their minimal incarnation fail to explain the diphoton excess. Hence, it is relevant to ask if the effect can be explained by extending our consideration to *R*-symmetric SUSY models. In this article, therefore, we consider the minimal *R*-symmetric supersymmetric SM (MRSSM) [28,29], which—apart from having the usual virtues of a SUSY model—provides a ready explanation for the nonobservance of SUSY-specific signals at the LHC. This is because the MRSSM permits us to have a spectrum where the gluinos are of Dirac nature and hence can be made very heavy (of the order of a few TeV) without making the squarks, especially of the third generation, correspondingly heavy. Thus, it is possible to have, in the MRSSM, stops and/or sbottoms light enough to be compatible with Higgs boson self-energy corrections and, at the same time, ensure that direct production cross sections for all these sparticles remain small enough to evade all LHC (and other) constraints. More details are given in the next section.

The inclusion of the *R* symmetry also has the effect of making the μ term vanish, which would make the Higgsinos of the theory massless in the case of unbroken electroweak symmetry. To create the μ term, it becomes necessary to incorporate two additional $SU(2)$-doublet chiral superfields $\hat{R}_u$ and $\hat{R}_d$ carrying nonzero *R* charges. This precludes them from coupling with most SM fields (except the Higgs sector), and hence they are dubbed "inert" doublets. It is one of the neutral components of these inert doublets that we propose as a candidate for the 750 GeV resonance.

The plan of this paper is as follows. In Sec. II, we describe the *R*-symmetric version of the MSSM, illustrating the role of the inert doublets mentioned above. We then go on, in Sec. III, to explain how *R* symmetry requires being broken in order to obtain a left-right mixing in the top-squark sector, which is vital to get a diphoton signal. Section IV is devoted to the details of how the diphoton excess arises in this model. In Sec. V, we summarize our results and mention some tests which may falsify this scenario in future runs of the LHC.

II. MRSSM FRAMEWORK

In *R*-symmetric models, one adds to the symmetries of the SM an extra $U(1)_R$ global symmetry, under which the superpartner fields transform but the SM fields do not. Among other things, this *R* symmetry prohibits the following terms which usually appear in the MSSM:

(i) Majorana gaugino mass terms: $m\bar{\lambda}^c\lambda$;

(ii) SUSY-breaking trilinear scalar couplings, $A_{ijk}\varphi_i\varphi_j\varphi_k$, where the φ generically stand for scalar fields in the theory;

(iii) the μ term: $\mu\hat{H}_u\hat{H}_d$.

Since gauginos cannot be massless, in the MRSSM, they will have to be Dirac fermions—to construct which one needs to introduce additional superfields $\hat{S}$, $\hat{T}$, and $\hat{\mathcal{O}}$ (with appropriate quantum numbers) to partner with the Bino, the Winos and the gluinos, respectively. We can easily make these Dirac masses large—of the order of several TeV—without running into problems. In the MSSM, where gauginos have Majorana mass terms, these Majorana masses appear on the right side of the renormalization group evolution equation for squarks, i.e.,

$$\frac{dm_{\tilde{q}}^2}{dt} = cM_{\tilde{g}}^2 + \cdots, \qquad (1)$$

where c is an appropriate coefficient, $M_{\tilde{g}}$ is the Majorana mass of the gluino, and $t = \ln Q^2/\mu^2$ as usual. Raising this Majorana mass to a few TeV would therefore push the squark masses also to a few TeV, and one would not then be able to get light stops and sbottoms. However, if the mass is of Dirac type, it does not appear at all on the right side of Eq. (1), and hence there is no obstacle to pushing the Dirac mass up to a few TeV [30,31] in the MRSSM.

An immediate consequence of this is that squarks coupling to quarks and a Dirac gluino have much lower production cross sections [31,32] at the LHC than they would have had in the usual case of a Majorana gluino. This is partly because of the large mass in the propagator but also because fermion-number-violating processes (e.g., $pp \to \tilde{q} + \tilde{q}$) are forbidden in the absence of exchanged Majorana fermions. This significantly weakens the rather tight constraints which have already been obtained at the LHC on light squarks in the MSSM. It may be noted in passing that multiple versions of models with Dirac gauginos can be found in the literature [30,33], where flavor and *CP*-violation constraints are suppressed [34] and issues pertaining to neutrino mass generation and dark matter can also be addressed [35,36].

