

A closer look to the sgoldstino interpretation of the diphoton excess

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ABSTRACT: We revisit the sgoldstino interpretation of the diphoton excess in the context of gauge mediation. While the bound on the gluino mass might seem to make the sgoldstino contribution to the diphoton excess unobservable, we show that the interpretation is viable in a thin, near critical region of the parameter space. This regime gives rise to drastic departures from the standard gauge mediation picture. While the fermion messengers lie in the (10–100) TeV range, some scalar messengers are significantly lighter and are responsible for the sgoldstino production and decay. Their effective coupling to the sgoldstino is correspondingly enhanced, and a non-perturbative regime is triggered when light and heavy messenger masses differ by a factor $\gtrsim 4\pi$. We also comment on the possible role of an R -axion and on the possibility to decouple the sfermions in this context.

KEYWORDS: Supersymmetry Phenomenology

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1 Introduction

The excess in the diphoton channel recently reported by the ATLAS [1] and CMS [2] experiments at the Large Hadron Collider (LHC) at an invariant mass of approximately 750 GeV has prompted a variety of possible interpretations. The interpretation in terms of the production and decay of a sgoldstino [3–7] places the underlying new physics into the wider context of supersymmetric extensions of the Standard Model (SM), thus going beyond the mere parameterisation of the effect in terms of ad hoc dynamics. In fact, the very dynamics responsible for the generation of gluino and photino masses through the F -term of the goldstino superfield also provide, as a consequence of supersymmetry, a gluon and a photon decay width for the sgoldstinos.

The connection between gaugino masses and decay widths is most easily illustrated in terms of an effective description of the interaction between the goldstino superfield and the SM gauge superfield strengths W_a^α (the index a labels the different gauge fields),

$$\mathcal{L}_{\text{eff}} = \frac{c_a}{\Lambda} \int d^2\theta X W_a^\alpha W_\alpha^a, \tag{1.1}$$

where Λ represents the scale at which the effective operator is generated and the dimensionless coefficient c_a takes everything else into account. If X is the only superfield getting an

F -term, its fermion component is the goldstino (the Goldstone of supersymmetry breaking) and its scalar partner is the sgoldstino. In terms of the F -term vev F , the gaugino masses are given by $M_a = 2c_a F/\Lambda$, and

$$\mathcal{L}_{\text{eff}} = \frac{M_a}{2F} \int d^2\theta X W_a^\alpha W_\alpha^a = \frac{M_a}{2} \lambda_a \lambda_a + \frac{M_a}{2\sqrt{2}F} (s v_a^{\mu\nu} v_{\mu\nu}^a - a v_a^{\mu\nu} \tilde{v}_{\mu\nu}^a) + \dots, \quad (1.2)$$

where λ_a and $v_a^{\mu\nu}$ are the gauginos and gauge field strengths associated to W_a^α , and the scalar component of X has been decomposed in its real and imaginary parts.

As we show below, the effective description in eq. (1.2) can hardly account for the diphoton excess in the context of concrete UV completions, in particular if gaugino masses originate from gauge mediation. The problem is not apparent in the effective description, especially if the coefficient of the effective operator is expressed in terms of M_a/F , as in eq. (1.2). The way out we present in this paper requires supersymmetry to be badly broken in the dynamics underlying the effective interaction, $\sqrt{F} \sim \Lambda$, in which case the relevant effects are not captured by the effective description in eq. (1.2), valid for $\sqrt{F} \ll \Lambda$. The problem, and its solution, are discussed in section 3 (after a few preliminaries in section 2), in the simple case in which the operator in eq. (1.2) originates from a loop of chiral messengers, as in minimal gauge mediation. The interpretation of the excess relies on the onset of a non-perturbative regime, when approaching a critical point where one of the scalar messengers becomes light. In section 4 we speculate on the possible role of such a regime in simple models of supersymmetry breaking. In section 5 and 6 we comment on the possible role of an R -axion and of D -terms raising the sfermion masses well above the experimental limits. In section 7 we summarize and conclude.

2 The effective description

In this section we take eq. (1.2) at face value and show how the size of the diphoton excess translates into constraints on its parameters. A similar derivation was done in refs. [3–7].

Although, needless to say, the very existence of the anomaly is not yet established, we will assume that it corresponds to the production of a scalar resonance, identified with a scalar component of X , decaying into two photons. The interaction in eq. (1.2) provides the necessary ingredients for the production of the resonance through gluon fusion and its decay into photons. We consider a reference value of $\sigma_{\gamma\gamma} \equiv \sigma(pp \rightarrow s \rightarrow \gamma\gamma) = 6 \text{ fb}$ for the cross section at 13 TeV, see e.g. [8–13]. In the light of the presently uncertain experimental situation, we do not aim at accounting for a possibly large width of the resonance.

As we will see, obtaining a large enough partial width $\Gamma(s \rightarrow \gamma\gamma) \equiv \Gamma_{\gamma\gamma}$ is not at all trivial. Let us then conservatively consider the minimum value of $\Gamma_{\gamma\gamma}$ necessary to account for the anomaly. It is easy to see that such a minimum value is obtained when i) $\gamma\gamma$ and the partons pp involved in the production are the only decay channels, so that $\Gamma_{\text{tot}} = \Gamma_{\gamma\gamma} + \Gamma_{pp}$, ii) Γ_{pp} dominates the width, and iii) the resonance is produced through

gluon fusion ($pp = gg$). Which happen to be quite plausible conditions. Then, one gets¹

$$\Gamma(s \rightarrow \gamma\gamma) \approx 0.3 \text{ MeV} \left(\frac{\sigma_{\gamma\gamma}}{6 \text{ fb}} \right). \quad (2.1)$$

In terms of the effective interaction in eq. (1.2), the prediction for the photon partial width is [17]

$$\Gamma(s \rightarrow \gamma\gamma) = \frac{m_s^3 M_\gamma^2}{32\pi F^2}, \quad M_\gamma = c_W^2 M_1 + s_W^2 M_2, \quad (2.2)$$

where $m_s \approx 750 \text{ GeV}$ is the mass of the resonance and M_γ is expressed in terms of the bino and wino masses M_1 and M_2 , and the Weinberg angle θ_W .² Comparing eqs. (2.1) and (2.2) we obtain

$$\sqrt{F} \lesssim 5 \text{ TeV} \left(\frac{M_\gamma}{200 \text{ GeV}} \right)^{1/2} \left(\frac{6 \text{ fb}}{\sigma_{\gamma\gamma}} \right)^{1/4}. \quad (2.3)$$

In this effective approach, the size of the diphoton excess points to a very low scale of supersymmetry breaking. It is not easy to deal with such a low scale, as we expect gauge mediation to provide the main source of gaugino masses at this scale and gaugino masses to be loop suppressed, as we discuss in the next section.

3 A simple ultraviolet completion

Let us now discuss in greater detail the interpretation of the diphoton anomaly taking into account the origin of the effective interaction in eq. (1.2). We assume that gaugino masses are obtained at the one loop level through the exchange of messenger superfields, directly coupled to supersymmetry breaking, in vectorlike representations of the SM group G_{SM} , as in minimal gauge mediation. Note that on general grounds [4, 11, 12, 18–30], the interpretations of the diphoton anomaly also requires the existence of vectorlike fields, on top of the 750 GeV resonance, mediating its production and decay. The gauge mediation messengers play precisely that role, thus providing a wider context for the existence of the vectorlike fields as well.

To be specific, we add to the minimal supersymmetric SM (MSSM) field content a chiral superfield X , with non-vanishing scalar and F -term vevs

$$X = x + \sqrt{2}\psi\theta + f\theta^2, \quad \langle X \rangle = M + F\theta^2, \quad x = M + \frac{s + ia}{\sqrt{2}}. \quad (3.1)$$

The vev of X plays the role of the supersymmetry breaking spurion of minimal gauge mediation (M and F can be taken positive without loss of generality). On top of that, the dynamical degrees of freedom of X also play a role here. In particular, the 750 GeV resonance will be associated to the real scalar s . Both s and a are assumed for simplicity to be mass eigenstates. The origin of the supersymmetry breaking masses of s and a is a model-dependent issue, which we will not investigate in this section. As mentioned, if

¹We have used a K factor $K_{gg} \approx 2.8$, as in [14]; the Higgs cross-section $\sigma_{13\text{TeV}}(gg \rightarrow H(750 \text{ GeV})) \approx 736 \text{ fb}$ [15]; and $C_{gg} = 2137$ [10] for the gluon parton distribution function from NLO MSTW 2008 [16].

²Other open channels like the goldstino decay into a pair of goldstinos are negligible in the region of parameter space we are interested in, see discussion in section 5.

