Di-photon resonance around 750 GeV: shedding light on the theory underneath

Joydeep Chakrabortty, Arghya Choudhury, Pradipta Ghosh, Subhadeep Mondal, Tripurari Srivastava

Abstract

Both the ATLAS and CMS Collaborations have recently observed an excess in the di-photon invariant mass distribution in the vicinity of 750 GeV with a local significance of $\sim 3\sigma$. In this article we try to investigate this excess in the context of a minimal simplified framework assuming effective interactions of the hinted resonance with photons and gluons. We scrutinise the consistency of this observation with possible accompanying yet hitherto unseen signatures of this resonance. Subsequently, we try to probe the nature of new particles, e.g., spin, electric charge and number of colour, etc., that could remain instrumental to explain this excess through loop-mediation.

The recent observation by the LHC Collaborations [1,2,3,4], concerning an excess in the di-photon invariant mass distribution $m_{\gamma\gamma}$ near 750 GeV, has gained huge attention in the particle physics community. The ATLAS Group, using 3.2 fb$^{-1}$ of data with 13 TeV centre-of-mass energy (E$_{CM}$), has estimated a local (global) significance of 3.9$\sigma$ (2.0$\sigma$) for a mass of the resonance $M_X = 750$ GeV [3]. At the same time, the CMS Collaboration has noticed a local (global) significance of 2.8$\sigma$−2.9$\sigma$ (< 1.0$\sigma$) for $M_X = 760$ GeV [4] using 3.3 fb$^{-1}$ of data at E$_{CM}$ = 13 TeV. Combining with the run-I data (19.7 fb$^{-1}$ at E$_{CM}$ = 8 TeV), the CMS excess appears at $M_X = 750$ GeV [4] with a local (global) significance of 3.4$\sigma$ (1.6$\sigma$). The latter corresponds to a narrow width for the resonance, $\Gamma_X = 105$ MeV while interpretation with only 13 TeV data indicates $\Gamma_X = 10.6$ GeV. The ATLAS measurement, on the contrary, hints a large decay width $\Gamma_X = 45$ GeV [3].

This is the first surprise from LHC run-II with 13 TeV center-of-mass energy [3] which remains unexplained within the Standard Model (SM) framework. In other words, properties of the said resonance, as experimentally observed so far, e.g., excess in $\gamma\gamma$ only and nothing in $ZZ, Z\gamma$ or in di-jet ($jj$) channels, definitely demand physics beyond the SM (BSM). It is, thus, timely to explore the origin and associated consequences of this resonance although the possibility of losing this excess with more data-set can not be completely overlooked. A quest to accommodate this excess has already produced a handful of contemporary analyses [7,8,9,10,11,12,13,14,15,16,17,18,19,20,21,22,23,24] along with a few simultaneous [25,26,27,28,29] studies. Most of these analyses are proposed within the context of a specific theory framework, which often requires new decay modes (invisible for example) and thus, address other issues, for example the dark matter (see Refs. [8,9]). We, however, aim to investigate this excess with a simplified effective framework and will try to explore the nature of hitherto unseen particles which, while running in the loop, can appear instrumental to produce the observed di-photon excess.

With this idea we have used a generic Lagrangian which couples this new resonance $H_X$ with photons and gluons as shown by eq. (1). We have further assumed: (1) on-shell production of $H_X$ and (2) a scalar, i.e., spin-0, nature for $H_X$. The latter is one of the natural options to explain a resonance in di-photon channel, i.e., two identical massless spin-1 particles, as dictated by Landau-Yang theorem [30,31]. The effective minimal Lagrangian is written as:

$$L_{eff} = \kappa_s G_{\mu\nu}G^{\mu\nu}_0 H_X + \kappa_A B_{\mu\nu}B^{\mu\nu} H_X,$$  \hspace{1cm} (1)

where $G_{\mu\nu}$, $B_{\mu\nu}$ are the associated field strengths with “a” representing the relevant non-Abelian index. The effective $H_X$-$g$-$g$ interaction

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and $H_X$-$\gamma$-$\gamma$ vertices are parametrised as $\kappa_g$ and $\kappa_A$ which encapsulate the effect of new physics appearing in the loops. The latter is an absolute necessity since SM-like couplings between the SM-gauge bosons and $H_X$ appear inadequate \cite{32} to explain the observed sizable decay width $\Gamma_X$ \cite{11,14} and the production cross-section $\sigma(pp(gg) \to H_X \to \gamma\gamma) \sim O(10 \text{ fb})$ \cite{11,2,3,4}, consistent with the results of various other LHC searches. The observations from different LHC searches put strong constraints on the $\kappa_g$, $\kappa_A$ parameter space. The latter can be translated in terms of $H_X \to gg$, $H_X \to \gamma\gamma$ branching fractions ($Br$s) since they are $\approx 8\kappa_g^2$, $\kappa_A^2 \cos^2 \theta_W$, respectively. Moreover, the associated squared matrix elements are similar while the phase spaces are identical. The number $8'$ appears from the colour factor and $\theta_W$ is the Weinberg angle \cite{33}.