The SM gauge quantum numbers and $U(1)_R$ charges of all the chiral superfields in the MRSSM are shown in Table I, where the MSSM superfields are listed on the left, while the additional superfields in the MRSSM are listed on the right. As mentioned in the Introduction, a μ term can be generated by adding two new superfields $\hat{R}_u$ and $\hat{R}_d$ carrying nonzero $U(1)_R$ charges. We have noted that this is enough to ensure nonzero Higgsino masses or, equivalently, at the electroweak scale, to make the lightest chargino $\tilde{\chi}_1^\pm$ heavy enough to escape the LEP bounds. We note that the scalars R_u and R_d have the same $U(1)_R$ charge as the superfields $\hat{R}_u$ and $\hat{R}_d$, whereas the $U(1)_R$ charges of the fermions $\tilde{R}_u$ and $\tilde{R}_d$ are less by one unit. To have an invariant action, the superpotential has to have $U(1)_R$

TABLE I. The chiral superfields in the MRSSM, showing their gauge quantum numbers under the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ as well as their $U(1)_R$ charge assignments.

Superfields	SM rep	$U(1)_R$	Superfields	SM rep	$U(1)_R$
$\hat{H}_u$	(1, 2, 1)	0	$\hat{R}_u$	(1, 2, −1)	2
$\hat{H}_d$	(1, 2, −1)	0	$\hat{R}_d$	(1,2,1)	2
$\hat{Q}_i$	$(3, 2, \frac{1}{3})$	1	$\hat{S}$	(1,1,0)	0
$\hat{U}_i^c$	$(\bar{3}, 1, -\frac{4}{3})$	1	$\hat{T}$	(1,3,0)	0
$\hat{D}_i^c$	$(\bar{3}, 1, \frac{2}{3})$	1	$\hat{O}$	(8,1,0)	0
$\hat{L}_i$	(1, 2, −1)	1			
$\hat{E}_i^c$	(1, 1, 2)	1			

charge of two units. This superpotential can now be written as

$$W = \mu_d \hat{R}_d \hat{H}_d + \mu_u \hat{R}_u \hat{H}_u + \Lambda_d \hat{R}_d \hat{T} \hat{H}_d + \Lambda_u \hat{R}_u \hat{T} \hat{H}_u + \lambda_d \hat{S} \hat{R}_d \hat{H}_d + \lambda_u \hat{S} \hat{R}_u \hat{H}_u + Y_d \hat{Q}_i \hat{H}_d \hat{D}_i^c + Y_e \hat{E}_i^c \hat{L}_i \hat{H}_d + Y_u \hat{U}_i^c \hat{Q}_i \hat{H}_u. \quad (2)$$

where the μ's, Λ's, λ's, and Y's are real constants.

In the most general case, this model has a Higgs sector extended beyond the MSSM by the neutral scalar components S and T of the superfields $\hat{S}$ and $\hat{T}$, as well as an extended fermionic sector. Thus, the neutral scalars in the model will result from the mixing of (H_u, H_d, S, T). The neutral part of the scalar potential, consisting of F terms, D terms, and explicit soft SUSY-breaking terms, can be written as

$$V_{\rm neut} = (m_{H_d}^2 + \mu_d^2)|H_d^0|^2 + (m_{H_u}^2 + \mu_u^2)|H_u^0|^2 + \frac{1}{8}(g^2 + g'^2)(|H_d^0|^2 - |H_u^0|^2 - |R_d^0|^2 + |R_u^0|^2)^2 + (m_{R_u}^2 + \mu_u^2)|R_u^0|^2 + (m_{R_d}^2 + \mu_d^2)|R_d^0|^2 + (B\mu H_d^0 H_u^0 + {\rm H.c.}) + |\lambda_u R_u^0 H_u^0 - \lambda_d R_d^0 H_d^0|^2 + \left|\frac{\Lambda_d}{\sqrt{2}} R_d^0 H_d^0 + \frac{\Lambda_u}{\sqrt{2}} R_u^0 H_u^0\right|^2 + \cdots, \quad (3)$$

