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Threshold enhancement of diphoton resonances

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ABSTRACT

We revisit a mechanism to enhance the decay width of (pseudo-)scalar resonances to photon pairs when the process is mediated by loops of charged fermions produced near threshold. Motivated by the recent LHC data, indicating the presence of an excess in the diphoton spectrum at approximately 750 GeV, we illustrate this threshold enhancement mechanism in the case of a 750 GeV pseudoscalar boson A with a two-photon decay mediated by a charged and uncolored fermion having a mass at the $\frac{1}{2}M_A$ threshold and a small decay width, < 1 MeV. The implications of such a threshold enhancement are discussed in two explicit scenarios: i) the Minimal Supersymmetric Standard Model in which the A state is produced via the top quark mediated gluon fusion process and decays into photons predominantly through loops of charginos with masses close to $\frac{1}{2}M_A$ and ii) a two Higgs doublet model in which A is again produced by gluon fusion but decays into photons through loops of vector-like charged heavy leptons. In both these scenarios, while the mass of the charged fermion has to be adjusted to be extremely close to half of the A resonance mass, the small total widths are naturally obtained if only suppressed three-body decay channels occur. Finally, the implications of some of these scenarios for dark matter are discussed.

1. Introduction

There is presently considerable excitement in the particle physics community as the ATLAS and CMS collaborations have reported an excess in the data collected from LHC collisions at an energy of 13 TeV, corresponding to a possible resonance with a mass of 750 GeV decaying into two photons [1]. Assuming the resonance to be a scalar boson denoted Φ, the production cross section times the decay branching ratio for the final state $pp \to \Phi \to \gamma\gamma$ is very large, $= 6 \pm 2$ fb [2]. Such a rate is very difficult to accommodate in minimal versions of theories that are often considered to be attractive extensions of the Standard Model (SM). For example, it has been shown [3] that in its Minimal Supersymmetric extension (MSSM) [4], while there are additional Higgs bosons beyond the already observed one that can indeed be identified with the 750 GeV state, the diphoton rate cannot be generated using purely the MSSM particle content. Hence, the Φ resonance must be accompanied by additional massive charged particles to enhance the $\Phi\gamma\gamma$ decay amplitude and, eventually, the Φgg amplitude in the likely case where the resonance is produced via the gluon fusion mechanism, $gg \to \Phi$.

An interesting possibility would be that these additional particles are electrically charged and non-colored (generally vector-like) fermions that contribute only to the $\Phi \to \gamma\gamma$ decay; see e.g. Ref. [3]. However, in this specific case, the large enhancement of the $\Phi\gamma\gamma$ amplitude would require either *i*) several charged fermions, and/or *ii*) large electric charges, and/or *iii*) strong Yukawa couplings. All these requirements could unfortunately put perturbation theory under jeopardy [5]. The same problem occurs, although to a lesser extent, in certain scenarios where additional colored particles contribute to the Φgg vertex [2,3,5].

One means by which this problem could be alleviated would be to assume that Φ is a pseudoscalar state $\Phi \equiv A$ and that the charged fermions running in the $A\gamma\gamma$ loop have masses near the $m = \frac{1}{2}M_A$ kinematic threshold [3]. In this case, the form factor $A^{\Phi}_{1/2}$ [4,6] that characterizes the loop contributions of spin-$\frac{1}{2}$ fermions as functions of the scalar to fermion mass ratio (and which depends on the parity of the spin-zero state) becomes maximal and much larger than in the very heavy or very light fermion mass limits. Nevertheless, even in this particular case, the obtained $A^{A}_{1/2}$ form factor is not sufficient to explain the large diphoton rate

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in a minimal way and without endangering perturbation theory. We should note that in the context of the MSSM, even for masses close to M_A, the contribution of the two $\chi_1^\pm$ chargino states to the $A \to \gamma\gamma$ rate is too small as their couplings are very weak.

In this paper, still assuming a pseudoscalar resonance and a charged and uncolored fermion with a mass close to the $\frac{1}{2}M_A$ threshold, we invoke an additional mechanism to enhance the $A\gamma\gamma$ loop amplitude: the charged fermions will form S-wave (quasi) bound states resulting in a Coulomb singularity developing very close to this kinematic threshold [7–9]. This singularity is regulated by the total decay width of the charged fermion which, if it is very small, say $\Gamma \lesssim 1$ MeV, allows an enhancement of the $A\gamma\gamma$ amplitude by a large factor. Although the mechanism is in fact quite generic, we study its implications in the context of the potential 750 GeV resonance hinted at by the first 13 TeV LHC data. We show that with only one singly charged fermion having a reasonable Yukawa coupling to the resonance, one could generate a $A \to \gamma\gamma$ amplitude that is sufficiently large to accommodate the LHC diphoton signal. This interesting possibility will be studied in two specific examples.

We first reconsider the MSSM [3], assuming that the CP-odd A state corresponds to the 750 GeV diphoton resonance and has a strong top quark Yukawa coupling. This already allows for a significant cross section in the top induced $gg \to A$ process. The required $A \to \gamma\gamma$ decay rate is then generated by loops of charginos with a mass $m_{\chi_1^\pm} = \frac{1}{2}M_A$ for which the $A\gamma\gamma$ form-factor develops a Coulomb singularity that is regulated by a chargino width $\Gamma_{\chi_1^\pm} < 1$ MeV. Such a small decay width can be achieved naturally by imposing that the only possible decay mode, the one into the stable lightest neutralino and a fermion pair, occurs only at the three-body level and is strongly suppressed. One can then have a large threshold enhancement which easily explains the LHC diphoton data in this minimal supersymmetric model.

In a second scenario, we consider either the MSSM or a two Higgs doublet model (2HDM) [10] in which the A state still corresponds to the new resonance as discussed above but where the required $A \to \gamma\gamma$ decay rate is now generated by two vector-like doublets and singlets of heavy leptons [3]. The lightest charged lepton E has a mass very close to threshold $m_E = \frac{1}{2}M_A$ and the Coulomb singularity is again regulated by a small decay width that is obtained by assuming a lighter (possibly stable) neutral lepton N with a mass $m_E < m_N + M_W$ such that the only available decay mode is the suppressed three-body channel $E \to NW^* \to Nf\bar{f}'$. We then explore the regions in the parameters Γ_E and $m_E - m_N$ that allow us to obtain the enhancement factor which explains the ≈ 6 fb $gg \to A \to \gamma\gamma$ production rate at the LHC.

Finally, in both the MSSM with a stable lightest neutralino and in a 2HDM where the charged lepton E is accompanied by a stable neutral one N, we consider the tantalizing possibility that the neutral particles are viable dark matter candidates [11]. We determine the range of masses and couplings that would allow such a possibility, once the relevant experimental constraints from direct and indirect dark matter searches are imposed.