After the electroweak symmetry breaking, the second term of eq. \ref{eq1} generates effective interactions like $H_X\gamma\gamma$ and also $H_XZ\gamma$, $H_XZZ$, even with vanishing $\kappa_W$. Their strengths are $\approx \kappa_g \cos^2 \theta_W$, $\kappa_A \sin \theta_W \cos \theta_W$ and $\kappa_A \sin^2 \theta_W$, respectively. It is thus, important to note that a non-zero $Br(H_X \to \gamma\gamma)$ would also imply non-zero $Br(H_X \to Z\gamma, ZZ)$ values since all of them are connected to $\kappa_A$. Their relative magnitudes, however, remain different depending on the factor of $\sin \theta_W$ or $\cos \theta_W$. Measurements from the experimental collaborations for the said processes, using 13 TeV data, remain yet inadequate \cite{41}. Nevertheless, measured information for $H_X \to ZZ, Z\gamma$ and $H_X \to \gamma\gamma$ \cite{53} from the 8 TeV searches definitely constrain the range of $\kappa_A$, $\kappa_g$ parameters. For example, one obtains $\sigma(pp \to H_X \to ZZ) < 12 \text{ fb}$ \cite{39} and $\sigma(pp \to H_X \to Z\gamma) < 11 \text{ fb}$ \cite{40} from the similar searches performed by the ATLAS with 8 TeV data. The available parameter space is also constrained by the di-jets searches, given as $\sigma(pp \to jj) < 1.9 \text{ pb}$ \cite{41} such that the missing evidence of $pp \to H_X \to jj$ process at the 13 TeV appears consistent. needless to mention that the CMS collaboration has also made similar studies \cite{44,45,46,47}. Furthermore, if one wishes to account for a large $\Gamma_X$ by introducing new, e.g., invisible decays, one needs to incorporate the constraints from monojet searches accordingly \cite{48,49}.

In this article we have used the expected limits from 13 TeV LHC searches for $ZZ, Z\gamma, jj$ and $\gamma\gamma$ processes, derived using the 8 TeV results. We have used Madgraph v2.2.3 \cite{50,51} and observed that the production (via gluon fusion) cross-section with 13 TeV ECM is roughly five times of the same with 8 TeV ECM, i.e., $\sigma(pp \to H_X)_{13 \text{ TeV}}/\sigma(pp \to H_X)_{8 \text{ TeV}} \approx 5$, as also noted in Ref. \cite{15}. Further, we have also used the constraint from Ref. \cite{52} assuming that this resonance can also appear through photon fusion. In our numerical study we have used FeynRules 2.3 \cite{53} to implement eq. \ref{eq1} together with the SM Lagrangian. Subsequently, Madgraph v2.2.3 has been utilised to compute the production cross-section $\sigma(pp \to H_X)$ through gluon fusion and to calculate different partial decay widths of $H_X$. In this study we have utilised 3.2 fb$^{-1}$ of ATLAS data at 13 TeV to accommodate the observed resonance. In detail, we have used $\Delta N$, the discrepancy between the observed and expected number of events $= 13.6 \pm 3.69$. Further, for this purpose three 40 GeV bins are chosen for $690 \text{ GeV} \leq m_H^{\text{pp}} \leq 810 \text{ GeV}$ \cite{1} with an efficiency of 0.4.\cite{2}

In order to study the effect of BSM physics, we first show the variation of $\Gamma_X$ with changes in the new physics parameters, $\kappa_g$ (left), $\kappa_A$ (right), in Fig. \ref{fig1}. Here, we have varied $\kappa_g$, $\kappa_A$ in the span of $10^{-6} \text{ - } 1$. In these two plots the cyan coloured region represents the allowed $2\sigma$ range of $\Delta N$. The orange, golden and green coloured regions represent various zones in the $\Gamma_X - \kappa_g (\kappa_A)$ planes that are excluded from the 8 TeV LHC measurements of $H_X \to ZZ, Z\gamma, jj$ processes. The yellow coloured region remains excluded from the measurement of $H_X \to \gamma\gamma$ \cite{33} process at the ATLAS with 8 TeV centre-of-mass energy. Lack of precision measurements for the latter, assuming $\sigma(pp \to H_X)_{13 \text{ TeV}}/\sigma(pp \to H_X)_{8 \text{ TeV}} \approx 5$, predicts a $2\sigma$ upper bound \cite{6} on $\sigma(pp \to H_X \to \gamma\gamma)_{13 \text{ TeV}}$ inconsistent with the one observed with 13 TeV. We will discuss this later in detail. Finally, the grey coloured region remains excluded from the photon fusion process, i.e., $\gamma\gamma \to H_X \to \gamma\gamma$, \cite{52} which predicts a maximum for $Br(H_X \to \gamma\gamma)$, independent of $Br(H_X \to gg)$. The region excluded by the photon fusion process is estimated by assuming that $H_X$ has only two decay modes $gg, \gamma\gamma$, i.e., $Br(H_X \to gg) + Br(H_X \to \gamma\gamma) = 1$. The observed limits on the $Br$s are subsequently translated in terms of $\kappa_g$ and $\kappa_A$.