with the ellipsis indicating additional terms involving the S and T fields. This potential could now be minimized to find the scalar eigenstates of the model. However, if we further assume that the S and T scalars acquire negligibly small vacuum expectation values (vevs) and hence mix very mildly with the Higgs sector,[1] then the minimization condition automatically pushes the S and T masses to the decoupling limit, and we are left with a MSSM-like pair of Higgs doublets $H_{u,d}$, in addition to the inert doublets $R_{u,d}$—as shown explicitly in Eq. (3). Moreover, in this limit, we recover the MSSM prediction for the tree-level mass of the light Higgs scalar, i.e.,

$$M_h^2 \simeq M_Z^2 \cos^2 2\beta. \quad (4)$$

The R-charge assignments ensure that, even after electroweak symmetry breaking, the R_u^0 and R_d^0 scalars do not mix with the $H_{u,d}^0$, though they may mix with each other. Of course, R-scalar couplings to pairs of any SM particle vanish. Moreover, the R-charge assignments of these R scalars restrict their trilinear couplings to:

(a) sfermions and chargino/neutralino combinations, e.g., $R\tilde{\ell}\,\tilde{\ell}$, $R\tilde{q}\,\tilde{q}$, $R\tilde{\chi}_i^+\tilde{\chi}_j^-$, $R\tilde{\chi}_i^0\tilde{\chi}_j^0$,

(b) paired-R scalars to SM bosons, i.e., RRH and RRV, where $V = W^\pm, Z^0$.

The trilinear $R_{u,d}$-sfermion-antisfermion couplings which play a major role in our work are

$$\mathcal{L}_{\rm int} = -\mu_u Y_u R_u^0 \tilde{u}_R \tilde{u}_L^* - \mu_d Y_d R_d^0 \tilde{d}_R \tilde{d}_L^* - \mu_d Y_e R_d^0 \tilde{\ell}_R \tilde{\ell}_L^* + {\rm H.c.}, \quad (5)$$

where Y_u, Y_d, and Y_e are Yukawa couplings of the SM and a sum over generations is implicit. For third-generation quarks, we have $Y_t \gg Y_b$, and therefore we will mostly confine ourselves to the R_u^0 scalar.

One may naively think that if the R_u^0 scalar is to be a candidate for the 750 GeV resonance then it can couple to photon/gluon pairs through a one-loop diagram involving stops (staus) in the loop (see Fig. 1). However, it is important to note that in this R-conserving scenario the R_u^0 scalar couples only to $\tilde{q}_L - \tilde{q}_R^*$ pairs and not to $\tilde{q}_L - \tilde{q}_L^*$ or $\tilde{q}_R - \tilde{q}_R^*$ pairs. As a result, one cannot generate R_u^0 couplings to photon/gluon pairs through stops (staus) loops (see Fig. 1) unless there is a mixing between the $\tilde{t}_L$ and $\tilde{t}_R$ states (or $\tilde{\tau}_L$ and $\tilde{\tau}_R$). It is clear, therefore, that if we are to interpret the 750 GeV scalar as a resonant R_u^0 decaying to diphotons we must break the R symmetry somehow, even if the breaking is a small effect.

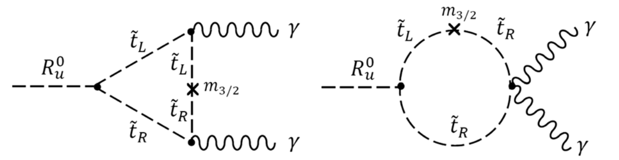

FIG. 1. Typical stop loops contributing to $R_u^0 \to \gamma\gamma$. Similar diagrams mediate gluon-gluon fusion $gg \to R_u^0$, where the photon lines are replaced by gluon lines.

III. R-SYMMETRY BREAKING

Apart from the above, there already exist fairly strong motivations for the spontaneous breaking of R symmetry from cosmological considerations [37]. We assume, therefore, that this happens in the hidden sector through some appropriate mechanism, leading to the gravitino acquiring a

[1] There is no *a priori* reason to assume this. However, a large vev of these scalars will be constrained from the Higgs boson signal strengths at the LHC.