2. Threshold enhancement of the diphoton width

Let us begin by discussing the possibility of a threshold enhancement in the general context of a spin-zero CP-even H or CP-odd A state with two-photon decays induced by loops of fermions with color number N_f^c, electric charge e_f and Yukawa couplings $\lambda_{\Phi ff}$ when normalized to their SM-like values, $\lambda_{\Phi ff}^{\rm SM} = \sqrt{2}m_f/v$ with $v = 246$ GeV. The two-photon partial decay widths read [6,4]

$$\Gamma(\Phi \to \gamma\gamma) = \frac{G_\mu \alpha^2 M_\Phi^3}{128\sqrt{2}\pi^3}\left|\sum_f N_f^c e_f^2 \lambda_{\Phi ff} A_{1/2}^\Phi(\tau_f)\right|^2 \quad (1)$$

where α is the QED fine structure constant, $\alpha = e^2/4\pi \approx 1/128$ at a scale M_Φ and G_μ is the Fermi constant. The form factors $A_{1/2}^\Phi(\tau_f)$ which depend on the variable $\tau_f = M_\Phi^2/4m_f^2$ are given, for the scalar and the pseudoscalar cases, by

$$A_{1/2}^H = 2\left[\tau_f + (\tau_f - 1)f(\tau_f)\right]\tau_f^{-2}, \quad A_{1/2}^A = 2\tau_f^{-1}f(\tau_f), \quad (2)$$

$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau} & \text{for } \tau \le 1, \\ -\dfrac{1}{4}\left[\log\dfrac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2 & \text{for } \tau > 1. \end{cases} \quad (3)$$

The amplitudes are real for Φ masses below the $M_\Phi = 2m_f$ kinematical threshold and develop an imaginary part above. When the loop fermion is much heavier than the Φ state, $m_f \to \infty$, one obtains $A_{1/2}^H = \frac{4}{3}$ and $A_{1/2}^A = 2$, while in the opposite limit, $m_f \to 0$, one has $A_{1/2}^\Phi \to 0$.

The maximal values of the form factors are attained near the mass threshold $m_f = \frac{1}{2}M_\Phi$ where one has: $\mathrm{Re}(A_{1/2}^H) \approx 2$ and $\mathrm{Re}(A_{1/2}^A) \approx \frac{1}{2}\pi^2 \approx 5$ for the real parts and $\mathrm{Im}(A_{1/2}^\Phi) \approx 0$. Hence, near threshold, the form-factor $A_{1/2}^\Phi$ is much larger for a CP-odd state and we will thus concentrate on this case in the rest of the discussion. Furthermore, we will only consider the case where color-neutral fermions contribute in the loops: heavy quarks would also contribute to the Agg loop-induced coupling[1] and would generate unacceptably large production rates in the situations which will be considered here (like in the $pp \to t\bar{t}$ process for instance [15]).

Nevertheless, the expressions Eqs. (2)–(3) above do not accurately describe the threshold region $m_f \approx \frac{1}{2}M_\Phi$ for the $A\gamma\gamma$ form factor. Indeed, for fermion masses just above but very close to threshold, a Coulomb singularity develops due to the fermions forming S-wave (quasi) bound states [16,17]. This can be taken into account, in a non-relativistic approach, by re-writing the form factor close to threshold as [8]

$$\tilde{A}_{1/2}^A = a + b \times G(0,0;E_f + i\Gamma_f), \quad (4)$$

where, to leading order, one has $a = \frac{1}{2}\pi^2$ and $b = 8\pi^2/m_f^2$ for the real and imaginary parts, $E_f = M_A - 2m_f$ for the distance from the threshold region and Γ_f is the total decay width of the fermion f running in the loop. Here a and b are the perturbatively calculable coefficients obtained from matching the non-relativistic theory to the full theory. $G(0,0;E_f)$ is the S-wave Green's function of the non-relativistic Schrödinger equation in the presence of a Coulomb potential $V(r) = -\alpha/r$. The fermion decay width Γ_f is introduced in order to regulate the Coulomb singularity in the Green's function with real and imaginary parts [9]

$$\mathrm{Re}\,G(0,0;E_f + i\Gamma_f) = -\frac{m_f p_-}{4\pi} + \frac{m_f p_0}{4\pi}\log\frac{m_f^2 D^2}{p_+^2 + p_-^2} + \frac{m_f p_0^2}{2\pi}\sum_{n=1}^{\infty}\frac{p_- - p_n}{n^2[(p_- - p_n)^2 + p_+^2]}, \quad (5)$$

[1] In principle, one expects these quarks to be rather heavy from LHC direct searches, $m_Q \gtrsim 800$ GeV [12] and hence beyond the $m_Q = \frac{1}{2}M_\Phi$ threshold. Nevertheless such a configuration could be possible in some special cases where bound states can form; see e.g. Ref. [13]. In the case where the resonance also couples to top quarks, there would also be a significant enhancement of the $A\gamma\gamma$ amplitude near the $M_\Phi \approx 2m_t$ threshold [8,14], but it is negligible in practice since a 750 GeV resonance is far from this configuration.

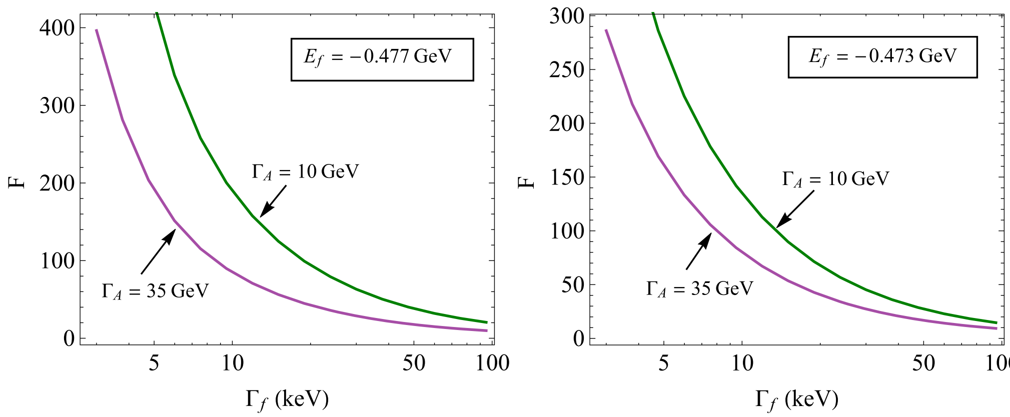

Fig. 1. The enhancement factor F of the cross section on including the threshold corrections as a function of the fermion total width Γ_f for two choices of $E_f = M_A - 2m_f$ of 0.477 GeV (left) and 0.473 GeV (right) and two values of the A total width Γ_A as indicated.

$$\mathrm{Im}\,G(0,0;E_f+i\Gamma_f) = -\frac{m_f\,p_+}{4\pi} + \frac{m_f\,p_0}{2\pi}\arctan\frac{p_+}{p_-} + \frac{m_f\,p_0^2}{2\pi}\sum_{n=1}^{\infty}\frac{p_+}{n^2[(p_- - p_n)^2 + p_+^2]}, \quad (6)$$

where $p_\pm = (\frac{1}{2}m_f(\sqrt{E_f^2+\Gamma_f^2} \pm E_f))^{1/2}$, $p_n = p_0/n$, $p_0 = \frac{1}{2}m_f\alpha$ and D is a renormalization constant which we set to unity in the following as this will only affect our results at higher orders[2] [8]. The three terms in the above expressions correspond to the lowest order contribution, a single Coulombic photon exchange and a sum over contributions involving the exchange of $n+1$ Coulombic photons. The position of the first pole in E_f can be obtained by inspecting the denominator of the $n=1$ contribution to the last terms of the equations above. Although the sum in n runs from 1 to ∞, the sum converges rather quickly and, in reality, it is sufficient for our purposes to calculate up to $n = 100$.

One should note that while the pole is present in both the real and imaginary parts, the numerator p_+ for the imaginary part is suppressed compared to that of the real part $p_- - p_0/n$ in the vicinity of the pole. The large enhancement of the two-photon form factor is therefore obtained from the real part of the Green's function.

As mentioned earlier, the bound-state formation results in poles in the form factor at energies just below the threshold for pair production of the fermions, regulated by the width of the fermions. The position of these poles in terms of the energy E_f as well as the size of the enhancement therefore depend on the size of the fermion decay width and the coupling to the photon.