F were the only source of supersymmetry breaking, ψ would be the goldstino and x the sgoldstino. We also add messenger superfields $\Phi_i, \bar{\Phi}_i$ in irreducible, conjugated (possibly real) representations of the SM group. In order to generate one-loop masses for the three gauginos, the messengers should have non-trivial transformations under all the three SM gauge factors. They are coupled to supersymmetry breaking through X only,

$$\mathcal{L}_\Phi = \int d\theta^2 \lambda_i X \Phi_i \bar{\Phi}_i + \text{h.c.}, \tag{3.2}$$

where the coupling can be taken diagonal and positive without loss of generality.³ In the following, we will denote by λ_m the minimum value of the couplings λ_i (at the messenger scale).

Gaugino masses arise from eq. (3.2) through the standard gauge mediation mechanism. Sfermion masses also get a contribution from eq. (3.2), which however is not necessarily the only, nor the dominant, one (see e.g. section 6). We will therefore assume that the sfermions do not play a role in the diphoton anomaly.

Let us now discuss whether the diphoton excess can be accounted for in this setup. Before discussing the scenario we are interested in, in which $F \sim \lambda M^2$, we show that this is not possible in the $F \ll \lambda M^2$ limit.⁴

3.1 $F \ll \lambda M^2$

In the limit where $F \ll \lambda_m M^2$, supersymmetry breaking can be neglected when integrating out the messengers $\Phi_i + \bar{\Phi}_i$, whose fermion and scalar components all have masses close to $\lambda_i M$. The effective interaction in eq. (1.2) follows, with

$$\frac{c_a}{\Lambda} = \frac{\alpha_a}{8\pi M} N_a, \quad \text{giving} \quad M_a = \frac{\alpha_a}{4\pi} \frac{F}{M} N_a, \tag{3.3}$$

where a labels the factor of the SM group ($U(1)_Y, SU(2)_L, SU(3)_c$ respectively), and $N_a = \sum_i N_{a,i}$, and $N_{a,i}$ is the corresponding Dynkin index of $\Phi_i + \bar{\Phi}_i$. For instance, if the messengers form complete $SU(5)$ multiplets, $N_a = 1$ for a $\mathbf{5} + \bar{\mathbf{5}}$ and $N_a = 3$ for a $\mathbf{10} + \bar{\mathbf{10}}$.

It is now easy to show that the numerical results obtained in the previous section are not phenomenologically viable in this context. The explicit expression of the photino mass is

$$M_\gamma = \frac{\alpha}{4\pi} \frac{F}{M} N_\gamma, \quad N_\gamma = 2 \text{Tr}(Q^2), \tag{3.4}$$

where $\alpha = e^2/(4\pi)$ and the Dynkin N_γ is obtained tracing on the Φ_i superfields only. Plugging the above expression in eq. (2.2), the dependence on F drops out,

$$\Gamma(s \rightarrow \gamma\gamma) = \frac{m_s^3}{M^2} \frac{\alpha^2}{(8\pi)^3} N_\gamma^2. \tag{3.5}$$

³In eq. (3.2) we have omitted explicit mass terms for the messengers. In the presence of the latter, the supersymmetric mass of the messengers is not necessarily related to the sgoldstino coupling; except if for instance the messengers and X are charged under an extra $U(1)$ symmetry, a scenario discussed in section 6.

⁴After our paper was submitted to arXiv, and before it became available, ref. [31] appeared, which also obtained this result.

Eq. (2.1) then gives an upper limit on the messenger scale

$$M \lesssim 70 \text{ GeV } N_\gamma \left(\frac{6 \text{ fb}}{\sigma_{\gamma\gamma}} \right)^{1/2} \Rightarrow \lambda_m N_\gamma \gtrsim 14 \frac{M_m}{\text{TeV}} \left(\frac{\sigma_{\gamma\gamma}}{6 \text{ fb}} \right)^{1/2}, \quad (3.6)$$

where $M_m = \lambda_m M$ is the mass of the lightest messengers. The experimental bounds on the latter⁵ require relatively large values of the Dynkin index. Such values can be achieved in the presence of a rich enough set of messengers. For example, a full family of messengers, filling a $\bar{\mathbf{5}} + \mathbf{10}$ representation of SU(5), would give $N_\gamma = 32/3$, and $\lambda_m > 1.3$ would suffice to allow TeV scale messengers.

The problem arises from the gaugino masses, the gluino mass M_3 in particular. Being loop suppressed with respect to the messenger masses M_m , the experimental bound on M_3 forces M_m to be in the $\mathcal{O}(100 \text{ TeV})$ region, barring unrealistic values of N_3 . We have in fact

$$M_3 = \frac{\alpha_3}{4\pi} \frac{F}{M} N_3 \ll \frac{\alpha_3}{4\pi} M_m N_3 \Rightarrow M_m \gg \frac{130 M_3}{N_3}. \quad (3.7)$$

When plugged in eq. (3.6), such large messenger masses require unrealistic values of N_γ .

In the expressions above, we have assumed that only s contributes to the diphoton anomaly. The possibility that both s and a contribute is often considered, also in connection to the possibility of explaining a possibly sizeable width of the 750 GeV resonance [3, 6, 7]. The presence of both contributions would enhance the photon width by a factor of two, but would not change our conclusions.

3.2 $F \sim \lambda M^2$

Drastic departures from the grim predictions of the previous subsection arise in the regime in which supersymmetry breaking is sizeable, and the effective description in eq. (1.2) does not apply. In order to obtain the expressions for the partial widths $\Gamma_{\gamma\gamma}$, Γ_{gg} in this regime, we first write the relevant interactions. Omitting for simplicity the messenger index i , the mass terms for the fermion ($\psi, \bar{\psi}$) and scalar ($\phi, \bar{\phi}$) components of the messengers $\Phi, \bar{\Phi}$ are

$$-\mathcal{L}_{\text{mess}}^{(2)} = (\lambda M \psi \bar{\psi} + \text{h.c.}) + \lambda (\phi^\dagger, \bar{\phi}) \begin{pmatrix} \lambda M^2 & F \\ F & \lambda M^2 \end{pmatrix} \begin{pmatrix} \phi \\ \bar{\phi}^\dagger \end{pmatrix}. \quad (3.8)$$

The fermion messengers ψ and $\bar{\psi}$ form Dirac spinors with mass $M_m = \lambda M$. The scalar mass eigenstates are

$$\phi_{h,l} = \frac{\phi \pm \bar{\phi}^\dagger}{\sqrt{2}}, \quad \text{with masses } m_{h,l}^2 = M_m^2 \pm \lambda F. \quad (3.9)$$

We assume to be in the regime $F \leq \lambda_m M^2$, so that no messenger develops a vev.

⁵The experimental bounds on the mass of the lightest messengers depend on their decay mode. One possibility is that they decay through a small coupling to the MSSM fields. For example, this is the case if the presence of a superpotential interaction $W_{\text{LQ}} = \bar{\Phi} Q L$. Then, at TeV energies these interactions lead to leptoquark type couplings involving the lightest fields ϕ_l . The lower bounds on their mass are around 650 GeV [32, 33]. In various instances below we consider higher multiplicities in the number of messengers, therefore leading to higher multiplicities of the leptoquark couplings. We will therefore use 1 TeV as a reference lower bound.

The partial widths $\Gamma_{\gamma\gamma}$, Γ_{gg} and the gaugino masses M_a are generated by the gauge interactions and the trilinear messenger interactions

$$-\mathcal{L}_{\text{mess}}^{(3)} = \left(\lambda \frac{s + ia}{\sqrt{2}} \psi \bar{\psi} + \text{h.c.} \right) + \sqrt{2} \lambda^2 M s (|\phi_l|^2 + |\phi_h|^2). \quad (3.10)$$

In particular, the decay widths of s (but not of a) get a contribution from trilinear interactions with the scalar messengers with masses $m_{h,l}^2 = \lambda^2 M^2 \pm \lambda F$. Parametrically, the strength of the trilinear coupling of the lightest messenger is measured by the effective coupling

$$g_{\text{eff}} = \lambda \frac{M_m}{m_l} = \frac{\lambda^2 M}{(\lambda^2 M^2 - \lambda F)^{1/2}}. \quad (3.11)$$

The crucial observation is that there exists a small region of the parameter space, near the critical point $M = \sqrt{F/\lambda}$, where the lighter scalar messenger is significantly lighter than its natural scale M_m , and its effective coupling to the scalar resonance s is correspondingly enhanced. This is the regime in which the interpretation of the diphoton excess has a chance to be phenomenologically viable, and that we will study in detail. We will call it the “near-critical” regime. When the enhancement becomes very large, the system enters a strongly interacting regime.

A few comments are in order.