It is evident from Fig. \ref{fig1} that expecting $\Gamma_X$ as large as 45 GeV or more is perfectly consistent with the observed limits on $ZZ, \gamma\gamma, Z\gamma$ searches at the 8 TeV LHC. However, it is the di-jet searches which rules out the region of parameter space with $\Gamma_X > 3 \text{ GeV}$ (right plot), corresponding to $\kappa_g \gtrsim O(0.001)$ (left plot). The observed behaviour is well expected as $\Gamma_X \sim gg$ and thus, $\Gamma_X$ grows rapidly with $\kappa_g$, compared to that with $\kappa_A$, i.e., $\Gamma_X \sim \gamma\gamma$ since the latter is suppressed by a factor of $\cos^4 \theta_W/8$. For $\kappa_g$ (estimated from $Br(H_X \to \gamma\gamma)$), the most stringent bound is coming from the photon fusion process which is represented by the grey coloured region. For the photon fusion process, $Br(H_X \to \gamma\gamma) \propto 1/\sqrt{\Gamma_X}$ \cite{52} and thus, smaller upper bound on $Br(H_X \to \gamma\gamma)$ and hence, on $\kappa_g$ is expected for larger $\Gamma_X$. This is evident from the right plot of Fig. \ref{fig1}. It is important to note that the photon fusion process can also provide an indirect bound on $Br(H_X \to gg)$, i.e., on $\kappa_g$, assuming $Br(H_X \to \gamma\gamma) + Br(H_X \to gg) = 1$. It is also apparent that

\cite{7}The ATLAS and CMS collaborations have recently reported $ZZ$ \cite{54,55} and $Z\gamma$ \cite{56,57} search results with early 13 TeV data.

\cite{8}For di-jet searches, early 13 TeV results are also available \cite{42,43}.
the photon fusion process discards $\Gamma_X \gtrsim 0.3$ GeV which is 10 times smaller than the one predicted from the di-jet search limit. Hence, given the observed large $\Gamma_X$ from the ATLAS, one needs almost the equal amount of $\Gamma_X$ from the hitherto unseen decay modes of this resonance, e.g., invisible decays. Here, we use $\Gamma_X = \Gamma_{H_X \to \gamma\gamma} + \Gamma_{H_X \to ZZ} + \Gamma_{H_X \to jj} + \Gamma_{H_X \to \ell\ell} + \cdots$, as expected from eq. \ref{eq:GammaX}, to estimate $\Br(H_X \to \gamma\gamma)$ for the photon fusion process. It is now clear that in the chosen setup, no realistic values of $\kappa_\Delta$, $\kappa_\tau$ parameters can account for a total $\Gamma_X \gtrsim 0.3$ GeV. Thus, the presence of a huge additional decay width is essential for the studied construction which will be tightly constrained from the dark matter and monojet searches.

The discussion presented so far concerning the photon fusion process has one caveat related to the estimation of $\Br(H_X \to gg)$. So far, we have used eq. \ref{eq:GammaX} to estimate $\Gamma_X$, however, while evaluating the effect of photon fusion process on $\Br(H_X \to gg)$, i.e., on $\kappa_\tau$ (left plot of Fig. \ref{fig:GammaX}), we have used $\Br(H_X \to \gamma\gamma) + \Br(H_X \to gg) = 1$ which is apparently contradictory. At this point one must note that in the given construction the quantities $\Br_s(H_X \to ZZ, ZZ \to H_X)\text{,}$ as already explained, are suppressed compared to $\Br(H_X \to \gamma\gamma)$. Moreover, so far we have no information available for processes like $YZ, ZZ \to H_X$. Thus, the assumption $\Br(H_X \to gg) = 1 - \Br(H_X \to \gamma\gamma)$ remains useful for estimating the scale of $\Br(H_X \to gg)$. Using the all available branching fractions instead would yield weaker upper bounds on $\Br(H_X \to gg)$, i.e., on $\kappa_\tau$.