Besides the fact that the dominant contribution to the enhancement is from the real and not the imaginary part of the form factor, one should note that in the CP-even scalar case, the $H\gamma\gamma$ form factor is P-wave and highly suppressed at the threshold; therefore, the bound state formation can be neglected as in this case the possible enhancement is negligible.

In the absence of threshold enhancement the diphoton spectrum can be well-described by a standard Breit–Wigner distribution centered around the resonance mass M_A and with a width that is equal to its total width Γ_A. The threshold enhancement mechanism, on the other hand, is only efficient over a small energy range close to the fermion pair-production threshold which we choose to coincide with the peak of the Breit–Wigner distribution. The modification of the *total* diphoton cross section therefore is affected by the fact that the width of M_A is large in comparison to this narrow range.[3] Assuming for simplicity that the gluon parton distribution functions remain constant over the width of the Breit–Wigner, an approximation which is reasonable since we will not be considering widths larger than a few tens of GeV, and ignoring normalization factors which cancel, the modification of the total diphoton cross section can be quantified by defining an enhancement factor F as

$$F = \int_{M_A-\Delta E}^{M_A+\Delta E} d\hat{s}\,\sqrt{\hat{s}}\,\frac{\left|\tilde{A}^A_{1/2}\left(\frac{\hat{s}}{4m_f^2}\right)\right|^2}{(\hat{s}-M_A^2)^2 + M_A^2\Gamma_A^2} \Bigg/ \int_{M_A-\Delta E}^{M_A+\Delta E} d\hat{s}\,\sqrt{\hat{s}}\,\frac{\left|A^A_{1/2}\left(\frac{\hat{s}}{4m_f^2}\right)\right|^2}{(\hat{s}-M_A^2)^2 + M_A^2\Gamma_A^2}, \quad (7)$$

i.e. as the convolution of the enhanced form factor with the Breit–Wigner distribution divided by the corresponding quantity in the un-enhanced case. The integration region in Eq. (7) extends from $M_A - \Delta E$ to $M_A + \Delta E$, where ΔE is some appropriate energy interval allowing us to include the bulk of the diphoton events, thus ensuring that F actually describes the *total* cross section enhancement. In our numerical analysis we take $\Delta E = 20$ GeV, which corresponds to half of the bin width exhibiting the largest excess of diphoton events. In Fig. 1, we display the quantity F as a function of the fermion width Γ_f for a resonance mass $M_A = 750$ GeV, two values of the resonance width $\Gamma_A = 10$, 35 GeV and for two values of the mass difference $E_f = M_A - 2m_f$.

As we shall see in the following section, a factor F of 400 seen for $\Gamma_f \lesssim 3$ keV for $\Gamma_A = 35$ GeV could explain the LHC results in the MSSM case, whereas in our 2HDM scenario a factor $F \approx 50$ would be sufficient, requiring $\Gamma_f \lesssim 15$ keV. Note that it is possible to get the desired enhancement for larger values of Γ_f if smaller Γ_A is chosen. The enhancement factor is of course dependent on the choice of E_f or conversely m_f, i.e. whether or not E_f corresponds exactly to the position of the pole. This is the Achilles heel of our scenario as some "fine-tuning" is thus necessary.

[2] In principle, one could calculate higher order corrections to the coefficients a and b; however, these are not needed for this preliminary study as, in particular for the QED case, they should be highly suppressed.

[3] We are grateful to the referee for valuable comments on this issue.

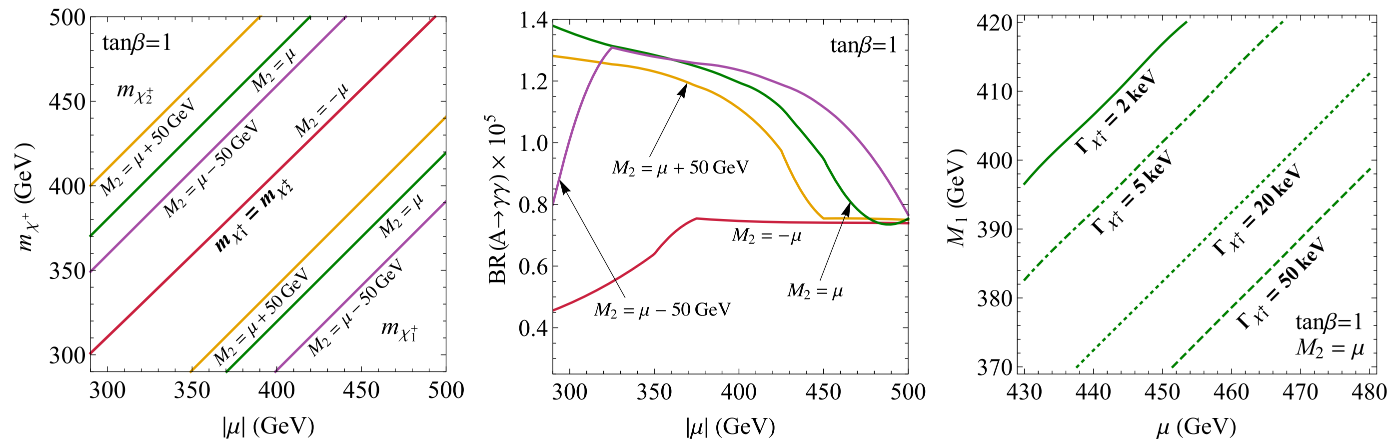

Fig. 2. The two chargino masses (left) and the branching ratio BR($A \to \gamma\gamma$) when the chargino contributions are included (center) as a function of μ for the values of $M_2 = \mu, \mu + 50$ GeV and $\mu - 50$ GeV (for $\mu > 0$) and $M_2 = -\mu$ (for $\mu < 0$). Contours in the $[\mu, M_1]$ plane for which we obtain a total chargino width of $\Gamma(\chi_1^\pm \to \chi_1^0 f \bar{f}') = 2, 5, 20$ and 50 keV (right). The MSSM with $\tan\beta = 1$ and $M_A = 750$ GeV (and heavy sfermions) is assumed in all cases.

3. Implications for diphoton resonance scenarii

Let us now discuss the implications of this possible threshold enhancement in some scenarii for the 750 GeV $\Phi = A$ resonance, starting with the plain MSSM scenario.

In the MSSM, two Higgs doublets Φ_u and Φ_d are required to break the electroweak symmetry leading to five physical states, two CP-even h and H, a CP-odd A and two charged $H^\pm$ bosons [4]. In the so-called decoupling limit, $M_A \gg M_Z$, the lighter h state is the Higgs boson observed at the LHC in 2012 and subsequently determined to have SM-like properties, while the Φ resonance is a superposition of the heavier neutral CP-even H and CP-odd A that are nearly degenerate in mass (as is also the case of the charged Higgs boson). The two states have zero tree-level couplings to vector bosons and similar couplings to fermions. The latter are controlled by the ratio of vacuum expectation values $\tan\beta = v_u/v_d$ which is the only relevant parameter in the Higgs sector of the model besides $M_A \approx M_H$. For $\tan\beta \approx 1$, the only important Yukawa coupling is the one of the top quark, $y_t = \sqrt{2}m_t/(v\tan\beta) \approx 1$.