nonzero mass $m_{3/2} = \Lambda^2/M_P$, where Λ is the SUSY-breaking scale and M_P is the Planck scale. This R-breaking information then gets communicated to the visible sector by anomaly mediation. As a result, the R-symmetry breaking sector in the Lagrangian can be written

$$\mathcal{L}_{\not R} = A_u \tilde{u}_R \tilde{u}_L^* H_u^0 + A_d \tilde{d}_R \tilde{d}_L^* H_d^0 + A_l \tilde{l}_R \tilde{l}_L^* H_d^0 + \text{H.c.} + M_1' \tilde{B}\,\tilde{B} + M_2' \tilde{W}\,\tilde{W} + M_3' \tilde{g}\,\tilde{g}, \tag{6}$$

where the gaugino masses (in addition to the large Dirac mass parameters M_i) get small R-breaking contributions M_i',

$$M_i' = \frac{g_i^2}{16\pi^2} b_i m_{3/2} \qquad (i = 1, 2, 3), \tag{7}$$

with beta functions

$$b_1 = 33/5 \qquad b_2 = 1 \qquad b_3 = -3. \tag{8}$$

In this case, we also generate R-symmetry-breaking trilinear terms,

$$\begin{aligned} A_t &= \frac{m_{3/2}}{16\pi^2}\left(-\frac{13}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + 6Y_t^2 + Y_b^2\right) \\ A_b &= \frac{m_{3/2}}{16\pi^2}\left(-\frac{7}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + Y_t^2 + 6Y_b^2 + Y_\tau^2\right) \\ A_\tau &= \frac{m_{3/2}}{16\pi^2}\left(-\frac{9}{5}g_1^2 - 3g_2^2 + 3Y_b^2 + 4Y_\tau^2\right), \end{aligned} \tag{9}$$

with the others being negligible.

The gravitino mass $m_{3/2}$—which sets the scale for the R-symmetry breaking—can be taken to be very small; e.g., a mass in the ballpark of hundreds of keV to a few MeV is permitted by experimental data [38]. In this case, most of the R-symmetry-breaking effects will also be small. The largest of these is the A_t term, which generates a mixing between the left- and the right-chiral stops after electroweak symmetry breaking. The mass-squared matrix for the stops $\tilde{t}_L$ and $\tilde{t}_R$ takes the form

$$\mathcal{M}_{\tilde{t}}^2 = \begin{bmatrix} (\mathcal{M}_{\tilde{t}}^2)_{11} & (\mathcal{M}_{\tilde{t}}^2)_{12} \\ (\mathcal{M}_{\tilde{t}}^2)_{21} & (\mathcal{M}_{\tilde{t}}^2)_{22} \end{bmatrix}, \tag{10}$$

where

$$\begin{aligned} (\mathcal{M}_{\tilde{t}}^2)_{11} &= \frac{1}{8}\left(g^2 + \frac{g'^2}{3}\right)(v_d^2 - v_u^2) + m_{\tilde{t}_L}^2 + \frac{1}{2}Y_t^2 v_u^2, \\ (\mathcal{M}_{\tilde{t}}^2)_{12} &= (\mathcal{M}_{\tilde{t}}^2)_{21} = A_t v_u, \\ (\mathcal{M}_{\tilde{t}}^2)_{22} &= \frac{g'^2}{6}(v_d^2 - v_u^2) + m_{\tilde{t}_R}^2 + \frac{1}{2}Y_t^2 v_u^2, \end{aligned} \tag{11}$$

in terms of the vevs v_u and v_d of the two Higgs doublets H_u and H_d, respectively. For the physical states $\tilde{t}_1$ and $\tilde{t}_2$, the mixing angle $\theta_{\tilde{t}}$ is now given by

$$\tan 2\theta_{\tilde{t}} = \frac{2 v_u A_t}{\frac{1}{8}(v_d^2 - v_u^2)(g^2 - g'^2) + (m_{\tilde{t}_L}^2 - m_{\tilde{t}_R}^2)}. \tag{12}$$

A similar expression occurs in the MSSM (with an additional μ term in the numerator), where a large value of A_t ensures a large mixing angle. In the present model, however, A_t is proportional to $m_{3/2}$ and is hence very small, but one can still obtain *maximal* mixing if the denominator of Eq. (12) can be arranged to vanish; i.e., we tune the parameters such that

$$m_{\tilde{t}_L}^2 = m_{\tilde{t}_R}^2 + \frac{1}{8}(v_u^2 - v_d^2)(g^2 - g'^2). \tag{13}$$

This large mixing between $\tilde{t}_L$ and $\tilde{t}_R$ is crucial for our analysis, since it permits the R_u^0 scalar to decay to diphotons and to be produced through gluon fusion by stop loops at the required rates.