It is evident from Fig. \ref{fig:GammaX} that $\Gamma_X \gtrsim 0.3$ GeV appears excluded from the relevant existing LHC limits and from the constraint of photon fusion process. This observation demands the existence of huge additional decay width to reach the target of 45 GeV. If we call this additional width as $\Gamma_X^{\text{add}}$, without specifying the origin, then one can write $\Gamma_X \equiv \Gamma_X^{\text{add}} \equiv \Gamma_X^{\text{add}} + \Gamma_{H_X \to \gamma\gamma} + \Gamma_{H_X \to ZZ} + \Gamma_{H_X \to jj} + \Gamma_{\ell\ell}$. This approach will modify all the associated branching ratios as will be explored subsequently by choosing three different values of the total decay widths: (1) 1 GeV (small width), (2) 10 GeV (moderate width) and (3) 45 GeV (large width).

The subsequent effects of the aforesaid construction are explored in Fig. \ref{fig:GammaX} where we have investigated the impact of diverse LHC and photon fusion constraints in the $\Br(H_X \to gg)$ - $\Br(H_X \to \gamma\gamma)$ plane. These two $\Br_s$ are expected to show some kind of correlation between them since the observed excess appears through $gg \to H_X$ process followed by $H_X \to \gamma\gamma$ decay. It is also possible to observe a similar correlation in the $\kappa_\tau, \kappa_\Delta$ plane since $\Br(H_X \to gg)$, $\Br(H_X \to \gamma\gamma) \propto \kappa_\tau^2, \kappa_\Delta^2$, respectively. In Fig. \ref{fig:GammaX} the black coloured line represents the best-fit value corresponding to $\Delta N = 13.6$ while the cyan and blue coloured bands represent the $1\sigma$ ($9.91 \lesssim \Delta N \lesssim 17.29$) and $2\sigma$ ($6.22 \lesssim \Delta N \lesssim 20.98$) allowed regions in the concerned planes, respectively. The orange, green, red and yellow coloured regions, similar to Fig. \ref{fig:GammaX} represent various zones in the concerned plane that are excluded from 8 TeV LHC limits on $H_X \to ZZ, ZY, jj$ and $\gamma\gamma$ processes. In the case of $H_X \to \gamma\gamma$ process, assuming $\sigma(pp \to H_X)\|_{13\text{ TeV}}/\sigma(pp \to H_X)\|_{\text{ATLAS}} = 5$, one would expect a 2$\sigma$ upper bound \cite{sphalerone} on $\sigma(pp \to H_X \to \gamma\gamma)\|_{13\text{ TeV}}$ as 10 fb using the ATLAS data. This is in tension with the 13 TeV ATLAS observation \cite{atlasphoton} and rules out higher values of the observed $\sigma(pp \to H_X \to \gamma\gamma)$, starting from the central one. A similar analysis using the CMS data \cite{cms} excludes the higher values of the observed $\sigma(pp \to H_X \to \gamma\gamma)\|_{13\text{ TeV}}$ beyond 1$\sigma$.

Lastly, the photon fusion process at the LHC, which predicts a maximum for $\Br(H_X \to \gamma\gamma)$ independent of $\Br(H_X \to gg)$, rules out the grey coloured region in the $\Br(H_X \to gg)$ - $\Br(H_X \to \gamma\gamma)$ plane. It is interesting to note that the constraint for the photon fusion was derived with the assumption of $\Br(H_X \to gg) + \Br(H_X \to \gamma\gamma) = 1$ which discards a region where $\Br(H_X \to gg) + \Br(H_X \to \gamma\gamma) > 1$. For the three chosen values of $\Gamma_X$, the maximum $\Br(H_X \to \gamma\gamma)$ is estimated \cite{photonfusion} as $\sim 0.42, 0.13, 0.06$, respectively and thus, the regions with $\Br(H_X \to gg) > 0.58$ (left plot of Fig. \ref{fig:GammaX}), $\Br(H_X \to gg) > 0.87$ (middle plot of Fig. \ref{fig:GammaX}), $\Br(H_X \to gg) > 0.94$ (right plot of Fig. \ref{fig:GammaX}) remain ruled out. The upper limits of $\Br(H_X \to gg)$, as depicted in Fig. \ref{fig:GammaX} are purely illustrative. This is because, following our earlier discussion, $\Br(H_X \to gg) = 1 - \Br(H_X \to \gamma\gamma)$ estimated in a regime when $\Gamma_{H_X}^{\text{add}} \gg \Gamma_X \equiv \Gamma_X$ appears simply illustrative. For the rest of the processes the primary productions are driven by the gluon fusion process. The latter gives a high value for $\Br(H_X \to gg)$ with increasing $\kappa_\tau$ and as a consequence remains excluded from the di-jet search limits, especially for moderate to large $\Gamma_X$. For example, for the choice of $\Gamma_X = 10$ GeV one gets $\Br(H_X \to gg)\max \sim 0.40$ (middle plot of Fig. \ref{fig:GammaX}) while for the choice of $\Gamma_X = 45$ GeV one ends up with $\Br(H_X \to gg)\max \sim 0.20$ (right plot of Fig. \ref{fig:GammaX}). In the case of small decay width (left plot of Fig. \ref{fig:GammaX}) constraint from the di-jet searches remains ineffectual.