At the LHC, the H/A states are mainly produced in the $gg \to \Phi$ fusion mechanism that is mediated by a top quark loop with an amplitude that is given by an expression similar to that of Eqs. (1)–(3) except for some color factors and the replacement of α with α_s [6]. The cross sections are such that $\sigma(gg \to A) \approx 2\sigma(gg \to H)$ and for $M_\Phi \approx 750$ GeV and $\tan\beta \approx 1$, one obtains $\sigma(A+H) \approx 2$ pb at the $\sqrt{s} = 13$ TeV LHC [18]. The $\Phi = H/A$ states will then mainly decay into top quark pairs with partial ($\approx$ total) widths that are of order $\Gamma_\Phi \approx 35$ GeV. Concentrating on the pseudoscalar A resonance, if the two-photon decay is generated by the top quark loop only, the branching ratio for the relevant inputs is BR($A \to \gamma\gamma) \approx 7 \times 10^{-6}$ [19]. One thus has a resonance production times decay rate of about $\sigma(gg \to A) \times$ BR($A \to \gamma\gamma) \approx 10^{-2}$ fb. Note that smaller widths can also be achieved with a similar result by increasing $\tan\beta$; the decrease in the production cross section will be compensated by an increase in the branching ratio. For $\tan\beta = 3$ the total width Γ_A is around 5 GeV.

In Ref. [3], the possible loop contributions of the supersymmetric particles have been analyzed and found to be far too small to explain the large two-photon decay rate. Here, we will reconsider the chargino loop contribution to $A \to \gamma\gamma$ in light of the possible threshold enhancement discussed above. These contributions are briefly summarized below.

The general chargino mass matrix, in terms of the wino and higgsino mass parameters M_2 and μ in the limit $\tan\beta \approx 1$ in which we specialize, is given by

$$\mathcal{M}_C = \begin{bmatrix} M_2 & \sqrt{2}M_W \sin\beta \\ \sqrt{2}M_W \cos\beta & \mu \end{bmatrix} \overset{\tan\beta=1}{\to} \begin{bmatrix} M_2 & M_W \\ M_W & \mu \end{bmatrix} \quad (8)$$

The two physical chargino states $\chi_1^\pm, \chi_2^\pm$ and their masses are determined through unitary matrices U and V defined and given by (σ_3 is the Pauli matrix with diagonal values $+1, -1$)

$$U^* \mathcal{M}_C V^{-1} \; : \; V = \mathcal{O}_+ \, ,$$
$$U = \begin{cases} \mathcal{O}_- & \text{if } \det \mathcal{M}_C > 0 \\ \sigma_3 \mathcal{O}_- & \text{if } \det \mathcal{M}_C < 0 \end{cases} , \quad \mathcal{O}_\pm = \begin{bmatrix} \cos\theta_\pm & \sin\theta_\pm \\ -\sin\theta_\pm & \cos\theta_\pm \end{bmatrix} . \quad (9)$$

The coupling of the pseudoscalar A boson to pairs of the same chargino is given by

$$g_{A\chi_i^-\chi_i^+} = -\frac{e}{\sqrt{2}\sin\theta_W}\left[\sin\beta V_{i1}U_{i2} + \cos\beta V_{i2}U_{i1}\right]$$
$$\overset{\tan\beta=1}{\to} -\frac{e}{2\sin\theta_W}\left[V_{i1}U_{i2} + V_{i2}U_{i1}\right] \quad (10)$$

As can be seen, this coupling is maximal for equal admixtures of higgsinos and winos, $|\mu| \approx M_2$. Taking the limit $|\mu| = M_2 \gg M_W$ for simplicity and still $\tan\beta = 1$ (a choice which maximizes the A production cross section), the matrix Eq. (8) is easy to diagonalize and $\theta_\pm = \pi/4$. For $\mu \geq 0$, one obtains for the masses of the charginos and their couplings to the A state

$$\mu > 0 : m_{\chi_{1,2}^\pm} \simeq \mu \mp M_W \, , \; g_{A\chi_1^-\chi_1^+} \simeq -g_{A\chi_2^-\chi_2^+} \simeq e/(4\sin\theta_W) \quad (11)$$

For large μ values, the two states $\chi_{1,2}^\pm$ have thus masses that are close to each other and couplings to A of opposite sign; their loop contributions to the $A\gamma\gamma$ amplitude will therefore interfere destructively with each other reducing the $\chi^\pm$ impact in $A \to \gamma\gamma$ decays.

In turn, for $\mu < 0$, one obtains for the masses and couplings in the limit $|\mu| \approx M_2 \gg M_W$,

$$\mu < 0 : m_{\chi_{1,2}^\pm} \simeq |\mu| \, , \; g_{A\chi_1^-\chi_1^+} \simeq -g_{A\chi_2^-\chi_2^+} = \mathcal{O}(M_W^2/\mu^2) \quad (12)$$

Hence, the two charginos are nearly degenerate in mass and have suppressed couplings to the A state. Some numerical examples [19] for the two options of sign(μ) are shown in Fig. 2 where, in the left-hand side, we display the chargino masses for the four possibilities $\mu = M_2, M_2 \pm 50$ GeV when $\mu \geq 0$ and $\mu = -M_2$ for $\mu < 0$. Indeed, the masses behave as described in Eqs. (11) and (12) for the $\mu > 0$ and $\mu < 0$ cases respectively. In the central frame, we show the branching ratio BR($A \to \gamma\gamma$) when the (un-enhanced) chargino loop contributions are included. As can be seen, the contributions are not the largest for $m_{\chi_1^\pm} = 375$ GeV, as one would naïvely expect since the form factor $A_{1/2}^A \approx \frac{1}{2}\pi^2$ is maximal. This effect is due to the negative interference of the two chargino loops for $\mu > 0$ and the small Higgs couplings to charginos in the $\mu < 0$ case.

We now focus on the case $\mu = M_2$ with $\mu > 0$ and describe the spectrum for $\tan\beta = 1$ when the lightest chargino is at the $m_{\chi_1^\pm} = \frac{1}{2}M_A = 375$ GeV threshold (which, here, occurs for $\mu =$ 455 GeV). We would have $m_{\chi_2^0} \approx m_{\chi_1^\pm} \approx \mu - M_W$, $m_{\chi_3^0} \approx \mu$ and $m_{\chi_4^0} \approx \mu + M_W$ and the lightest neutralino mass can then be chosen via the remaining input that enters the chargino–neutralino sector: the bino mass parameter M_1. Here, a careful choice ensures that the total width for the chargino $\chi_1^\pm$ is small, providing the threshold enhancement factor of the $A\chi_1^\pm\chi_1^\mp$ loop form-factor needed to explain the LHC data. Indeed, in the R-parity conserving scenario that we consider here, the only possible decay of the chargino $\chi_1^\pm$ will be into the lightest neutralino χ_1^0 (which is stable) and a W boson. If the mass splitting $m_{\chi_1^\pm} - m_{\chi_1^0}$ is small, the W boson is off-shell and decays into light fermions through the three-body decay $\chi_1^\pm \to \chi_1^0 W^* \to \chi_1^0 f\bar{f}'$, which has a very small partial (= total) width.

In the right panel of Fig. 2 we show contours of $\chi_1^\pm$ total decay width $\Gamma_{\chi_1^\pm} = 2$, 5, 20 and 50 keV, in the $[\mu, M_1]$ plane assuming $M_2 = \mu$ and $\tan\beta = 1$. Is it clear that it is possible to simultaneously attain $m_{\chi_1^\pm} = 375$ GeV and a very small width, $\Gamma_{\chi_1^\pm} \lesssim 2$ keV, which allows a sufficient enhancement of the $\chi_1^\pm$ loop contribution to the $A\chi_1^\pm\chi_1^\mp$ amplitude to reproduce the diphoton rate measured at the LHC. Hence, the situation is not desperate in the MSSM and there is a way to explain the properties of the diphoton resonance in this context.