- In the near-critical regime, supersymmetry is maximally broken, $F \approx \lambda M$, and the dynamics is far from being described, even qualitatively, by the effective approach in eq. (1.2).
- The near-critical region is fine-tuned, as F and λM need to be close, with the fine-tuning parameter given by $\Delta = (M_m/m_l)^2 = (g_{\text{eff}}/\lambda)^2$. In section 4 we will speculate on a possible dynamical origin of such a degeneracy.
- A lower bound to the potential is guaranteed by supersymmetry, independently of the size of the trilinear coupling. Assuming the potential is stabilized in a nearly critical point, we expect the minimum to be meta-stable, with a model dependent lifetime.
- In the near-critical regime, the gaugino masses are not drastically enhanced, for given M (a moderate enhancement comes from the $F \sim \lambda M^2$ corrections to the standard $F \ll \lambda M^2$ expressions). On the other hand, the diphoton anomaly is controlled by the lightest scalar messenger, and it is now possible to keep its mass light (to get a sizeable diphoton signal) while allowing M to be much larger (to get a gluino mass above experimental bound).
- In the presence of multiple messengers, if none of the couplings λ_i are (approximately) degenerate, only one messenger, the one with $\lambda_i = \lambda_m$, can benefit from an enhanced coupling. On the other hand, in the presence of an (approximate) degeneracy of different λ_i , e.g. consequence of a symmetry, more scalar messengers can be light at the same time. Such a degeneracy should involve messenger with same quantum

numbers under the SM gauge group, as gauge radiative corrections could otherwise spoil the degeneracy. For example, in the context of unified theories, the different SM components of a unified multiplet would have equal couplings λ_i at the grand unification (GUT) scale, but the different RGE running would lift the degeneracy at low energy.

The broad picture that emerges has therefore: a resonance at 750 GeV associated to the sgoldstino s ; a messenger scale M_m of a few tens of TeV, so that the loop suppressed gluino mass can be above the experimental bound; a number of messengers (with same SM quantum numbers) with a near critical coupling $\lambda \approx F/M^2$ with an anomalously light scalar component (at the 1–2 TeV scale), responsible of the production and decay of the ~ 750 GeV resonance, and with all the other components at the scale M_m ; TeV scale gaugino masses generated by both light and heavy messengers. With this broad picture in mind, let us now proceed to a more detailed discussion.

The general one-loop expressions for the partial decay widths into gluons and photons of s and a , due to the loop of scalar and fermion messengers are given in the appendix. There, we also provide the expressions for the decays into ZZ , $Z\gamma$, WW . The decay into two Higgses is absent at the one-loop level. In the limit in which the sgoldstino is lighter than the messengers, $m_s^2 \ll 4M_m^2, 4m_l^2$, those expressions become

$$\Gamma(s \rightarrow gg) = \frac{m_s^3}{M^2} \frac{4}{9} \frac{8\alpha_3^2}{(8\pi)^3} \left| \sum_i N_{3,i} \left[1 + \frac{1}{4} \left(\frac{\lambda_i^2 M^2}{m_l^2} + \frac{\lambda_i^2 M^2}{m_h^2} \right) \right] \right|^2, \quad (3.12)$$

$$\Gamma(a \rightarrow gg) = \frac{m_a^3}{M^2} \frac{8\alpha_3^2}{(8\pi)^3} N_3^2, \quad (3.13)$$

$$\Gamma(s \rightarrow \gamma\gamma) = \frac{m_s^3}{M^2} \frac{4}{9} \frac{\alpha^2}{(8\pi)^3} \left| \sum_i N_{\gamma,i} \left[1 + \frac{1}{4} \left(\frac{\lambda_i^2 M^2}{m_l^2} + \frac{\lambda_i^2 M^2}{m_h^2} \right) \right] \right|^2, \quad (3.14)$$

$$\Gamma(a \rightarrow \gamma\gamma) = \frac{m_a^3}{M^2} \frac{\alpha^2}{(8\pi)^3} N_\gamma^2. \quad (3.15)$$

As mentioned, we expect M to be of order of a few tens of TeV or more, in order for the loop suppressed gluino mass to be above the experimental bounds. Eqs. (3.12)–(3.15) then show that the field a cannot play a role in the diphoton anomaly, as the corresponding widths are suppressed by M^2 . Hence in this context a possibly sizeable width of the diphoton resonance cannot be explained in terms of the production of two resonances close in mass associated to the fields s and a [3, 6, 7]. On the other hand, when m_l is around the TeV scale, the corresponding terms in eqs. (3.12) and (3.14) dominate and we have

$$\Gamma(s \rightarrow gg) \approx \frac{m_s^3}{m_l^2} \frac{4}{9} \frac{8\alpha_3^2}{(8\pi)^3} \left(\bar{N}_3 \frac{\lambda_m^2 M}{4m_l} \right)^2, \quad (3.16)$$

$$\Gamma(s \rightarrow \gamma\gamma) \approx \frac{m_s^3}{m_l^2} \frac{4}{9} \frac{\alpha^2}{(8\pi)^3} \left(\bar{N}_\gamma \frac{\lambda_m^2 M}{4m_l} \right)^2, \quad (3.17)$$

where $\bar{N}_3, \bar{N}_\gamma$ are the Dynkin indices summed only over the messengers in the near-critical regime, i.e. with a light scalar degree of freedom (assumed for simplicity to have all the same mass m_l), which all necessarily have $\lambda_i \approx \lambda_m$.

In the expressions for the partial widths in eqs. (3.16) and (3.17), the (multi-TeV)² suppression M^2 has been replaced by m_l^2 , which is allowed to be close to its $\mathcal{O}(\text{TeV}^2)$ experimental bounds. On top of that, the light scalar messengers further enhance the signal through the additional factor $(\lambda_m^2 M/m_l)^2$, which corresponds to an enhanced effective coupling g_{eff} , see eq. (3.11). To get a feeling of the size of the effect, we observe that the same enhancement could be obtained in the case of a standard fermion loop (with same SM quantum numbers, mass, and in the same limit used in eqs. (3.16) and (3.17)) using a Yukawa coupling $\lambda_f = \lambda_m^2 M/(4m_l)$ (with the Yukawa normalised as in eq. (3.10)). Large trilinears have also been considered in refs. [34–36].

Note that $\lambda_m^2 M$ cannot be taken arbitrarily larger than m_l , as the light messenger interactions become non-perturbative for large g_{eff} . The qualitative naive dimensional analysis (NDA) estimate of the onset of the strongly interacting regime is

$$g_{\text{eff}} \sim g_{\text{eff}}^* \equiv 4\pi. \tag{3.18}$$

As a consequence, we are only allowed to consider values of the heavy messenger scale bound by $M_m \lesssim (g_{\text{eff}}^*/\lambda_m)m_l$. Notice in particular that lowering λ_m allows to rise the upper bound on M_m .⁶

For completeness, let us exemplify the discussion above in a simple setup. We consider a set of n messenger pairs $V + \bar{V}$ with SM quantum numbers $(\mathbf{3}, \mathbf{2})_{-5/6} + (\bar{\mathbf{3}}, \mathbf{2})_{5/6}$. In the left hand side of figure 1 we show the values of the gluino mass (green) and of the mass of the lightest set of scalar messengers ϕ_l (dashed blue) fitting the signal strength as a function of $F/(\lambda M^2)$. Each contour line corresponds to a different number of messengers, as indicated on top of each curve. The plot is done using the full one-loop decay widths given in the appendix.

The lower bound on the gluino mass ($M_3 \gtrsim 1.7 \text{ TeV}$) requires increasing F/M , while M needs to be increased accordingly in order to keep a constant signal strength, see eqs. (3.16) and (3.17). In fact, the plot shows that the parameters of the theory are pushed into the near-critical regime $F/(\lambda M^2) \approx 1$ for realistic gluino and messenger masses. In the right hand side of figure 1 we zoom into the critical region. We find that, as long as we accept a significant tuning, this toy example is able to both fit the signal and exceed the present lower bounds on scalar messengers and gluinos.

On the other hand, the plot shows that in the effective regime of eq. (1.2), i.e. for $F \ll \lambda M^2$, the gluino mass cannot be accounted for. For high enough n , the plot in figure 1 is only indicative, as the theory presents close Landau poles in both the running of λ and the gauge couplings, which is discussed in detail in section 3.4. Also notice that for $F/(\lambda M^2)$ close enough to 1 the trilinear coupling becomes non-perturbative $g_{\text{eff}} \sim g_{\text{eff}}^*$.

⁶We can refine the perturbativity bound of the EFT containing the sgoldstino and the ϕ_l field. With the normalization $\delta\mathcal{L} = \sqrt{2}\lambda^2 M s |\phi_l|^2$ one gets $g_{\text{eff}}^* \sim 4\pi/(4dn)^{1/4}$, where n is the number of messenger fields and d the length of the representation. Although the bound is obtained under the assumption $m_l \gg m_s$, a similar bound can be obtained for $m_l \ll m_s$. See also ref. [37].