It is evident from eq. \ref{eq:GammaX} that $\Br(H_X \to ZZ)$, $\Br(H_X \to ZY)$ are suppressed compared to $\Br(H_X \to \gamma\gamma)$ by factors of $\tan^2 \theta_W$ and $\tan^2 \theta_W$ (numerically $\sim 0.09$ and 0.3), respectively which is also apparent from Fig. \ref{fig:GammaX}. Thus, unless one introduces interaction like $\kappa_\rho W^\mu_{\rho\nu} W^{\nu\nu}_\rho$ ($W^\nu_{\rho\nu}$ as the $SU(2)$ field strengths) these modes remain sub-leading. One can, nevertheless, compensate these deficits with a larger $\Br(H_X \to gg)$, assuming $gg \to H_X$ to be the leading production channel. These behaviours are reflected in Fig. \ref{fig:GammaX} where the regions excluded from ZZ and ZY searches appear with lateral shifts towards larger $\Br(H_X \to gg)$ values compared to $\Br(H_X \to \gamma\gamma)$ values, required to reproduce the observed excess. Larger $\Br(H_X \to gg)$ and hence larger $\Gamma_X$ appear naturally for higher $\kappa_\tau$ values which are in tension with the di-jet searches. Increasing $\kappa_\Delta$ receives constraint from the photon fusion process. The ZZ and ZY constraints, as already mentioned, require large values for both of the $\Br(H_X \to gg)$ and $\Br(H_X \to \gamma\gamma)$. The former faces tension from the di-jet search limits (moderate and large $\Gamma_X$ scenarios) while the latter, if not excluded by the photon fusion constraint, might give larger $\Delta N$ than actually observed. Hence, the parameter space ruled out by these constraints do not affect the signal region compatible for explaining the observed excess. The key feature of Fig. \ref{fig:GammaX} is the prediction...
Now we are ready to discuss the presence of other BSM particles that are essential to explain this excess through higher order processes. Information about these states are encapsulated in the Higgs potential up to the Planck scale [55, 56, 57, 58]. This has additional phenomenological advantages, e.g., stability of the SM-Higgs. From eq. (2) it appears that a larger $\Delta M$ is useful to produce a bigger $\Gamma_{H\gamma\gamma}$ in a test of perturbative suppression and at the same time, must possess sizable couplings with $H_X$ to reproduce the detected excess. Concerning the leading production, i.e., $gg \to H_X$, the possible candidate(s) is(are) either new coloured scalar(s) $\Phi$ or additional coloured fermion(s) $F$, possibly vector-like. These new particles must simultaneously couple to gluons as well as to $H_X$ and, are possibly embedded in a representation of some larger symmetry group. If these new scalars/fermions are also responsible for producing an enhanced $Br(H_X \to gg)$, they must carry electrical charges to get coupled to a photon. However, the non-minimal possibility is to consider another set of uncoloured but electrically charged fermion(s), scalar(s) or gauge boson(s) (appears in theories with extended non-Abelian gauge sector). Note that contributions from new chiral fermions produce a destructive effect compared to the bosonic contributions and thus, often are not compatible with the observed excess. On the other hand, vector-like fermions remain a viable alternative. The presence of an extended scalar sector has additional phenomenological advantages, e.g., stability of the SM-Higgs potential up to the Planck scale [55, 56, 57, 58]. This argument also holds true for new gauge boson(s). We, however, do not consider them in this article since they are hinted to be rather heavy $\gtrsim 2.5$ TeV [59, 60]. In a nutshell, we conclude that to accommodate the observed di-photon excess one needs sizable couplings between $H_X$ and the new particles, for which coloured and/or electrically charged scalars or fermions remain the realistic options. Moreover, in the presence of the said new states, an enhanced $Br(H_X \to \gamma\gamma)$ is more anticipated compared to an enlarged $Br(H_X \to gg)$ as for the latter experimental evidences are still missing.