Let us now turn to the case of a two-Higgs doublet model (2HDM) [10], again identifying the pseudoscalar state A with the 750 GeV resonance. The phenomenology of A is exactly the same as in the MSSM, in particular in the 2HDM alignment limit in which the lighter h state is SM-like and assuming the charged $H^\pm$ boson to be heavy enough. The A production mode at the LHC is the same as discussed above, but for the two-photon decay the only contribution will be that coming from top quark loops which, much like in the MSSM case, is again too small: an enhancement factor of at least ≈ 400 is required to obtain the resonance cross section of $\sigma(gg \to A \to \gamma\gamma) \approx 6 \pm 2$ fb measured at the LHC. Part of this enhancement can be obtained by introducing a doublet and two singlets of heavy vector-like leptons[4]

$$L_{L/R} = \begin{pmatrix} N \\ E \end{pmatrix}_{L/R}, \; E'_{L/R}, \; N'_{L/R}, \tag{13}$$

with the minimal Lagrangian describing their Yukawa couplings in the interaction basis

$$-\mathcal{L}_Y = \left\{ \frac{y_L^E}{\sqrt{2}} \bar{L}_L \Phi_d E'_R + \frac{y_L^N}{\sqrt{2}} \bar{L}_L \Phi_u N'_R + \mathrm{L} \leftrightarrow \mathrm{R} + m_L \bar{L}_L L_R \right. \\ \left. + m_N \overline{N'}_L N'_R + m_E \bar{E}'_L E'_R \right\} + \text{h.c.} \tag{14}$$

The Yukawa terms will result in mixing between the doublet and singlets, with the mixing matrix of the neutral/charged leptons taking the form:

$$\mathcal{M}_N = \begin{pmatrix} m_N & \frac{1}{\sqrt{2}} y_L^N v_u \\ \frac{1}{\sqrt{2}} y_L^N v_u & m_L \end{pmatrix}, \quad \mathcal{M}_E = \begin{pmatrix} m_L & \frac{1}{\sqrt{2}} y_L^E v_d \\ \frac{1}{\sqrt{2}} y_L^E v_d & m_E \end{pmatrix}. \tag{15}$$

On diagonalizing these matrices with angles $\theta_{N,E}$, the mass eigenstates can be written as

$$N_1 = \cos\theta_N N' + \sin\theta_N N\,, \; N_2 = \cos\theta_N N - \sin\theta_N N'\,,$$
$$\tan 2\theta_N = \sqrt{2} y_L^N v_u/(m_L - m_N),$$
$$E_1 = \cos\theta_E E + \sin\theta_E E'\,, \; E_2 = \cos\theta_E E' - \sin\theta_E E\,,$$
$$\tan 2\theta_E = \sqrt{2} y_L^E v_d/(m_E - m_L).$$

In light of the data on the diphoton resonance, one then assumes that the lepton E_1 has a mass $m_{E_1} \approx 375$ GeV and a Yukawa coupling $y_L^E \approx 2$, a value that is slightly below the perturbative limit. This allows an initial enhancement of the form factor $A_{1/2}^A$ of the $A \to \gamma\gamma$ amplitude yielding, for a total width $\Gamma_A \sim 35$ GeV, a branching ratio BR($A \to \gamma\gamma$) $\approx 4 \times 10^{-5}$. Then, to arrive at the observed σ for the diphoton rate, a further enhancement factor $F \sim 50$ is needed, i.e. $\Gamma_f \lesssim 15$ keV. In the original scenarios, see e.g. Ref. [3], several replicas of the above spectrum were required, leading to a model that is not entirely minimal. This additional factor can be now generated by the threshold enhancement of $A_{1/2}^A$ as discussed earlier.

Indeed, if one assumes $m_{E_1} = \frac{1}{2}M_A$ (up to a few MeV) and a small E_1 total decay width $\Gamma_{E_1} \ll 1$ MeV, an order of magnitude enhancement of the $A \to \gamma\gamma$ amplitude can be obtained with the minimal lepton spectrum of Eq. (13). The small width Γ_{E_1} can be obtained simply by ensuring that the mass difference $m_{E_1} - m_{N_1}$ is small and positive (this near mass degeneracy is required anyway in order to comply with precision electroweak data [3]). This ensures that the only possible E_1 decay is the three-body mode $E_1 \to N_1 W^* \to N_1 f\bar{f}'$ which requires a highly virtual W boson, strongly suppressing the decay width.

These decays of the heavy leptons have been discussed in Ref. [20] and using the relevant formulae for the three-body $E_1 \to N_1 W^* \to N_1 f\bar{f}'$ channel provided in the papers above, we show in the left-hand side of Fig. 3, the partial decay width (which in the absence of fermion mixing corresponds to the total width), $\Gamma_{E_1} = \Gamma(E_1 \to N_1 W^* \to N_1 f\bar{f}')$, as a function of $m_{E_1} - m_{N_1}$, assuming $\cos\theta_E = 1$. As can be seen a small width of about $\Gamma_{E_1} \approx 1$ keV can be obtained for a 50 GeV mass difference when $\sin\theta_N \simeq 0.033$. Note also that the presence of the light neutral state N_1 does not affect the total width of A, which is dominated by the $A \to \bar{t}t$ contribution, as the ratio of the two partial widths,

$$\frac{\Gamma(A \to N_1 N_1)}{\Gamma(A \to t\bar{t})} = \frac{1}{3} \frac{y_L^{N2} s_{\theta_N}^2 c_{\theta_N}^2}{y_t^2} \frac{\sqrt{1 - 4\frac{m_{N_1}^2}{m_A^2}}}{\sqrt{1 - 4\frac{m_t^2}{m_A^2}}}, \tag{16}$$

is of $\mathcal{O}(10^{-4})$ for $\sin\theta_N \simeq 0.033$ and $m_{N_1} \sim 325$ GeV.

On the right-hand side of Fig. 3, the mixing angle $\sin\theta_N$ is shown as a function of the mass difference $m_{E_1} - m_{N_1}$ for various values of the width Γ_{E_1}. From Fig. 1, we have seen that for widths below ~ 15 keV the desired enhancement factor can be obtained. Therefore, a mass difference in the range $m_{E_1} - m_{N_1} = 40$–80 GeV could easily lead to the enhancement factor needed to explain the LHC diphoton data, assuming that the A resonance is produced via gluon fusion.

[4] This is needed in order, first to cancel the chiral anomalies and second, to arrange that the lightest Higgs coupling to two photons, which is measured to be SM-like, is not significantly altered; see Ref. [3].

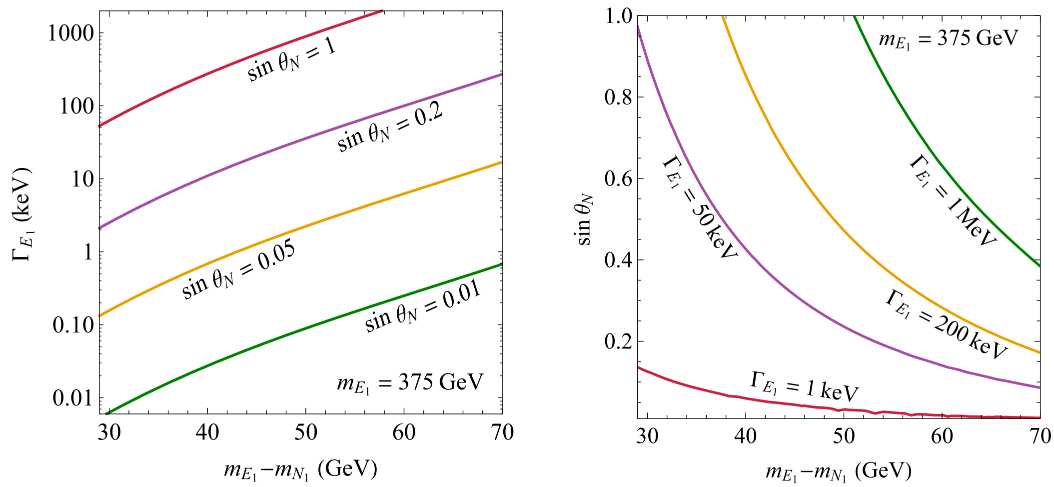

Fig. 3. Left: the sum of partial three-body decay widths $\Gamma(E_1 \to N_1 W^* \to N_1 f \bar{f}')$ in keV as a function of the mass difference $m_{E_1} - m_{N_1}$ for the values of $\sin\theta_N$ as shown. Right: the mixing angle between the doublet and singlet as a function of the mass difference $m_{E_1} - m_{N_1}$ which results in the width Γ_{E_1} as indicated.