It may be noted that the second term on the right of Eq. (13) is small, and hence we must have $m_{\tilde{t}_L} \simeq m_{\tilde{t}_R}$. In turn, if we plug this into Eq. (11) and note that the off-diagonal terms are small, it also follows that $m_{\tilde{t}_1} \simeq m_{\tilde{t}_2}$. This is unlike the MSSM, where a large splitting between stop masses is demanded by the Higgs mass constraint. In the MRSSM, however, additional contributions to the Higgs boson mass can be obtained [29] from loops involving the exchange of scalar partners of the S and T fields, rendering nearly degenerate stops a perfectly viable scenario.

We also note in passing that the breaking of R symmetry creates some small mixings between the H_u^0 and H_d^0 with the R_u^0 and R_d^0. However, as these will be generically proportional to the ratio $m_{3/2}/v \sim 10^{-6}$ ($m_{3/2} \sim 0.1$ MeV), they will not have any observable effects on the phenomenology of the Higgs boson sector.

IV. FITTING THE DIPHOTON SIGNAL

We are now in a position to study the observed diphoton signal as an R_u^0 resonance in the MRSSM. The dominant one-loop contributions to the decay of the R_u^0 scalar to a $\gamma\gamma$ pair come from the diagrams shown in Fig. 1, as well as a crossed diagram. Similar diagrams would also mediate its decay to a gg pair.

Evaluation of these diagrams, with both $\tilde{t}_1$ and $\tilde{t}_2$ running in the loops, leads to the partial widths

$$\Gamma(R_u^0 \to \gamma\gamma) \simeq \frac{\alpha^2 N_c^2 Q_{\tilde{t}}^4}{1024\pi^3} \sum_{i=1}^{2} \frac{M_R^3 \mu_{\rm eff}^2}{m_{\tilde{t}_i}^4} |F(\tau_i)|^2$$
$$\Gamma(R_u^0 \to gg) \simeq \frac{\alpha_s^2}{512\pi^3} \sum_{i=1}^{2} \frac{M_R^3 \mu_{\rm eff}^2}{m_{\tilde{t}_i}^4} |F(\tau_i)|^2, \qquad (14)$$

where N_c is the color factor, $Q_{\tilde{t}} = 2/3$ is the fractional charge of the stops, and M_R is the mass of the R_u^0 scalar. In the above formulas, $\mu_{\rm eff}$ is an effective coupling defined as

$$\mu_{\rm eff} = \frac{1}{4} \mu_u Y_t \sin^2 2\theta_{\tilde{t}}, \qquad (15)$$

and $F(\tau)$ is the loop integral function

$$F(\tau) = \left(\tau \sin^{-1} \frac{1}{\sqrt{\tau}}\right)^2 - \tau, \qquad (16)$$

where $\tau_i = 4m_{\tilde{t}_i}^2 / M_R^2$. This particular form of $F(\tau)$ arises only in the case $2m_{\tilde{t}} > M_R$—which is anyway assumed by us to preclude the decay $R_u^0 \to \tilde{t}\tilde{t}^*$.

We now compare the predictions of this model with the experimental results quoted in the Introduction. It is important to note, at this stage itself, that we must assume that the sparticle mass spectrum is such that all tree-level decays of the R_u^0 scalar (which occur through the trilinear couplings described in Sec. II) are kinematically disallowed. This is indicated by (a) the nonobservability of all decay channels other than the diphoton and (b) the requirement that the diphoton branching ratio should be large enough to yield the observed event rate at the LHC. We have already seen that such a sparticle mass spectrum does not conflict with any known theoretical or experimental requirements.

It is most convenient to treat the two widths $\Gamma_{\gamma\gamma} = \Gamma(R_u^0 \to \gamma\gamma)$ and $\Gamma_{gg} = \Gamma(R_u^0 \to gg)$ as correlated variables and study the plane formed by plotting them against each other. The production cross section for the R_u^0 scalar will be given in terms of Γ_{gg} and $\Gamma_{\gamma\gamma}$ by

$$\sigma_R = \frac{\pi^2}{8 s M_R} \Gamma_{gg} C_{gg} K_{gg} + \frac{8\pi^2}{s M_R} \Gamma_{\gamma\gamma} C_{\gamma\gamma} K_{\gamma\gamma}, \qquad (17)$$