The fact that the same product is $\sim O(10^{-11})$ [54] for a 750 GeV state with SM-like properties justifies BSM nature of this excess. In the presence of a new BSM scalar $\Phi$, with mass $M_\Phi$, effective charge $Q_\Phi$ and number of colour $N_\Phi$, the $Br(H_X \to \gamma\gamma)$ can be written as [32, 61],

$$Br(H_X \to \gamma\gamma) = \frac{\alpha^2 M_\Phi^3}{16\pi^3} \frac{\sin^2 \theta_{\Phi H_\gamma}}{M_\Phi^2} \frac{Q_\Phi^4}{N_\Phi^2} \frac{1}{\Gamma_{H\gamma\gamma}}. \quad (2)$$

Here, $\alpha_{em}$ is the electromagnetic coupling constant, $\theta_{\Phi H_\gamma}$ represents the coupling between $\Phi$ and $H_X$ and the detail of $A_0(x_\Phi)$ function, where $x_\Phi = 4M_\Phi^2/M^2$ is given in Ref. [32]. Keeping in mind the issue of perturbativity we choose $-\sqrt{4\pi} \leq \theta_{\Phi H_\gamma} \leq \sqrt{4\pi}$, in our numerical analyses. The quantity $\theta_{\Phi H_\gamma}$ parametrises the information about the vacuum expectation value (VEV) of $H_X$ and the amount of possible mixing between $H_X$ and the SM-Higgs. From eq. (2) it appears that a larger $\theta_{\Phi H_\gamma}$ is useful to produce a bigger $Br(H_X \to \gamma\gamma)$. In reality, however, such scenarios are unrealistic as they correspond to either experimentally challenging large mixing within $H_X$ and the SM-Higgs or a large VEV for $H_X$ inconsistent with the electroweak precision tests [63]. It is apparent from eq. (2) that depending on the values of $M_\Phi$, $Q_\Phi$ and $N_\Phi$, the quantity $Br(H_X \to \gamma\gamma)$ can receive sizable enhancement. An enlargement is also possible if the future LHC observation confirms a smaller $\Gamma_X$. In our numerical analyses we choose $400$ GeV $\ll M_\Phi$ $\lesssim 1000$ GeV, consistent with the existing collider bounds on such exotic particles [62, 63, 64]. A sample variation of $Br(H_X \to \gamma\gamma)$ in the $M_\Phi - \theta_{\Phi H_\gamma}$ plane for a colour singlet ($N_\Phi = 1$), triply charged ($Q_\Phi = 3$), scalar with different $\Gamma_X$, 1 GeV (left) and 45 GeV (right) is shown in Fig. [3]. It is evident from Fig. [3] that an experimentally viable light, i.e., $M_\Phi = 400$ GeV, colour singlet $\Phi$ with $Q_\Phi = 3$ can produce at most a $Br(H_X \to \gamma\gamma) \sim O(10^{-8})$ (left plot) when $\Gamma_X$ is small, i.e., 1 GeV. Choosing $\Gamma_X = 45$ GeV instead one faces a reduction by a factor of 45 (right plot). From eq. (2) we see that $Br(H_X \to \gamma\gamma) \propto Q_\Phi^4$. Thus, even for an exotic colour singlet $\Phi$ with $Q_\Phi = 3$, one would expect a maximum $Br(H_X \to \gamma\gamma) \sim O(10^{-8})$ keeping $\Gamma_X$, $M_\Phi = 1$ GeV, 400 GeV. Now, from our previous discussion in the context of Fig. [2] we have estimated $Br(\gamma\gamma \to gg)$ as $\sim O(10^{-8})$ and $\sim O(10^{-5})$ for the choice of $\Gamma_X = 1$ and 45 GeV, respectively from the best-fit value of $\Delta N$. Hence, the maximum $Br(H_X \to \gamma\gamma)$, extracted from Fig. [3] using eq. (2) for a 400 GeV $\Phi$ with $N_\Phi$, $Q_\Phi = 1$, 10, would give an unrealistic $Br(H_X \to gg) \sim 400(180)$ for $\Gamma_X = 1(45)$ GeV scenarios. One may try to consider a similar but coloured (say...
Figure 3: Plots showing the variation of $Br(\gamma\gamma) = Br(H_X \rightarrow \gamma\gamma) \times 10^6$ in the $M_{\Phi} - |g_{\Phi\PhiHX}|$ plane for $N_\Phi^c = 3$ with $Q_\Phi = 1$, $\Gamma_X = 1$ GeV (left) and 45 GeV (right). The chosen ranges for $M_\Phi$ and $g_{\Phi\PhiHX}$ are explained in the text. The multiplicative factor $n = 9(11)$ for $\Gamma_X = 1(45)$ GeV.

Figure 4: Variation of $Br(\gamma\gamma) = Br(H_X \rightarrow \gamma\gamma) \times 10^6$ in the $N_{\Phi}^c - Q_\Phi$ plane for $M_\Phi = 600$ GeV, $g_{\Phi\PhiHX} = 1$, $\Gamma_X = 1$ GeV(left) and 45 GeV(right). Here, $n = 5(7)$ for $\Gamma_X = 1(45)$ GeV.

Figure 5: Plots showing the variation of $Br(\gamma\gamma) = Br(H_X \rightarrow \gamma\gamma) \times 10^6$ in the $M_F - |g_{FFH_X}|$ plane for $N_F^c = 1$, $Q_F = 1$ with $\Gamma_X = 1$ GeV (left) and 45 GeV (right). The chosen ranges for $M_F$ and $g_{FFH_X}$ are explained in the text. Here, $n = 3(5)$ for the left(right) plot.