4. Implications for dark matter

Since in both the scenarios we have studied, the lightest non-SM particles (the lightest neutralino and the mostly singlet lightest heavy neutrino) are charge and color-neutral, it is tempting to examine their viability as dark matter candidates. In order to simplify the discussion we assume as usual some discrete symmetry (R-parity in the MSSM and a $\mathcal{Z}_2$ symmetry in the 2HDM case under which the heavy leptons are even whereas all other fields are odd) that renders the lightest new state completely stable. This assumption has actually already, explicitly or implicitly, been made in the previous sections.

We begin with the 2HDM scenario, where the situation turns out to be more straightforward. Focusing on the regime where the dark matter annihilation is mediated by the s-channel exchange of the pseudoscalar A state and including all relevant interactions between the dark matter particles and the SM ones, notably those mediated by the Z boson, the relevant part of the Lagrangian Eq. (14) in terms of mass eigenstates can be written as

$$\mathcal{L} = \mathcal{L}_{\rm SM} - i\frac{y_L^N}{\sqrt{2}} s_{\theta_N} c_{\theta_N} A \bar{N}_1 \gamma^5 N_1 - i\frac{y_t}{\sqrt{2}} A\bar{t}\gamma^5 t - i\frac{y_b}{\sqrt{2}} A\bar{b}\gamma^5 b + \frac{e}{2 s_W c_W} s^2_{\theta_N} \bar{N}_1 \gamma^\mu N_1 Z_\mu \tag{17}$$

where $s_{\theta_N} \equiv \sin\theta_N$, and similarly for the charged sector. Note that we have ignored A couplings to light fermions since their contributions, being Yukawa suppressed, are much smaller than those of top and bottom quarks as well as those mediated by the Z-boson. Since, however, the mass splitting between the E_1 and N_1 states is relatively small, coannihilation processes could become important. These can be described by the Lagrangian

$$\mathcal{L}_{\rm coann} = \frac{g_2}{\sqrt{2}} s_{\theta_N} c_{\theta_E} \left(\bar{E}_1 \gamma^\mu N_1 W^-_\mu + h.c.\right) + \frac{g_2}{2 c_W} c^2_{\theta_E} \bar{E}_1 \gamma^\mu E_1 Z_\mu + i\frac{y_L^E}{\sqrt{2}} s_{\theta_E} c_{\theta_E} A \bar{E}_1 \gamma^5 E_1 \tag{18}$$

where for simplicity we assume that the charged scalars are heavy. To study the dark matter aspects of the model, we have implemented the Lagrangians Eqs. (17) and (18) in the public code MicrOMEGAs [21] with the help of the FeynRules package [22].

Our results are depicted in Fig. 4, where we highlight the $(m_{N_1}, \sin\theta_N)$ combinations for which the latest limits on dark matter abundance from the Planck mission [23] can be satisfied

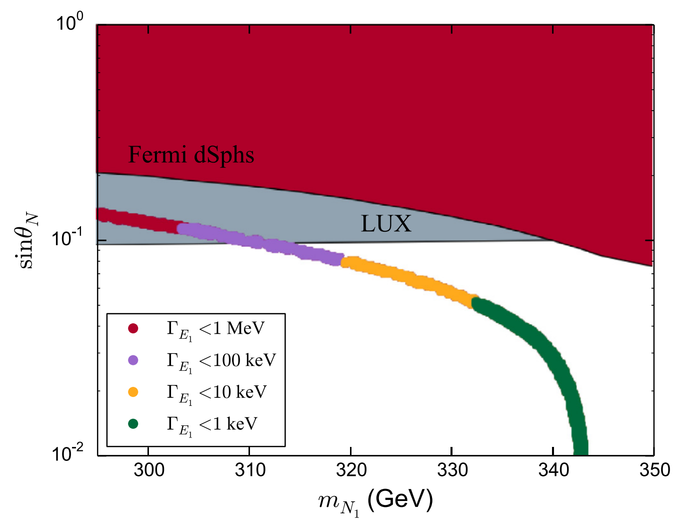

Fig. 4. Mixing angle versus the mass of the lightest heavy neutrino for which the Planck bound on the dark matter abundance in the universe is satisfied (colored band) for different widths of the lightest heavy electron, Γ_{E_1}. In the gray shaded region, the predicted spin-independent scattering cross-section off nucleons is in conflict with the latest LUX limits. The red shaded region depicts the limits from the Fermi satellite searches for dark matter annihilation – induced continuum gamma rays in dSphs. (For interpretation of the references to color in this figure, the reader is referred to the web version of this article.)

according to standard thermal freeze-out. The different colorings correspond to different ranges for the predicted width of the lightest heavy electron E_1 from the decay mode $E_1 \to N_1 W^* \to N_1 f \bar{f}'$. The other parameters entering the Lagrangian Eq. (17) have been set to the values $y_L^N = y_t = 1$, while for the charged sector parameters we have chosen $m_{E_1} = 375.003$ GeV, $y_L^E = 1$ and $s_{\theta_E} = 0.1$. The heavier neutrino mass m_{N_2} has been set to a large value in this analysis. The abrupt decrease in the required values of $\sin\theta_N$ in order to satisfy the latest Planck limits appearing around $m_{N_1} \sim 335$ GeV is due to coannihilation processes becoming important. In fact, by inspecting Fig. 4 one might think that the relic density constraint cannot be satisfied for $m_{N_1} \gtrsim 345$ GeV. It should however be kept in mind that these results also depend on the choice of y_L^E and s_{θ_E} and that the correct relic density can be easily obtained for larger N_1 masses by decreasing the former and/or increasing the latter. Such a choice would, nonetheless, be rather detrimental to the model's capacity to reproduce the LHC diphoton excess which is maximized for large values of y_L^E and small values of s_{θ_E}. In this sense, the parameter space depicted in Fig. 4 constitutes a reasonable compromise between the requirements for the model to reproduce the diphoton excess and to obtain the correct dark matter abundance in the Universe.

A well-known constraint on dark matter scenarios involving vector-like couplings to the Z boson comes from direct detection; see for example Ref. [24]. We have computed the predicted spin-independent scattering cross section off nucleons[5] and compared it to the updated analysis performed by the LUX collaboration [25]. The excluded regions of parameter space are depicted by the gray-shaded area of Fig. 4. As expected, the direct detection limits restrict the mixing of the singlet and doublet heavy neutrinos to small values, where the annihilation is mostly dominated by A-boson exchange. For our choice of parameters, this constraint also forces small values for the total width of the heavy electron Γ_{E_1}. Note that from Eq. (17) one can see that the spin-dependent WIMP-nucleon scattering cross section scales with $s^4_{\theta_N}$, implying that the direct detection limits on the neutral sector mixing angle are expected to evolve rather slowly. Concretely, the projections for the planned LZ experiment [26] indicate an eventual sensitivity to scattering cross sections of the order of 7×10^{-12} pb for a WIMP mass of 350 GeV, a factor ~ 500 lower than the current LUX limits. This translates to an expected factor ~ 5 improvement on the limits on s_{θ_N}.