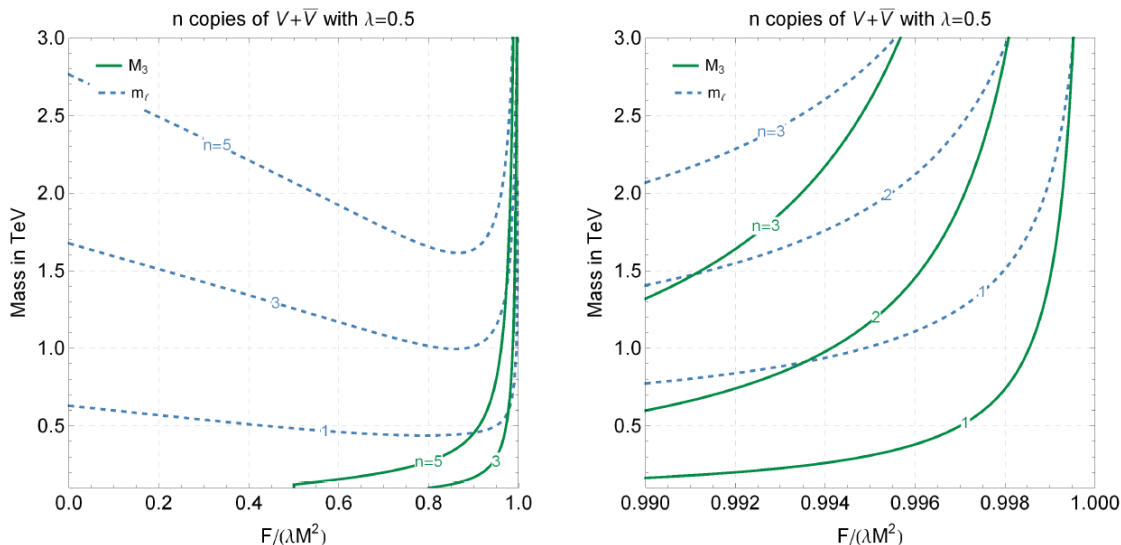


Figure 1. *Left:* values of the lightest messenger mass m_l (dashed blue) and gluino mass M_3 (green) fitting the signal strength (with $\sigma_{\gamma\gamma} = 6$ fb), as a function of the parameter $F/(\lambda M^2)$. The n messengers are assumed to be in the $(\mathbf{3}, \mathbf{2})_{\pm 5/6}$ SM irreps. *Right:* left plot zoomed in.

3.3 Quantitative analysis

Let us now show quantitatively that the diphoton excess can indeed be reproduced while keeping m_l and M_3 above the experimental limits. As mentioned, the excess is controlled by m_l , while gaugino masses depend on M_m . We want to keep m_l around the TeV scale, above its experimental bound, which helps fitting the excess. Unlike in the case of section 3.1, this does not make the gluino mass unacceptably small, as we can now take M_m significantly heavier. Using eq. (3.17) and (2.1), we find that fitting the diphoton excess requires

$$\frac{g_{\text{eff}}}{g_{\text{eff}}^*} \approx \frac{6.9}{\bar{N}_\gamma} \left(\frac{m_l}{\text{TeV}} \right) \left(\frac{\sigma_{\gamma\gamma}}{6 \text{ fb}} \right)^{1/2}. \quad (3.19)$$

In order to keep g_{eff} below g_{eff}^* , while keeping $m_l \gtrsim 1$ TeV, the messengers with light scalar components should have $\bar{N}_\gamma \gtrsim 7$. We will discuss examples in section 3.4. Making the effective coupling g_{eff} large only gives a moderate gain with respect to the case in section 3.1. In fact, the results for the diphoton cross section are the same (for equal mass and SM quantum numbers of the relevant degrees of freedom) when $g_{\text{eff}} = 6 \lambda_m$, and the factor 6 limits the gain.

On the other hand, the previously hopeless situation with the gluino mass is now completely different. To start with, the standard expression in eq. (3.3) for the gluino mass gets a $\mathcal{O}(1)$ enhancement in the $F \sim \lambda M^2$ regime. In the near-critical region, $F \approx \lambda M^2$, the enhancement is given by a factor $\log 4 \approx 1.4$. Also, extra messengers not in the near-critical regime can contribute to the gluino mass, while being negligible in the diphoton signal. We therefore have

$$M_3 = \frac{\alpha_3}{4\pi} M_m \bar{N}_3 \log 4 + \Delta M_3, \quad (3.20)$$

where ΔM_3 is the contribution of the non-critical messengers, giving

$$M_m \approx \frac{100 \text{ TeV}}{\bar{N}_3} \left(\frac{M_3 - \Delta M_3}{\text{TeV}} \right). \quad (3.21)$$

Most important, such a large value of M_m is now allowed, as long as g_{eff} it is not too large. The value of g_{eff} required by eq. (3.21) is

$$\frac{g_{\text{eff}}}{g_{\text{eff}}^*} \approx \frac{8\lambda_m}{\bar{N}_3} \left(\frac{M_3 - \Delta M_3}{m_l} \right). \quad (3.22)$$

We therefore conclude that we can make M_3 large enough, while not exceeding the NDA bound on g_{eff} , if $\bar{N}_3 \gtrsim 8\lambda_m$. We will discuss examples in the next subsection. Note that smaller values of λ_m help with perturbativity.

3.4 Examples

Let us discuss a few examples of viable messenger field content, leading to $g_{\text{eff}} \lesssim g_{\text{eff}}^*$ in eqs. (3.19), (3.22). As mentioned, we prefer the near-critical messengers Φ_i to be given by n copies of the same SM irreducible representation, to guarantee that the near-equality of their coupling to X , possibly consequence of a symmetry, is not spoiled by gauge radiative corrections. In order to induce the diphoton signal, we need them to be colored and electrically charged. Other, non-critical, messengers can also be around. Their contribution to the diphoton signal will be negligible, but they can play a role in gauge coupling unification. On the other hand, if a Landau pole for the gauge coupling is to be avoided below the GUT scale, the total number of messengers cannot be too large.

Different model building avenues are available, depending on whether or not one aims at the perturbativity of gauge couplings up to the GUT scale and at gauge coupling unification. If the perturbativity of gauge couplings is not an issue, the bounds in eqs. (3.19), (3.22) can be easily satisfied while maintaining the light scalar interactions semi-perturbative. Let us then aim at models with perturbative gauge couplings up to the GUT scale (see [38, 39] for a related discussion). This requires $N_1, N_2, N_3 \lesssim 5$, where N_a are the Dynkin indices of all messengers (near-critical and not).

If gauge coupling unification is not an issue, extra, non near-critical superfields are not required, and we can assume $\bar{N}_i = N_i$. Addressing the gluino mass constraints while keeping g_{eff} under control is not an issue: the UV perturbativity condition $N_3 \lesssim 5$ is always compatible with $g_{\text{eff}} < g_{\text{eff}}^*$, in eq. (3.22), for an appropriate value of λ_m . The only drawback of a smaller λ_m is the higher fine-tuning necessary for near-criticality. As for fitting the signal strength, the relation $\bar{N}_\gamma = \bar{N}_2 + (5/3)\bar{N}_1$ shows that one can obtain relatively large values of N_γ in eq. (3.19) while keeping $N_1, N_2 < 5$.

To stick to known SM representations, let us consider as a first example n copies of $U^c + \bar{U}^c$, where U^c has the same quantum numbers as the SM up quark singlets. The Dynkin indices are $(N_1, N_2, N_3) = n(8/5, 0, 1)$. As the perturbativity of g_1 forces $n \leq 3$, we can take $n = 3$. We then have $N_3 = 3$ in eq. (3.22), which, for $M_3 - \Delta M_3 = m_l$, is compatible with $g_{\text{eff}} < g_{\text{eff}}^*$ when $\lambda_m \lesssim 0.4$. As for eq. (3.19), the three copies of $U^c + \bar{U}^c$ give $N_\gamma = 8$, which is compatible with $g_{\text{eff}} \lesssim g_{\text{eff}}^*$.

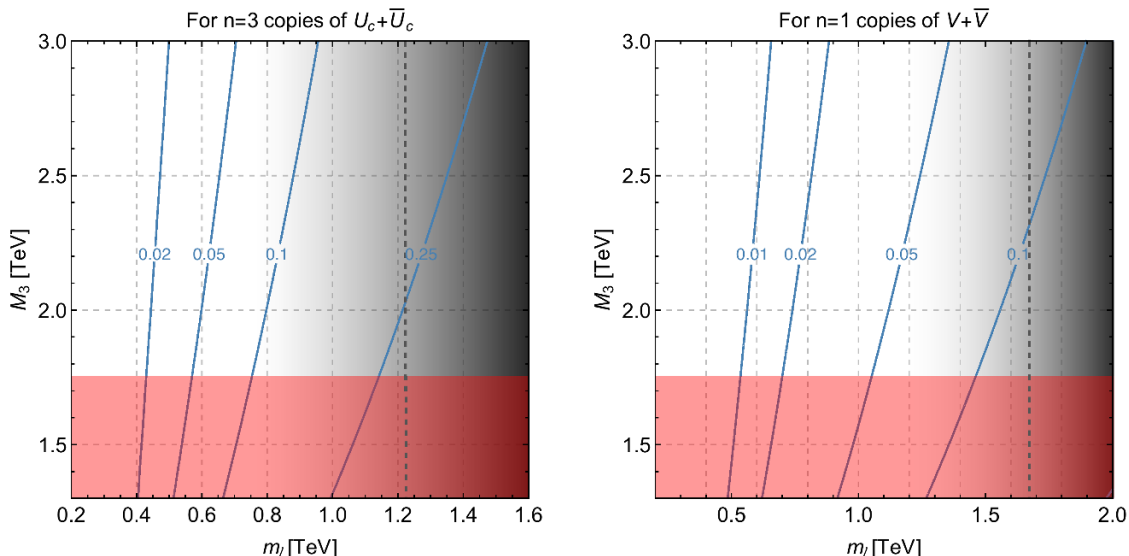


Figure 2. Contour lines of constant λ_m (blue) fitting the signal, for $\sigma_{\gamma\gamma} = 6$ fb, in the plane of gluino and light messenger masses. We have shaded in red the gluino mass exclusion and in a gray gradient the non-perturbativity region $g_{\text{eff}} \gtrsim g_{\text{eff}}^*$. The dashed curve corresponds to the isocurve $g_{\text{eff}} = g_{\text{eff}}^*$.