$N_\Phi^c = 3$) $\Phi$ which predicts a maximum $Br(H_X \rightarrow \gamma\gamma) \sim O(10^{-3})$ for $\Gamma_X = 1$ GeV. However, one still needs an unrealistic $Br(H_X \rightarrow gg) \sim 50$ in this scenario. Moreover, for a $\Phi$ with non-zero colours one must carefully investigate the $H_X \rightarrow jj$ constraint, even for a realistic $Br(H_X \rightarrow gg)$, especially for moderate to large $\Gamma_X$.

From the last discussion it appears that the use of new BSM scalar is not adequate to explain the observed excess. In order to explore this further we have plotted the variation of $Br(H_X \rightarrow \gamma\gamma)$ in the $N_{\Phi}^c - Q_\Phi$ plane in Fig. 3 with $M_\Phi = 600$ GeV, 1 for the choice of $\Gamma_X = 1$ GeV (left) and 45 GeV (right). Here, we vary both $N_{\Phi}^c$, $Q_\Phi$ in the range of 1 : 20 and the chosen values of $M_\Phi$, $g_{\Phi\PhiHX}$ are purely illustrative. It is apparent from both of these plots that to satisfy $Br(\gamma\gamma \times gg)$, consistent with the observation of Fig. 2, one should have an unrealistic $Br(H_X \rightarrow gg) \sim 10(5)$ for $\Gamma_X = 1(45)$ GeV. Adopting smaller $M_\Phi$ (say 400 GeV) simultaneously with a larger $g_{\Phi\PhiHX}$ (say $\pm 3$) one can reach a maximum $Br(H_X \rightarrow \gamma\gamma) \sim 0.012$ and $\sim 0.00025$ for $\Gamma_X = 1$ and 45 GeV, respectively considering $N_\Phi^c \geq 14$, $Q_\Phi \geq 16$. Here, we have used eq. (2) and information from Fig. 2. So, apparently these exotic scenarios can give a realistic $Br(H_X \rightarrow gg) \sim O(0.1)$, consistent with the di-jet searches (see Fig. 2). However, this moderate $Br(H_X \rightarrow gg)$ value may get excluded from the future LHC searches with expected higher sensitivity. Moreover, one must carefully re-evaluate the maximum value for $g_{\Phi\PhiHX}$ in a consistent theory framework. It is now evident from the last discussion that the presence of BSM $\Phi$s, instrumental to reproduce the observed excess, requires really high electric and colour charges. Particles with such high colour charges are expected to be produced amply at the LHC, unless very massive and hence, rather stringent constraints are expected on their existence. We thus, leave our discussion about the BSM scalars without further detail. We note in passing that $Q_\Phi$ value as high as 20 can be interpreted as an effective electric charge, keeping $N_\Phi^c$ fixed. For example data-set with $Q_\Phi = 20$ for a fixed $N_\Phi^c$, using eq. (2), can be thought of as a coloured/uncoloured multiplet with members of almost the same masses and having electric charges from $\pm 1$ to $\pm 10$.

Let us now investigate a similar scenario in the presence of new BSM vector-like fermion, $F$. For a fermion with mass $M_F$, electric charge $Q_F$ (in the units of $|e|$), number of colours $N^c_F$, the quantity $Br(\gamma\gamma \rightarrow \gamma\gamma)$ is expressed as $Br(\gamma\gamma \rightarrow \gamma\gamma) = \frac{g^2 M_X^2}{16\pi^2 f^2} \frac{2 g_{FFH_X} M_F}{M_F^2} N^c_F Q_F^2 A_F(x_f)^2 \sim O(10^{-3})$.

The sample variation of $Br(H_X \rightarrow \gamma\gamma)$ in the $M_F - g_{FFH_X}$ plane for a colour singlet doubly charged ($Q_F = 2$) fermion is shown in Fig. 5 for $\Gamma_X = 1$ (left) and 45 GeV (right). It is easy to see from these plots that the presence of BSM fermions is more efficient to raise $Br(H_X \rightarrow \gamma\gamma)$ compared to the BSM scalars. For example a colour singlet doubly charged fermion can produce $Br(H_X \rightarrow \gamma\gamma)$ as high as 0.007 and $\sim 10^{-4}$ for $\Gamma_X = 1$ and 45 GeV, respectively. These numbers are orders of magnitude larger compared to the same from Fig. 2 and, as stated before, can only be achieved for a $\Phi$ with very high $Q_\Phi$ and $N_\Phi^c$. These enhanced $Br(H_X \rightarrow \gamma\gamma)$ values are also useful to estimate realistic values of $Br(H_X \rightarrow gg)$, using the information from Fig. 2. As an example, from Fig. 5 with the maximum of $Br(H_X \rightarrow \gamma\gamma)$, one can estimate $Br(H_X \rightarrow gg) \sim 0.14(0.063)$ for $\Gamma_X = 1(45)$ GeV using the derived bound on $Br(\gamma\gamma \times gg)$. Clearly, one can easily reproduce the observed excess, especially for smaller $\Gamma_X$, without any difficulty. However, for larger $\Gamma_X$, depending on its value, some of the $Br(H_X \rightarrow gg)$ values
branching ratio \([66]\) is given as:
\[
Br = \frac{\alpha_f^2 M_X^4}{512 \pi^2 \Gamma_X} \left| \frac{g_{FFH}}{M_F} A_f(x_F) \right|^2.
\] (4)