where C_{gg} and $C_{\gamma\gamma}$ are the gluon and photon fluxes given by

$$C_{gg} = \int_\delta^1 \frac{dx}{x} f_{g/p}(x) f_{g/p}\left(\frac{\delta}{x}\right)$$
$$C_{\gamma\gamma} = \int_\delta^1 \frac{dx}{x} f_{\gamma/p}(x) f_{\gamma/p}\left(\frac{\delta}{x}\right), \qquad (18)$$

with $\delta = M_R^2/s$, where $\sqrt{s} = 13$ TeV and M_R is set to 750 GeV. The functions $f_{g/p}(x)$ and $f_{\gamma/p}(x)$ are, of course, the gluonic and photonic parton-density functions, respectively. K_{gg} is a QCD correction factor, which we take to be approximately 1.5 [39], while $K_{\gamma\gamma}$ is set to unity. Using the CTEQ-6 [40] set of structure functions, we evaluate $C_{gg} \approx 2914$, and we take $C_{\gamma\gamma} \approx 11$ [19]. It follows that the production cross section is

$$\sigma_R \approx 12.4 \text{ nb} \times \frac{\Gamma_{gg}}{M_R} + 2.0 \text{ nb} \times \frac{\Gamma_{\gamma\gamma}}{M_R}. \qquad (19)$$

We now have the cross section for the following:

(i) Diphotons, given by

$$\sigma_{\gamma\gamma} = \sigma_R \frac{\Gamma_{\gamma\gamma}}{\Gamma_R}, \qquad (20)$$

where $\Gamma_R = \Gamma_{gg} + \Gamma_{\gamma\gamma}$ is the total width of the R_u^0 resonance, assuming that no other decay modes are available to the R_u^0 scalar—which will be the case if $2M_{\tilde{t}} > M_R$, as assumed.

(ii) Dijets, given by

$$\sigma_{gg} = \sigma_R \frac{\Gamma_{gg}}{\Gamma_R}; \qquad (21)$$

as the R_u^0 scalar has no coupling with quarks and it is lighter than all the squark pairs, we can safely assume that the decay of a R_u^0 to dijets is completely dominated by the gg mode.

Our analysis is then based on the following constraints:

(iii) The total width Γ_R of the R_u^0 scalar should satisfy

$$\Gamma_R < 50 \text{ GeV}. \qquad (22)$$

Since the best-fit width is about 45 GeV, the value 50 GeV chosen above seems to provide a reasonable leeway for errors.

(iv) The dijet cross section observed at the LHC in the 13 TeV run is consistent with the SM prediction of about 12.5 ± 1.2 pb [41]. Thus, we must demand that the dijet excess arising from decay of the R_u^0 satisfies

$$\sigma_{gg} < 2.5 \text{ pb} \qquad (23)$$

assuming agreement with the SM at the 95% confidence level.

(v) The diphoton excess must be consistent with the observed values as presented by the ATLAS and CMS collaborations (see the Introduction). If we consider the 95% confidence level, the ATLAS results require

$$4 \text{ fb} < \sigma_{\gamma\gamma} < 16 \text{ fb}, \qquad (24)$$

and the CMS results may be taken to require

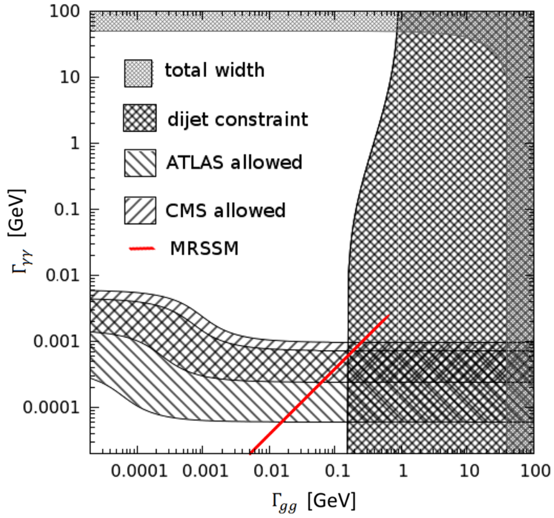

FIG. 2. Illustrating allowed regions as well as constraints on the $\Gamma_{\gamma\gamma}$-Γ_{gg} plane from the experimental data alongside the predictions of the MRSSM model (red line).