A second example, involving a SM representation included in the adjoint of SU(5), is the case of n copies of $V + \bar{V}$, where V is a fundamental of SU(3)_c and SU(2)_L with hypercharge $Y = -5/6$. The Dynkin indices are $(N_1, N_2, N_3) = n(5, 3, 2)$. As the perturbativity of g_1 forces $n \lesssim 1$, we take $n = 1$. We then have $N_3 = 2$ in eq. (3.22), which, for $M_3 - \Delta M_3 = m_l$, is compatible with $g_{\text{eff}} < g_{\text{eff}}^*$ when $\lambda_m \lesssim 0.25$. And in eq. (3.19) we have $N_\gamma = 34/3 \approx 11$, which is compatible with $g_{\text{eff}} \lesssim g_{\text{eff}}^*$.

In figure 2 we show two plots with contour lines in the (M_3, m_l) -plane corresponding to different values of λ_m fitting the signal. The plots are done with the one-loop formulas given in the appendix. On the left hand side for $n = 3$ copies of $U^c + \bar{U}^c$ fields, and as discussed above for this choice of parameters the model is on the edge of non-perturbativity, shown with a dashed curve. On the right hand side the plot is done for one set of $V + \bar{V}$ messengers in the adjoint of SU(5).

The two above examples have different predictions for the $pp \rightarrow \gamma Z, ZZ, WW$ rates. The first example ($U^c + \bar{U}^c$) predicts them to be well below the present limits: $\Gamma_{ZZ}/\Gamma_{\gamma\gamma} \approx 0.08$, $\Gamma_{Z\gamma}/\Gamma_{\gamma\gamma} \approx 0.6$. The second example ($V + \bar{V}$) predicts higher rates, but also below the present experimental limits, with an accidental suppression of the $Z\gamma$ rate: $\Gamma_{ZZ}/\Gamma_{\gamma\gamma} \approx 1.3$, $\Gamma_{Z\gamma}/\Gamma_{\gamma\gamma} \approx 0.02$, $\Gamma_{WW}/\Gamma_{\gamma\gamma} \approx 2.8$. The previous ratios have been obtained using the formulas in the appendix.

Neither of the previous examples preserves the successful gauge coupling unification achieved in the MSSM. The simplest way to preserve unification is to add extra fields (non near-critical) so that the messengers form complete SU(5) multiplets. This is possible only for the second example. The first example would require completing the $3 \times (U^c + \bar{U}^c)$ into $3 \times (\mathbf{10} + \bar{\mathbf{10}})$ of SU(5). In this case, however, the total Dynkin would be $N_1 = N_2 = N_3 = 9$, well above what perturbative gauge unification requires.

The second example, on the other hand, requires completing $V + \bar{V}$ to a full $SU(5)$ adjoint, with $N_1 = N_2 = N_3 = 5$, at the boundary of gauge coupling unification. In this case, the fields $V + \bar{V}$ are accompanied by an adjoint of $SU(3)_c$, Σ , an adjoint of $SU(2)_L$, W , and a singlet. If X is a singlet of $SU(5)$, the λ couplings of those fields are the same at the GUT scale, $\lambda_V = \lambda_\Sigma = \lambda_W$. The RGE running to low energies makes the triplet coupling λ_W lower than λ_V , which prevents $V + \bar{V}$ from being near-critical (only the triplet can be). On the other hand, if X is the singlet component of a $SU(5)$ adjoint, then $\lambda_V = \lambda_W/6 = \lambda_\Sigma/4$ at the GUT scale. For perturbative values of the couplings at the GUT scale, λ_V remains the smallest coupling during the whole running, and $V + \bar{V}$ can play the role of the near-critical fields.⁷

Other solutions can be obtained in the context of flipped $SU(5)$ [40, 41] or using messenger spectra that preserve gauge coupling unification but are not in full $SU(5)$ multiplets [42].

In summary, it is possible to choose a messenger spectrum such that the bounds in eqs. (3.19), (3.22) are satisfied with $g_{\text{eff}} < g_{\text{eff}}^*$ and that gauge couplings unify at the GUT scale, where they are at the border of perturbativity. The simplest example we found is the case of messengers forming a whole adjoint representation of $SU(5)$, with the $Y \neq 0$ components near-critical and the $Y = 0$ components off near-criticality.

3.5 Strong coupling

From the discussion of possible models in the previous section, the following dichotomy emerges between UV and IR non-perturbativity. On one side, one can choose to have large representations or a large number of messenger fields. This allows to interpret the diphoton excess at the expense of Landau poles at some tens or hundreds of TeVs, thus having strong dynamics in the UV, but a weakly coupled EFT at TeV energies.

On the other side, one can avoid Landau poles by incorporating a smaller number of messenger fields at the expense of tuning the trilinear λ_m close to criticality, and hence approaching strong dynamics in the IR. As an infrared effect, it does not give rise to Landau poles and it does not spoil the nice UV properties of supersymmetric theories.

In the IR non-perturbative regime, the numerical results showed above may receive large corrections. Moreover, the trilinear interaction leads to an attractive force between the light scalar messengers and, in the IR strong coupling regime $g_{\text{eff}} \gtrsim g_{\text{eff}}^*$, one expects a tower of bound states. In fact, a similar phenomenon is argued to happen in the MSSM if the trilinear interaction $A_t H \tilde{q}_L \tilde{u}_R$ becomes strong [43]. Thus we expect that the light scalar messengers form an S-wave color-singlet bound state S_b . The resonance S_b would be a tightly bound state, as the binding energy is controlled by m_t . And it would be a true bound state, as the formation time is controlled by the inverse of the binding energy, and the decay by perturbative QCD interactions. Then, the bound state S_b would mix with the sgoldstino and, since the constituents of S_b are colored, this would give rise to a direct coupling between gluons and the physical state. To our knowledge, the details of

⁷The adjoint containing $V + \bar{V}$ also contains a gauge singlet. Its Yukawa coupling is likely to be the smallest one at low energy. Therefore, when $V + \bar{V}$ is near critical, the singlet will be in the broken phase and develop a vev. This is not a problem, as the vev would not break the SM gauge symmetry.

the phenomenology and the interplay of the sgoldstino and the possible singlet resonance is far from settled. It would be interesting to further explore such phenomenology, perhaps through lattice techniques, especially if in the future the sgoldstino scenario near criticality will gain further support from the experiments.

4 Speculations on the origin of the near-critical regime

The near-critical regime requires an apparent fine-tuning making $m_l \ll \lambda M$ and enhancing g_{eff} in eq. (3.11). In this section, we speculate on a possible connection between a dynamical origin of near-criticality and the strong regime it induces.

Let us consider the simplest possible completion of the basic model building block considered in the previous section, eq. (3.2). The gauge group is G_{SM} . Besides the MSSM ones, the chiral superfields are X and a vectorlike set of messengers $\Phi + \bar{\Phi}$. For simplicity, we take Φ in a single irreducible SM representation. The superpotential is

$$W = W_{\text{MSSM}} + \lambda X \Phi \bar{\Phi} - F X, \tag{4.1}$$

where $\lambda > 0$, $F > 0$ with no loss of generality. While the above system is simple and well known, we are not aware of a thorough discussion of its behaviour near the critical point.

Let us recall the main features of the above system. We expand X as in eq. (3.1). Then the system has a supersymmetric minimum for $\langle x \rangle = 0$, $|\langle \Phi \rangle| = |\langle \bar{\Phi} \rangle| = \sqrt{F/\lambda}$. The system has two phases, controlled by the scalar vev of x , $\langle x \rangle = M$, which can also be taken real and non-negative. When $M > \sqrt{F/\lambda}$, the critical point, the messengers have no vev, the gauge symmetry is unbroken, and supersymmetry is broken by the F -term vev of X , $\langle f \rangle = F$. In such an unbroken phase, the tree level potential is flat, $V = F^2$. On the other hand, when $M < \sqrt{F/\lambda}$, both messengers get a vev of size $|\langle \Phi \rangle| = |\langle \bar{\Phi} \rangle| = \sqrt{F/\lambda - M^2}$ and the potential, as a function of M , is

$$V = F^2 - (F - \lambda M^2)^2, \tag{4.2}$$

which forces $M = 0$, where the potential has its absolute minimum, the gauge symmetry is broken, and supersymmetry is unbroken.