Here, \(\alpha_f\) is the strong coupling constant. In our numerical analysis we have multiplied \(Br(H_X \rightarrow gg)\), as shown in eq. (4), by a factor of 1.5, relevant for the higher order effects of strong interactions. Using eqs. (3) and (4) simultaneously we have studied a sample variation of \(\Gamma_Xg_{gF}(\gamma \gamma x)\) in the \(M_F-g_{FFH}g\) plane as shown by Fig. 4. For this figure \(M_F, g_{FFH}\) are varied as of Fig. 2. We work with \(N_F = 3, Q_F = 3\). The observed behaviours of \(\Gamma_Xg_{gF}(\gamma \gamma x)\) with different parameters, i.e., \(M_F, g_{FFH}\), and \(\Gamma_X\) are expected from eqs. (3) and (4). For example, both \(Br(H_X \rightarrow \gamma \gamma)\) and \(Br(H_X \rightarrow gg)\) are \(\propto \Gamma_X\) and thus, shrinking of the allowed parameter space, compatible with the observed excess, for larger \(\Gamma_X\), 45 GeV. (right plot of Fig. 7) is anticipated. At the same time, these two branching ratios are \(\propto g_{gF}(H, \gamma \gamma x, M_F^2)\) (see eqs. (3), (4)). Hence, apparent lowering of \(Br(H_X \rightarrow gg)\) for larger \(M_F\) values must be compensated with larger \(g_{FFH}\) values in order to remain compatible with the excess. This feature is depicted in Fig. 7, notably for the left one. The most useful aspect of Fig. 7 is connected with the estimation of future detection possibility for the production \(gg \rightarrow H_X \rightarrow FF\). Assuming that the future measurements indicate a narrow width for this excess, say 1 GeV, then the room for measuring \(\sigma(gg \rightarrow H_X \rightarrow FF)\) is less promising for two reasons: (1). The expected enhancement in the production for low \(M_F\) region is ameliorated with a relatively small \(g_{FFH}\). (2). In the high \(g_{FFH}\) regime, the same logic remains applicable through heavier \(M_F\). These two features are visible from the left plot of Fig. 7. On the contrary, a more stringent limit, i.e., \(Br^2(\gamma \gamma x) \sim O(10^{-3})\), for larger \(\Gamma_X = 45\) GeV prefers smaller \(M_F\) and larger \(g_{FFH}\) (see right plot of Fig. 7). Both of these would appear useful to enhance \(\sigma(gg \rightarrow H_X \rightarrow FF)\).

**Conclusions:** To summarise, the LHC run-II has already observed an excess in the di-photon invariant mass distribution near 750 GeV. This excess, as argued in this article, definitely requires BSM physics. In this article we tried to explore this excess, assuming a spin-0 nature, using an simplified effective Lagrangian, sensitive to new physics effects. The chosen framework helped us to estimate a lower bound of \(\Gamma_X\), consistent with the different LHC constraints and photon fusion process, for changes in the new physics parameters, \(\kappa_F, \kappa_A\). We have also explored the possible correlation between \(Br(H_X \rightarrow \gamma \gamma)\) and \(Br(H_X \rightarrow gg)\) in the light of the observed excess.
and diverse possible constraints. This correlation provides a model-independent but $\Gamma_{\gamma\gamma}$-dependent bound on $\text{Br}(H_K \to \gamma\gamma) \times \text{Br}(H_K \to gg)$. Subsequently, we have utilised this correlation to scrutinise the effect of other BSM scalars, fermions with various electric charge, number of colour which simultaneously couple to $H_K$ and $gg$, $\gamma\gamma$ and might appear instrumental to reproduce this excess through higher order processes. Our analyses show that to accommodate the observed excess, the presence of additional BSM fermions are preferred compared to the scalars. Moreover, detecting these new fermions in the future is more anticipated for a large width of the observed excess. In conclusion, given this diphoton excess survives with more data-set, this cannot be an isolated surprise. Rather, this must be the pioneering evidence of a BSM mass spectrum while other heavier members are awaiting to be detected.

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