$$1 \text{ fb} < \sigma_{\gamma\gamma} < 12 \text{ fb}. \tag{25}$$

Combining all these constraints, we display our results in Fig. 2, which shows the $\Gamma_{\gamma\gamma}$-Γ_{gg} plane for a wide range of values from 10^{-5} to 10^2. The dark gray-shaded strip along the top and right edges of this panel represents the range ruled out by the total width constraint in Eq. (22). The larger hatched region on the right side of the panel represents the dijet constraint in Eq. (23); i.e., all points in the region would have lead to an observable dijet signal at the 13 TeV run. The overlapping more-or-less horizontal strips with parallel hatching depict the values *allowed* by the ATLAS and CMS observations (see the key on the figure), with the overlap region appearing cross-hatched. Finally, the solid (red) oblique line close to the lower end of Fig. 2 represents the predictions of the MRSSM model, as we vary $\mu_{\rm eff}$ from 0.3 to 2.5 TeV and the (lighter) stop mass from $M_R/2$ to about 600 GeV. The rationale of choosing such ranges in explained below. However, for the moment, it is immediately clear that our predictions are nicely consistent with both the ATLAS and CMS observations, as the line passes clearly through both the allowed strips before entering the region ruled out by the dijet constraints. We may claim, therefore, to have a neat explanation of the observed diphoton excess (and the absence of other signals) in the MRSSM, without having had to extend the field content of the model specifically for this purpose.

It remains to be verified that the parameter choices in the MRSSM as made by us do not run into conflict with the major theoretical and experimental hurdles which affect other new physics models, especially the MSSM, which our model resembles in its basic aspects. We have already seen, in the Introduction, that constraints from direct searches for gluinos and squarks of the first two generations can be easily accommodated by making these sparticles heavy, without pushing up the stop mass(es) or inducing a large stop mixing. There is no restriction on making the sleptons in the model too heavy for observation at the LHC—at least with the existing data. We have mentioned that the model contains light stops and other scalars S and T, which, together, contribute enough to explain the Higgs mass constraint. It is also well known [42] that the mass of the stop can be as low as around 300 GeV if it is nearly degenerate with the neutralino $\tilde{\chi}_1^0$. We have, of course, demanded that the stop masses be above $M_R/2 \simeq 375$ GeV. The value of $\mu_{\rm eff} \lesssim \sqrt{4\pi}M_R \lesssim 2.6$ TeV is well within the perturbative limit. A more serious constraint on μ is known to arise *in the MSSM* from the vacuum stability issue, where the presence of terms like $-\mu Y_t \tilde{t}_L \tilde{t}_R^* H_d$ in the Lagrangian induces vacuum instability in the large-μ limit [9]. However, in the MRSSM, such terms are forbidden by R symmetry, and therefore the issue does not arise.[2] Moreover, as the mixing of the R_u^0 with the Higgs scalar is very small (see Sec. III), we do not expect constraints from the measured Higgs boson signal strengths, either.

In Fig. 3, we show the allowed region in the $\mu_{\rm eff}$-$m_{\tilde{t}_1}$ plane, the one-dimensional projection of which is the solid (red) line in Fig. 2. The region shaded black on the left corresponds to $m_{\tilde{t}_1} < M_R/2$ and is precluded by the requirement that the R_u^0 decays dominantly to photon and gluon pairs. The large triangular region, shaded gray, on the right side indicates too low a diphoton signal at the LHC (corresponding to the bottom section of Fig. 2), while the hatchings show the experimentally allowed regions using the same conventions as in Fig. 2. The small gray wedge on the top left of the panel indicates the dijet constraint, which is not very strong for the MRSSM solution. It may be seen from the figure that, even for reasonably low values of the stop mass and the $\mu_{\rm eff}$ parameter, we do get a viable solution for the diphoton excess.

We note, however, that this MRSSM solution leads to the prediction of a somewhat low width of 100 MeV or less for the R_u^0 resonance. This, while definitely larger than the Higgs boson width in the SM (4 MeV), is still small compared to the widths of the W and Z bosons. We can attribute the corresponding long life of the R_u^0 to the fact that it can only decay through one-loop diagrams. After all, it is an inert scalar. A small decay width is not a problem for the model at this stage of experimentation, since the kind of low statistics available at the moment leads to very poor estimations of the decay width. It is also important to note that larger widths of 200 MeV or more are incompatible with the nonobservation of a dijet excess—this is a generic feature of models having a scalar decaying *exclusively* to $\gamma\gamma$ and gg modes.