On the left of the critical point, $M < \sqrt{F/\lambda}$, the tree level potential provides a sufficiently accurate description (except perhaps very near the critical point, where it is nearly flat). On the right side, on the other hand, the flat direction for $M > \sqrt{F/\lambda}$ is lifted by the one-loop correction to the potential, which acquires a positive slope and let M slide towards the critical point, see figure 3.

In a region around the critical point, though, the system hides a non-perturbative regime, triggered by the growth of the coupling in eq. (3.11). Interestingly, this happens only in a tiny region, characterised by

$$|\lambda M^2 - F| \lesssim \lambda^2 \frac{F}{(4\pi)^2}. \tag{4.3}$$

We can then speculate on the possibility that the effective potential generated by strong interactions has a negative slope in some point of the near-critical region. If that were

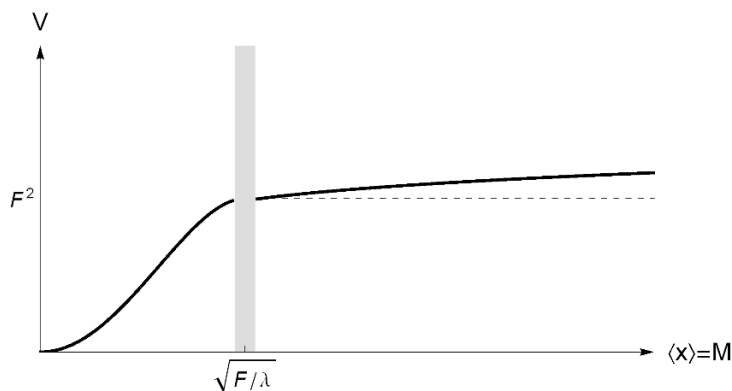


Figure 3. Schematic representation of the one-loop corrected potential in section 4 as a function of $\langle x \rangle$. The tree-level potential in the pseudo-flat direction is plotted with a dashed line. The shaded region corresponds to the region of near-criticality (the size is largely exaggerated).

the case, a metastable minimum would be generated for the potential in the near-critical region, thus providing a dynamical origin for the apparently fine-tuned closeness of λM^2 and F (and for the origin of supersymmetry breaking). But this is of course just wishful thinking. Still, we consider worth pointing out the existence of an ephemeral, but potentially interesting, non-perturbative regime in a simple and well known model, which might deserve further investigation.

If it stabilised in the near-critical region, the system we considered has an obvious, spontaneously broken R -symmetry. The phase of x would play the role of the R -axion and would be massless, up to supergravity R -breaking corrections. In the next section, we will make a few considerations on the possible implications of such a light R -axion, in a more general context.

5 The R -axion

The discussion in section 3 assumed that the only relevant interactions of the diphoton resonance are the ones with the messengers in eq. (3.2). On the other hand, the hidden sector dynamics can in principle give rise to alternative decay channels that could compete with the photon and gluon ones, thus affecting our phenomenological analysis. However, it is not unreasonable to assume that most of the hidden degrees of freedom are irrelevant because characterised by large, $\mathcal{O}(\sqrt{F})$ masses.

Even if that is the case, there are two possible decay channels that can not be ignored. First, the SUSY breaking mass of the sgoldstino $m_s^2/F^2|X|^4|_D$ leads to a decay of the sgoldstino into two goldstinos. Such a decay is negligible with respect to $\Gamma(s \rightarrow \gamma\gamma)$, see eq. (3.17), in the regions of parameter space that we consider. Then, a wide class of supersymmetry breaking models predicts the existence of a light degree of freedom, the R -axion, which could be relevant. In this section, we shortly outline the possible role of a light R -axion in the diphoton phenomenology. The R -symmetry plays a central role in most supersymmetry breaking models. If present, its spontaneous breaking is welcome to

allow for Majorana gaugino masses. The R -axion is the Goldstone associated to such a spontaneous breaking. As such, it is massless, up to the small explicit breaking of the R -symmetry provided by supergravity [44]. This breaking gives it a mass that, in our parameter space, is $\mathcal{O}(100 \text{ MeV})$ [45], with a non-negligible dependence on the hidden sector dynamics. If X is the only source of spontaneous R -symmetry breaking, as in the toy model in the previous section, the R -axion is the phase a of its scalar component, $x = (M+s/\sqrt{2})e^{ia/(\sqrt{2}M)}$ (where with an abuse of notations we are using the same notations for the fields s, a that we used for their linearised versions in eq. (3.1)). In general, the R -axion will have at least a component in a , if X has a non-vanishing R -charge. For simplicity, and to maximise the role of the R -axion in the diphoton phenomenology, we will assume that the R -axion coincides with a .

The radial component s , i.e. the $\sim 750 \text{ GeV}$ resonance, can then decay into two R -axions, with a partial width

$$\Gamma(s \rightarrow aa) = \frac{1}{64\pi} \frac{m_s^3}{M^2}. \tag{5.1}$$

Depending on the subsequent fate of the R -axion and on the relative size of the above width and the $gg, \gamma\gamma$ ones, the above channel, if present, can affect the discussion in section 3.

The mass of the R -axion is in the ballpark of the pion mass. If $m_a > 2m_\pi$, it will dominantly decay into two pions.⁸ If $m_a < 2m_\pi$, the relevant channels are R -axion decays into two photons, electrons or muons. The lepton decays are proportional to $m_a(m_f/M)^2$, see ref. [45], and can be suppressed with respect to the decay into photons for the messenger mass scales that we consider. Then, the decay into two photons is in principle relevant, as the values of the R -axion mass just happens to be in the window in which the two photons are collimated enough to be seen at the LHC as a single photon [9, 14, 46–56]. As a consequence, the decay $s \rightarrow aa$ could in principle also account for the diphoton signal, through the subsequent decay of the R -axions into two collimated photons. Unfortunately, the lifetime of the R -axion is too long for the decay to take place before hitting the detector. In order for the two photons to be collimated enough, the mass of the R -axion should conservatively be below 200 MeV , and this is already in tension with the possibility that the decay is induced by dynamics at the TeV scale [14]. As in our case the decay is induced by dynamics at the few $\mathcal{O}(10 \text{ TeV})$ scale (unlike the resonance s , the R -axion has no trilinear coupling to the light scalar messengers, eq. (3.10)), there is no chance that it will be fast enough to give rise to the diphoton signal. This is the case except if the R -axion mass happens to be very close to the pion mass (or the η mass, but that value of the masses might be too large [14]) [51, 52]. In the latter case, a non-negligible mixing with the pion would allow the R -axion to decay as a pion, well before hitting the detector. In summary, the fate of the R -axion is either to contribute to the invisible width of s or, in a fine-tuned window for its mass, to contribute to the diphoton signal.

As for the relative size of the widths, the gg and $\gamma\gamma$ widths are suppressed by a loop factor, compared to eq. (5.1), but the aa width is suppressed by a higher scale, M^2 versus m_l^2 . The relative sizes of the widths then depends on the specific values of the parameters

⁸The R -axion decay into two gravitinos is very much suppressed in our region of parameter space [44, 45].

one considers, and is controlled by $\lambda M/m_l = g_{\text{eff}}/\lambda$,

$$\frac{\Gamma(s \rightarrow gg)}{\Gamma(s \rightarrow aa)} \approx \frac{\alpha_3^2}{36\pi^2} \left(\frac{g_{\text{eff}}}{\lambda}\right)^4 \bar{N}_3^2, \quad \frac{\Gamma(s \rightarrow \gamma\gamma)}{\Gamma(s \rightarrow aa)} \approx \frac{\alpha^2}{288\pi^2} \left(\frac{g_{\text{eff}}}{\lambda}\right)^4 \bar{N}_\gamma^2. \quad (5.2)$$

Assuming for simplicity that the R -axion mass is below 200 GeV, we have three regimes (assuming $N_\gamma \lesssim 34 N_3$).

- $80/\sqrt{\bar{N}_\gamma} \lesssim (g_{\text{eff}}/\lambda)$

In this regime, the decay in R -axions is subdominant to both the decay into gg and $\gamma\gamma$. Therefore, it does not affect the discussion in section 3.