Additional constraints on the scenario come from searches for dark matter annihilation – induced gamma rays and in particular from the Fermi satellite searches for continuum gamma rays in Dwarf Spheroidal Galaxies (dSphs) [27] and for gamma ray lines from the Galactic center [28]. The former are depicted by the red band in Fig. 4. As for the latter, we find them to be subdominant throughout our parameter space even assuming a realistic halo profile for our Galaxy [30]. This might appear to be slightly counter-intuitive, since we are invoking here a mechanism that significantly boosts the diphoton signal. However, the threshold enhancement which is effective in the LHC environment is irrelevant in the case of indirect detection, due to the fact that the center of mass energy is not sufficient to produce the mediator A on-shell. This is in fact an attractive up-shot of the threshold enhancement mechanism invoked in our work, which relieves the tension that has been shown to exist [29] in models attempting to relate the 750 GeV diphoton excess with dark matter.

All in all, we see that the relic density, direct and indirect detection constraints are compatible with the small E_1 width values needed in order to reproduce the LHC diphoton excess, yielding a viable dark matter candidate under the form of a mostly singlet heavy neutrino. Moreover, at least for the parameter ranges depicted in Fig. 4, we expect that direct and indirect detection experiments will be able to probe a substantial part of the Planck compatible parameter space region within the next few years.

We now turn to the MSSM case and for our computations, we again employ MicrOMEGAs. We fix, as discussed previously, the lightest chargino mass at $m_{\chi_1^\pm} = \frac{1}{2}M_A = 375$ GeV and require a mass difference $m_{\chi_1^\pm} - m_{\chi_1^0} < M_W$ to ensure a small decay width for $\chi_1^\pm$. The lightest neutralino turns out to be an admixture of bino, higgsino and wino, with an under-abundant relic density $\Omega h^2 \sim 10^{-3}$–10^{-2}, on one hand as a result of the relatively strong couplings of mixed neutralino scenarios to A, and on the other hand of the neutralino mass being relatively close to $\frac{1}{2}M_A$, i.e. the so-called "funnel region". Thus, in this scenario, thermal relic neutralinos cannot account for more than $\mathcal{O}(10\%)$ of the total dark matter in the universe. Note that even for under-abundant dark matter components direct detection bounds do apply, upon appropriate rescaling of the limits. We find that in the relevant region of parameter space, the combination $\sigma_{\rm SI} \times \Omega_{\rm MSSM}/\Omega_{\rm Planck}$ lies below the LUX bounds.

In light of these findings, an interesting possibility would be to consider the option of gravitino dark matter, potentially upon embedding of our MSSM scenario in a (most likely general [31]) gauge-mediated supersymmetry breaking framework. Given that the gravitino abundance from neutralino decays will in general decrease as $m_{\tilde{G}}/m_{\chi_1^0}$ with respect to the – already under-abundant – neutralino relic density, the most likely scenario would in fact be thermal gravitino production; see e.g. Ref. [32] and references therein. Performing such an analysis goes beyond the scope of our study. In any case, gravitino dark matter with a general neutralino next-to-lightest superparticle has been extensively studied in Ref. [33]. Moreover, the under-abundance of these neutralinos should help relax the tension with Big Bang Nucleosynthesis constraints, see e.g. the recent discussion in Ref. [34].

5. Conclusions

In this paper, we have discussed the possibility of threshold enhancement of the branching ratio of pseudoscalar resonances into photon pairs. We focused on a 750 GeV pseudoscalar boson A, in light of the recent experimental hints of an excess in the diphoton spectrum at the LHC. If the loop mediating the $A \to \gamma\gamma$ decay was to contain new fermions at approximately half the mass of the resonance, i.e. $\frac{1}{2}M_A \sim 375$ GeV, then this decay could be significantly enhanced. The precise value of the enhancement factor was shown to depend on the width and the mass of these new fermions. Concretely, we found that fermion widths smaller than $\Gamma_f < 100$ keV, naturally occurring in 3-body decay processes, and masses m_f of a few MeV above 375 GeV could lead to an enhancement of the observed diphoton rate by two orders of magnitude. We then applied this idea to two concrete new physics scenarios where A could be produced via the gluon fusion mechanism: the minimal supersymmetric model and a two-Higgs doublet model augmented with one vector-like doublet and two singlets of leptons.

In the MSSM case, we found that a chargino with mass $m_{\chi_1^\pm} \approx 375$ GeV can provide the necessary enhancement factor to attain a diphoton cross section of the order of 4–8 fb as favored by ATLAS and CMS. We examined the chargino total decay width, where the chargino decays into a neutralino and a SM fermion pair through a sufficiently off-shell W boson. If this width lies below ~ 2 keV, a condition which can easily be satisfied for appropriate choices of the soft masses $M_{1,2}$ and the higgsino mass parameter μ, the necessary enhancement in the chargino loop in order to reproduce the observed diphoton excess, of the order of 400 at the cross section level, is obtained. To the best of our knowledge, this is the only explanation of the 750 GeV diphoton excess that has been proposed in the literature within the plain MSSM without any additional particle content (for an extension like the NMSSM with no additional particles, see for instance Ref. [35]).

We then discussed the threshold enhancement mechanism in the context of a basic 2HDM in which vector-like leptons are added to the spectrum. The lightest (mostly part of an isodoublet which also contains a neutral lepton) charged lepton E_1 is responsible for the threshold enhancement and its width is again given by the three-body decay $E_1 \to N_1 \bar{f} f'$ (with N_1 being mostly an isosinglet). In a large region of parameter space, it is found to lie in the desired region $\Gamma_{E_1} < 15$ keV, resulting in an enhancement factor of

[5] The spin-dependent scattering is found to be much weaker and will be ignored. Besides, we recall that pseudoscalar couplings of Dirac dark matter to the SM particles yield a negligible spin-independent scattering cross section, the latter being proportional to the momentum transfer which is extremely small compared to the mass M_A.

$F \sim 50$ of the cross section as required to explain the LHC diphoton data.

As a final exercise we studied whether the lightest neutral states of the new physics spectrum, the lightest neutralino in the MSSM and the lightest vector-like neutrino in our 2HDM variant, can play the role of dark matter in the Universe, where appropriate symmetries prevent the decay of these into SM particles. In our 2HDM scenario, we found that for the mass range from $m_{N_1} \approx 315$ to 350 GeV it is perfectly possible to satisfy the Planck constraints on the dark matter density in the Universe while being compatible with the LUX limits on the spin-independent scattering cross section off nuclei and the Fermi-LAT indirect searches for continuum γ-rays from dark matter annihilation in Dwarf Spheroidal Galaxies. Searches for gamma-ray lines were found to provide subleading constraints, since the threshold enhancement mechanism is not effective in dark matter annihilation at low velocities. In the MSSM case, the neutralino relic density turns out to be below the Planck value such that the direct detection constraints do not affect the region of interest.

Finally, let us note again that the scenarios exhibiting threshold enhanced diphoton signals are extremely contrived as they only occur in very narrow ranges of parameter space, i.e. of $m_f - \frac{1}{2}M_A$; therefore fine tuning at the 10 to 100 keV is required which may appear unnatural. Nevertheless, as it allows one to avoid complicated scenarios (with possibly a large multiplicity of new fermions) that are sometimes at the verge of being non-perturbative, Occam's razor leads us to believe that this "fine-tuned" scenario could constitute a plausible option. Of course, it might well be that this diphoton excess is simply a statistical fluctuation that will disappear with more data.