[2]Any small terms of this nature generated by the *breaking* of R symmetry will be suppressed by the small gravitino mass $m_{3/2}$ and would not have any significant phenomenological consequences.

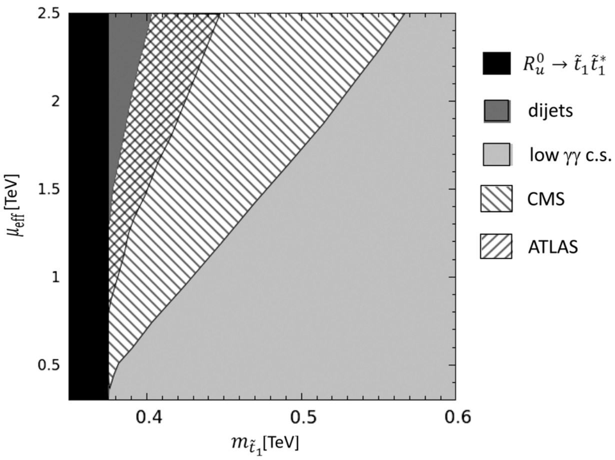

FIG. 3. Illustrating the allowed parameter values for the mass of the light stop $\tilde{t}_1$ (near degenerate with the $\tilde{t}_2$ in this model) and the effective coupling $\mu_{\rm eff}$. The solid shaded regions are ruled out by some constraint or the other, while the hatched regions are permitted by the experimental data.

V. CRITICAL SUMMARY

In this article, therefore, we have shown that, among the various possible explanations of the diphoton excess observed at the LHC, there exists the possibility of a SUSY solution which invokes an extra symmetry—the $U(1)_R$ symmetry—but does not require us to postulate new fields specifically to explain the effect. As for the model itself, apart from introducing a pair of new scalars and some superfields to convert the gauginos from Majorana to Dirac fermions, it retains the MSSM field content. We also obtain a good explanation of the failure of the LHC to discover SUSY signals to date. In addition, we require the R symmetry to be broken by scalar trilinear operators (and Majorana mass terms), for otherwise the inert scalars could not be produced at all in hadron-hadron collisions.

An obvious question to be asked before concluding this analysis is whether there are any confirmatory tests which could be used to verify if the ideas presented here are indeed correct. This can be answered quite easily in the affirmative. We argue as follows. The straight line shown in Fig. 2 enters the allowed region only if the (lighter) top squark has a mass in the range of a few hundred GeV, which would bring it very much within the kinematic range accessible for discovery at the LHC Run 2. This is also apparent from Fig. 3. Moreover, the neutral scalars R_u^0 and R_d^0 will be accompanied by their charged counterparts $R_u^\pm$ and $R_d^\pm$, and one could perhaps expect the mass ranges not to be very different. Charged scalars, of course, are easy to detect, and if they lie within the kinematic range of the LHC (as we have good reason to suspect), it cannot be long before they will be discovered. Thus, we have a couple of very clear ways in which the model in question can be falsified. The truth will only be known when more data are acquired and analyzed, but for the moment, we may rest satisfied that the MRSSM has enough pleasing features to be taken very seriously as an explanation of the recent LHC enigma.

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Note added.—After the original version of this article appeared on the arXiv, an appraisal of this work (among several others) appeared in Ref. [43]. Three issues were raised. The first issue, or set of issues, enumerated 1–5 in Ref. [43], are all related to the breaking of R symmetry and can be explained away as small effects due to the small gravitino mass. A second, more serious, criticism was that in the original version we had proposed $v_u = v_d$ to get a large stop mixing angle, but this leads to $\tan\beta = 1$, which is phenomenologically disfavored. In this version, however, there is no longer such a requirement, as Eq. (13) makes clear. It is, in fact, enough to tune the soft SUSY-breaking parameters for arbitrary values of $\tan\beta$. Finally, the issue of vacuum stability for large values of A_t was raised, but A_t is also suppressed in our model by the smallness of the gravitino mass. There is thus no problem in exploring the phenomenology of this model using the package SARAH mentioned in Ref. [43], and the caveat mentioned therein can be safely sidestepped.

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