- $14/\sqrt{\bar{N}_3} \lesssim (g_{\text{eff}}/\lambda) \lesssim 80/\sqrt{\bar{N}_\gamma}$

In this regime, the decay width in R -axions is larger than the decay width in $\gamma\gamma$, but not of the decay width in gg . Therefore, it does not affect the discussion in section 3, except in the fine-tuned window in which it mixes with the pion. In such a case, it gives the dominant contribution to the diphoton signal, and the $\sigma_{\gamma\gamma}$ rate determines $\Gamma(s \rightarrow aa)$,

$$\Gamma(s \rightarrow aa) \approx 0.3 \text{ MeV} \left(\frac{\sigma_{\gamma\gamma}}{6 \text{ fb}}\right) \Rightarrow M \approx 80 \text{ TeV} \left(\frac{6 \text{ fb}}{\sigma_{\gamma\gamma}}\right)^{1/2}. \quad (5.3)$$

- $(g_{\text{eff}}/\lambda) \lesssim 14/\sqrt{\bar{N}_3}$

In this regime, the decay width in R -axions is larger than both the decay widths in gg and $\gamma\gamma$. The diphoton signal is then suppressed, compared to what found in section 3, which should be avoided. Except in the fine-tuned region in which the R -axion decay into two photons is enhanced by the mixing with the pion, in which case the $\sigma_{\gamma\gamma}$ rate determines $\Gamma(s \rightarrow gg)$,

$$\Gamma(s \rightarrow gg) \approx 0.3 \text{ MeV} \left(\frac{\sigma_{\gamma\gamma}}{6 \text{ fb}}\right) \Rightarrow \frac{g_{\text{eff}}}{g_{\text{eff}}^*} \approx \frac{0.2}{\bar{N}_3} \left(\frac{m_l}{\text{TeV}}\right) \left(\frac{\sigma_{\gamma\gamma}}{6 \text{ fb}}\right)^{1/2}. \quad (5.4)$$

6 Sfermion masses and D -terms

In the previous section, we have ignored the MSSM sfermions. On the other hand, if the only contribution to their mass was the minimal gauge mediation two-loop contribution that follows from eq. (3.2), we would expect the colored sfermions to be lighter than the gluino, in which case they would play a role at least in forcing the whole spectrum to be heavier in order to pass the experimental bounds. In this section we show that i) it is indeed easy to split the spectrum and make the sfermions parametrically heavier than the gauginos, so that they do not play a role in the diphoton phenomenology, and ii) the model building ingredients needed to make them heavy modify the dynamics discussed in section 3, but have a minor impact on the conclusions.

In order to make the sfermions parametrically heavier than the gauginos it suffices to make both X and the sfermions charged under a (non-anomalous) $U(1)_X$ gauge factor. As

they couple to X , the messengers are then also charged under X . The vev of X breaks $U(1)_X$ and supersymmetry at the same time. As a consequence, the sfermions (and the scalar messenger, or “smessengers”) get tree-level soft masses from the $U(1)_X$ D -term, parametrically larger than the loop-induced gaugino masses. In the near critical regime, in which, as we will see, we still have $F \sim \lambda^2 M$, the D -term contribution to the soft mass of the scalar φ (sfermion or smessenger), m_φ^2 , is of the order of the heavy messenger scale,

$$m_\varphi^2 = q_\varphi g_X D \sim \frac{F^2}{M^2} \sim \lambda^2 M^2 = M_m^2, \quad (6.1)$$

where q_φ is the $U(1)_X$ charge of the scalar field, g_X the gauge coupling, and D the D -term. With M_m in the (10–100) TeV range, we are dealing with a simple realisation of the split supersymmetry spectrum [57–59].⁹

The large soft terms have a relevant impact on both sfermions and smessengers dynamics. The sfermions are too heavy to affect the diphoton phenomenology, as desired. In order for them not to be tachyonic, their $U(1)_X$ charges need to have the same sign as the D -term, say positive for definiteness. As $U(1)_X$ is non-anomalous, the supertrace must vanish, and the tree-level scalar soft masses must add up to zero. The positiveness of the MSSM sfermion masses hence forces negative soft mass terms for some scalars. This had long been considered as an obstacle to tree-level mediation of supersymmetry breaking in non-anomalous, renormalizable theories. But it is not: the messengers are anyway forced to have (overall) negative soft mass terms, as they couple to a positively charged field (see below). That does not make them tachyonic, as the negative soft mass term is compensated by the positive, supersymmetric, mass term. And their soft mass can compensate the positive sfermion soft masses. Such a class of models, in which the sfermions with negative soft masses needed to satisfy the supertrace constraint get a large, positive, supersymmetric mass term from $U(1)_X$ breaking and play the role of chiral messengers of minimal gauge mediation has been studied in refs. [60–62]. The compensation, i.e. anomaly cancellation, can arise automatically if the $U(1)_X$ is embedded in non-abelian gauge groups.

Let us now consider the impact of the new D -term contributions on the smessenger dynamics, and show that the conclusions obtained in section 3 are unchanged. Let $q_X = 1$ be the charge of X ,¹⁰ $-q$, $-\bar{q}$ the charges of Φ , $\bar{\Phi}$ (neglecting again the messenger flavour index i), with $q + \bar{q} = 1$, so that the total messenger soft mass, $m_\phi^2 + m_{\bar{\phi}}^2 = -g_X D$, is negative. The messengers are then chiral under $U(1)_X$, which “protects” their masses, in the same sense in which the electroweak symmetry “protects” the SM fermion masses. Eq. (3.8) becomes

$$-\mathcal{L}_{\text{mess}}^{(2)} = (\lambda M \psi \bar{\psi} + \text{h.c.}) + (\phi^\dagger, \bar{\phi}) \begin{pmatrix} \lambda^2 M^2 - q g_X D & \lambda F \\ \lambda F & \lambda^2 M^2 - \bar{q} g_X D \end{pmatrix} \begin{pmatrix} \phi \\ \bar{\phi}^\dagger \end{pmatrix}, \quad (6.2)$$

⁹A family-dependent assignment of $U(1)_X$ charges would give rise to a simple realisation of natural SUSY spectrum.

¹⁰Up to normalisation, $q_X = \pm 1$. If X is the dominant source of supersymmetry breaking, $D > 0$ is obtained for $q_X = 1$.

where all terms in the smessenger mass matrix are of the same order. In order to avoid tachyons, we need $\lambda^2 M^2 \geq qg_X D$ (by assumption larger than $\bar{q}g_X D$) and

$$\lambda^2 F^2 \leq M_\phi^2 M_\phi^2, \quad (6.3)$$

where

$$\begin{aligned} M_\phi^2 &= \lambda^2 M^2 - qg_X D, \\ M_\phi^2 &= \lambda^2 M^2 - \bar{q}g_X D. \end{aligned} \quad (6.4)$$

Near criticality (and a small smessenger mass $m_l^2 \ll M_m^2$) is obtained when the condition in eq. (6.3) is just satisfied, with F^2 just below the upper limit. Note that the near-critical regime cannot be associated to a fine-tuned cancellation in M_ϕ^2 , as that would imply $F \ll \lambda M^2$ and $D \sim F^2/M^2 \ll \lambda^2 M^2$. As a consequence, $F \sim \lambda M^2$. The heavy and light mass eigenstates ϕ_h, ϕ_l have now mass

$$m_{h,l}^2 = \frac{M_\phi^2 + M_\phi^2}{2} \pm \left[\left(\frac{M_\phi^2 - M_\phi^2}{2} \right)^2 + \lambda^2 F^2 \right]^{1/2} \quad (6.5)$$

and are given by

$$\begin{aligned} \phi &= \cos \theta \phi_h - \sin \theta \phi_l \\ \bar{\phi}^* &= \sin \theta \phi_h + \cos \theta \phi_l, \end{aligned} \quad \sin 2\theta = \frac{2\lambda F}{m_h^2 - m_l^2}. \quad (6.6)$$

The diphoton signal is not affected by the D -term contributions to the smessenger masses. The angle describing the mixing in the smessenger sector, now possibly different from $\pi/4$, does not enter the relevant trilinear interactions, which still have the form in eq. (3.10). The decay widths are therefore unchanged (for given m_l), in the near critical limit in which the light smessenger exchange dominates the diphoton signal. In particular, the effective coupling of the resonance to the light smessenger is still given by eq. (3.11).

On the other hand, the D -term has a mild effect on the relation of the gluino mass to the smessenger masses. We have in fact

$$\begin{aligned} M_a &= \frac{\alpha_a}{4\pi} \frac{F}{M} N_a g \left(\frac{m_l^2}{M_m^2}, \frac{m_h^2}{M_m^2} \right), \\ g(x_1, x_2) &= \frac{2}{x_1 - x_2} \left(\frac{x_1 \log x_1}{x_1 - 1} - \frac{x_2 \log x_2}{x_2 - 1} \right), \end{aligned} \quad (6.7)$$

and, in the near-critical regime,

$$M_a = \frac{\alpha_a}{4\pi} M_m N_a \frac{\sqrt{(1-qr)(1-\bar{q}r)}}{1-r} 2 \log(2-r), \quad \text{with } r = \frac{g_X D}{M_m^2}, \quad (6.8)$$

and $M_m = (g_{\text{eff}}/\lambda)m_l$, as before. For $r \rightarrow 0$ (and $\Delta M_3 = 0$), eq. (6.8) reproduces eq. (3.20). Numerically, for given m_l and g_{eff} , the presence of D -terms, i.e. of a non-zero r , gives only slightly lower values of gluino masses. We therefore conclude that in the discussion of section 3 the sfermions can be easily made heavy, without significantly modifying the conclusions about the possibility to fit the signal within the constraints. On the other

hand, the mixing in eq. (6.6), could induce a decay into Higgses $s \rightarrow hh$. Depending on the values of the relevant parameters, the $s \rightarrow hh$ decay width can be smaller, comparable or even larger than the $s \rightarrow gg$ width. In particular, the $s \rightarrow hh$ width can be suppressed by taking g_X small (but large enough to keep the sfermions heavy) or the Higgs mixing angle α close to $\pi/4$. Also, when Φ and $\bar{\Phi}$ have the same charge under $U(1)_X$, $q = \bar{q} = 1/2$, the messenger contribution to $s \rightarrow hh$ vanishes. This is the case in the $V + \bar{V}$ example of section 3.4, as both fields originate from the same adjoint.