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References

[1] ATLAS Collaboration, ATLAS-CONF-2015-081;
CMS Collaboration, CMS-PAS EXO-15-004.
[2] See e.g., R. Franceschini, et al., arXiv:1512.04933 [hep-ph];
J. Ellis, S.A.R. Ellis, J. Quevillon, V. Sanz, T. You, arXiv:1512.05327 [hep-ph];
M.R. Buckley, arXiv:1601.04751 [hep-ph];
A. Djouadi, J. Ellis, R. Godbole, J. Quevillon, arXiv:1601.03696 [hep-ph];
F. Staub, et al., arXiv:1602.05581 [hep-ph].
[3] A. Angelescu, A. Djouadi, G. Moreau, arXiv:1512.04921 [hep-ph].
[4] A. Djouadi, Phys. Rep. 459 (2008) 1.
[5] See for instance M. Dhuria, G. Goswami, arXiv:1512.06782 [hep-ph];
F. Goertz, J.F. Kamenik, A. Katz, M. Nardecchia, arXiv:1512.08500 [hep-ph];
M. Fabbrichesi, A. Urbano, arXiv:1601.02447 [hep-ph];
E. Bertuzzo, P. Machado, M. Taoso, arXiv:1601.07508 [hep-ph];
A. Salvio, F. Staub, A. Strumia, A. Urbano, arXiv:1602.01460 [hep-ph].
[6] J. Ellis, M. Gaillard, D. Nanopoulos, Nucl. Phys. B 106 (1976) 292;
H. Georgi, S. Glashow, M. Machacek, D. Nanopoulos, Phys. Rev. Lett. 40 (1978) 692;
A.I. Vaĭnshteĭn, M. Voloshin, V. Zakharov, M. Shifman, Sov. J. Nucl. Phys. 30 (1979) 711;
J. Gunion, H. Haber, G. Kane, S. Dawson, The Higgs Hunter's Guide, Addison-Wesley, Reading, 1990;
A. Djouadi, Phys. Rep. 457 (2008) 1.
[7] See for instance, M. Drees, K.i. Hikasa, Phys. Rev. D 41 (1990) 1547.
[8] K. Melnikov, M. Spira, O. Yakovlev, Z. Phys. C 64 (1994) 401.
[9] V.S. Fadin, V. Khoze, Sov. J. Nucl. Phys. 48 (1988) 309, Yad. Fiz. 48 (1988) 487, JETP Lett. 46 (1987) 525, Pis'ma Zh. Eksp. Teor. Fiz. 46 (1987) 417.
[10] For a review on 2HDMs, see G. Branco, et al., Phys. Rep. 516 (2012) 1.
[11] See e.g. Y. Mambrini, G. Arcadi, A. Djouadi, Phys. Lett. B 755 (2016) 426, arXiv:1512.04913;
M. Backovic, A. Mariotti, D. Redigolo, arXiv:1512.04917 [hep-ph];
D. Barducci, A. Goudelis, S. Kulkarni, D. Sengupta, arXiv:1512.06842 [hep-ph];
P.B. Dev, D. Teresi, arXiv:1512.07243 [hep-ph];
P. Ko, T. Nomura, arXiv:1601.02490 [hep-ph].
[12] ATLAS Collaboration, J. High Energy Phys. 08 (2015) 105, arXiv:1503.05425;
CMS Collaboration, Phys. Lett. B 729 (2014) 149.
[13] M.x. Luo, et al., arXiv:1512.06670 [hep-ph];
C. Hand, et al., arXiv:1602.08100 [hep-ph];
Y. Kats, M. Strassler, arXiv:1602.08819 [hep-ph].
[14] M. Spira, A. Djouadi, D. Graudenz, P.M. Zerwas, Phys. Lett. B 318 (1993) 347;
M. Spira, A. Djouadi, D. Graudenz, P.M. Zerwas, Nucl. Phys. B 453 (1995) 17.
[15] ATLAS collaboration, J. High Energy Phys. 1508 (2015) 148;
CMS collaboration, arXiv:1506.03062.
[16] A. Sommerfeld, Ann. Phys. 11 (1931) 257.
[17] For a recent discussion, see for instance M. Beneke, et al., arXiv:1601.04718 [hep-ph].
[18] These numbers have been obtained in the context of the hMSSM discussed in
A. Djouadi, et al., J. High Energy Phys. 1506 (2015) 168;
A. Djouadi, et al., J. High Energy Phys. 1310 (2013) 028;
A. Djouadi, et al., Eur. Phys. J. C 73 (2013) 2650.
[19] The numerical analysis is performed using the program `HDECAY` and `FeynHiggs`: A. Djouadi, J. Kalinowski, M. Spira, Comput. Phys. Commun. 108 (1998) 56;
A. Djouadi, M. Muhlleitner, M. Spira, Acta Phys. Pol. B 38 (2007) 635;
S. Heinemeyer, W. Hollik, G. Weiglein, Comput. Phys. Commun. 124 (2000) 76.
[20] A. Djouadi, Z. Phys. C 63 (1994) 317;
G. Azuelos, A. Djouadi, Z. Phys. C 63 (1994) 327.
[21] G. Bélanger, F. Boudjema, A. Pukhov, A. Semenov, Comput. Phys. Commun. 92 (2015) 322.
[22] A. Alloul, N.D. Christensen, C. Degrande, C. Duhr, B. Fuks, Comput. Phys. Commun. 185 (2014) 2250.
[23] P.A.R. Ade, et al., Planck Collaboration, arXiv:1502.01589 [astro-ph.CO].
[24] G. Arcadi, Y. Mambrini, M.H.G. Tytgat, B. Zaldivar, J. High Energy Phys. 1403 (2014) 134.
[25] D.S. Akerib, et al., LUX Collaboration, arXiv:1512.03506 [astro-ph.CO].
[26] D.S. Akerib, et al., LZ Collaboration, arXiv:1509.02910 [physics.ins-det].
[27] M. Ackermann, et al., Fermi-LAT Collaboration, Phys. Rev. Lett. 115 (23) (2015) 231301.
[28] M. Ackermann, et al., Fermi-LAT Collaboration, Phys. Rev. D 91 (12) (2015) 122002.
[29] F. D'Eramo , J. de Vries, P. Panci, arXiv:1601.01571 [hep-ph].
[30] J.F. Navarro, C.S. Frenk, S.D.M. White, Astrophys. J. 490 (1997) 493;
A.W. Graham, D. Merritt, B. Moore, J. Diemand, B. Terzic, Astron. J. 132 (2006) 2685.
[31] P. Meade, N. Seiberg, D. Shih, Prog. Theor. Phys. Suppl. 177 (2009) 143.
[32] V.S. Rychkov, A. Strumia, Phys. Rev. D 75 (2007) 075011.
[33] L. Covi, J. Hasenkamp, S. Pokorski, J. Roberts, J. High Energy Phys. 0911 (2009) 003.
[34] For a recent discussion, see for instance A. Arbey, M. Battaglia, L. Covi, J. Hasenkamp, F. Mahmoudi, Phys. Rev. D 92 (11) (2015) 115008.
[35] U. Ellwanger, C. Hugonie, arXiv:1602.03344 [hep-ph];
F. Domingo, S. Heinemeyer, J.S. Kim, K. Rolbiecki, arXiv:1602.07691 [hep-ph];
M. Badziak, M. Olechowski, S. Pokorski, K. Sakurai, arXiv:1603.02203 [hep-ph].