7 Summary and outlook

We have revisited the possibility to associate the recently reported diphoton excess to the production of a sgoldstino of about 750 GeV. In this context, the new degree of freedom is not an ad hoc degree of freedom, it is ordered by the need to break supersymmetry, in the context of a theory with its own appeal; and the experiment measures the scale of mediation of supersymmetry breaking, which turns out to be very low, $\mathcal{O}(100 \text{ TeV})$ or less.

We assumed that supersymmetry breaking, and thus the sgoldstino resonance, is coupled to the MSSM fields through gauge mediation, which is appropriate at such low scales. The messenger superfields then provide the additional degrees of freedom needed for the decay and production of the resonance.

We showed that the experimental bounds on gaugino masses force the messenger scale M_m to be $\sim (10\text{--}100) \text{ TeV}$ and thus make the sgoldstino contribution to the diphoton excess unobservable, for a reasonable messenger content; except in a small region of the parameter space near the critical point beyond which the messengers get a vev, $F \approx \lambda M^2$.

The phenomenology in this thin, nearly-critical region drastically departs from the standard gauge mediation picture. One (or more) of the scalar messengers becomes much lighter than the heavy messenger scale. It can therefore lie at the TeV scale, as needed to account for the diphoton excess. At the same time, when the messenger becomes much lighter than M_m , its effective trilinear coupling gets enhanced by a factor $\lambda M_m/m_l$, where m_l is the light messenger mass, thus further helping to account for the excess. When the enhancement becomes very large, the system enters a strongly interacting regime. It is then not possible to further raise the gain through a larger hierarchy between the heavy and light messengers.

The IR non-perturbativity found at small $m_l \ll M_m$, associated to a large relevant coupling, is of quite a different nature than the usual UV non-perturbativity associated to irrelevant or marginal coupling. As an infrared effect, it does not give rise to Landau poles and it does not spoil the UV properties of supersymmetric theories. A quantitative analysis of the possibility to account for the diphoton effect showed a dichotomy between those two regimes. On the one hand, it is possible to account for the diphoton excess while avoiding the IR strong coupling by using a large enough set of messenger fields. This however forces Landau poles well below the GUT scale, and thus strong dynamics in the UV. On the other hand, it is possible to maintain the theory perturbative in the UV (up to the GUT scale) by having a lower number of messengers, but that forces a large trilinear coupling and induces strong dynamics in the IR.

The IR regime is particularly intriguing. It requires the system to be in a fine-tuned near-critical region, where the determinant of the scalar messenger mass matrix is small. In the context of the simplest possible structure of the hidden sector, $W_{\text{hidden}} = FX$, we observed that the near critical region is located on one end of the metastable flat direction associated to X , i.e. around the critical point, before the cascade to the supersymmetric minimum. The shape of the loop corrected effective potential along the flat direction is well known, it slowly pushes X towards the cascade. On the other hand, because of the non-perturbative regime arising there, the shape of the potential in the near critical region is not obvious. The obvious speculation is then that a metastable minimum could form in the near-critical region, thus providing a dynamical origin for the apparent fine-tuning we need, and for the origin of supersymmetry breaking. But this is of course just wishful thinking. In any case, an investigation of that region with non-perturbative methods would be welcome.

Back to phenomenology, we did not aim at accounting for a possibly large width of the resonance, relying of the presently uncertain experimental situation. In particular, the known interpretation of an apparent width in terms of the production of two resonances close in mass, identified with the scalar and pseudoscalar components of the sgoldstino, is not available here, as the pseudoscalar component has no (enhanced) trilinear coupling to the messengers.

In passing, we have commented on the role of a possible R -axion in the analysis of the diphoton excess. In this setup, the R -axion mass is in the ballpark of the pion mass. That is the window in which the decay of the sgoldstino into two R -axions, followed by the subsequent decay of each R -axion into two collimated photons, would mimic a diphoton signal. On the other hand, the lifetime of the R -axion would be too long for it to decay before the detector, except when a sizeable mixing with the pion arises.

Finally, we have shown that it is possible to give the sfermions a mass parametrically larger than the gauginos ones, so that they have no impact on our discussion, without altering our conclusions.

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A Partial widths

In this appendix we give the one loop expression for the partial decay widths of s and a into gg , $\gamma\gamma$, ZZ , $Z\gamma$, WW , neglecting the mass of the massive gauge bosons.

$$\Gamma(s \rightarrow gg) = m_s \frac{8\alpha_3^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{3,i} \sqrt{x_i} \left[P(x_i) + \frac{F(x_{i,l}) + F(x_{i,h}) - 2F(x_i)}{2} \right] \right|^2 \quad (\text{A.1})$$

$$\Gamma(a \rightarrow gg) = m_a \frac{8\alpha_3^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{3,i} \sqrt{x_i} P(x_i) \right|^2 \quad (\text{A.2})$$

$$\Gamma(s \rightarrow \gamma\gamma) = m_s \frac{\alpha^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{\gamma,i} \sqrt{x_i} \left[P(x_i) + \frac{F(x_{i,l}) + F(x_{i,h}) - 2F(x_i)}{2} \right] \right|^2 \quad (\text{A.3})$$

$$\Gamma(a \rightarrow \gamma\gamma) = m_a \frac{\alpha^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{\gamma,i} \sqrt{x_i} P(x_i) \right|^2 \quad (\text{A.4})$$

$$\Gamma(s \rightarrow ZZ) = m_s \frac{\alpha^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{Z,i} \sqrt{x_i} \left[P(x_i) + \frac{F(x_{i,l}) + F(x_{i,h}) - 2F(x_i)}{2} \right] \right|^2 \quad (\text{A.5})$$

$$\Gamma(a \rightarrow ZZ) = m_a \frac{\alpha^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{Z,i} \sqrt{x_i} P(x_i) \right|^2 \quad (\text{A.6})$$

$$\Gamma(s \rightarrow Z\gamma) = m_s \frac{2\alpha^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{Z\gamma,i} \sqrt{x_i} \left[P(x_i) + \frac{F(x_{i,l}) + F(x_{i,h}) - 2F(x_i)}{2} \right] \right|^2 \quad (\text{A.7})$$

$$\Gamma(a \rightarrow Z\gamma) = m_a \frac{2\alpha^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{Z\gamma,i} \sqrt{x_i} P(x_i) \right|^2 \quad (\text{A.8})$$

$$\Gamma(s \rightarrow WW) = m_s \frac{2\alpha_2^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{W,i} \sqrt{x_i} \left[P(x_i) + \frac{F(x_{i,l}) + F(x_{i,h}) - 2F(x_i)}{2} \right] \right|^2 \quad (\text{A.9})$$

$$\Gamma(a \rightarrow WW) = m_a \frac{2\alpha_2^2}{(4\pi)^3} \left| \sum_i \frac{\lambda_i}{\sqrt{2}} N_{W,i} \sqrt{x_i} P(x_i) \right|^2. \quad (\text{A.10})$$

In the above expressions, F is the (off-shell) scalar loop function of s , P is the (off-shell) fermion loop function of a , and the fermion loop function of s , S , has been expressed in terms of the previous two,

$$P(x) = \arctan^2 \frac{1}{\sqrt{x-1}}, \quad F(x) = xP(x) - 1, \quad S(x) = P(x) - F(x). \quad (\text{A.11})$$

The arguments of the loop functions are

$$x_i = 4 \frac{\lambda_i^2 M^2}{m_{s,a}^2}, \quad x_{ih,il} = 4 \frac{m_{ih,il}^2}{m_{s,a}^2}. \quad (\text{A.12})$$

Finally, the Dynkin coefficients are

$$\begin{aligned} N_\gamma &= \frac{5}{3} N_1 + N_2, \\ N_Z &= \frac{5}{3} \tan^2 \theta_W N_1 + \cot^2 \theta_W N_2, \\ N_{Z\gamma} &= \frac{5}{3} \tan \theta_W N_1 - \cot \theta_W N_2, \\ N_W &= N_2, \end{aligned} \quad (\text{A.13})$$

in terms of the SM Dynkin indices $N_{1,2,3}$.

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