Interpreting the 750 GeV digamma excess: a review

ALESSANDRO STRUMIA
CERN, INFN and Dipartimento di Fisica, Università di Pisa

We summarise the main experimental, phenomenological and theoretical issues related to the 750 GeV digamma excess.

The first LHC data about $pp$ collisions at $\sqrt{s} = 13$ TeV agree with the Standard Model (SM), except for a hint of an excess in $pp \rightarrow \gamma\gamma$ peaked at invariant mass around 750 GeV [1]. We denote the new resonance with the symbol, $\mathcal{F}$, used in archaic Greek as the digamma letter and later as the number 6 $\approx M_\mathcal{F}/M_h$, but disappeared twice. New data will tell if the $\mathcal{F}$ resonance disappears or is confirmed. In the meantime, the $\mathcal{F}$ excess attracted significant theoretical interest [2–370]. Indeed, unlike many other anomalies that disappeared, the $\gamma\gamma$ excess cannot be caused by a systematic issue, neither experimental nor theoretical. Theoretically, the SM background is dominated by tree-level $q\bar{q} \rightarrow \gamma\gamma$ scatterings, which cannot make a $\gamma\gamma$ resonance. Experimentally, one just needs to identify two photons and measure their energy and direction. The $\gamma\gamma$ excess is either the biggest statistical fluctuation since decades, or the main discovery.

1 Data

During the Moriond 2016 conference CMS presented new data taken without the magnetic field; ATLAS presented a new analysis with looser photon selection cuts (called ‘spin 2’ analysis to distinguish it from the earlier ‘spin 0’ analysis); furthermore both collaborations recalibrated photon energies in a way optimised around 750 GeV rather than around $M_h = 125$ GeV. As a result, the statistical significance of the $\gamma\gamma$ excess increased slightly, both in CMS and in ATLAS.

See [302, 346, 365] for attempts of finding a Standard Model interpretation.
Fig. 1a shows the $\gamma\gamma$ spectra: we consider the ‘spin 0’ ATLAS analysis and the sum of CMS photon categories. Both ATLAS and CMS find the most statistically significant $\gamma\gamma$ excess around 750 GeV. Their consistency can be seen from the peak in fig. 1b where we summed ATLAS and CMS event counts. The width of the resonance ranges between 0 and 100 GeV, and can be larger (‘broad’) or smaller (‘narrow’) than the experimental resolution of about 6−10 GeV. The best-fit width is $\Gamma \sim 45 \text{ GeV} \sim 0.06 M_f$. The total rates in the two cases, narrow and broad, are:

$$
\begin{array}{c|cc|cc}
 & \sqrt{s} = 8 \text{ TeV} & & \sqrt{s} = 13 \text{ TeV} \\
 & \text{narrow} & \text{broad} & \text{narrow} & \text{broad} \\
\hline
\text{CMS} & 0.63 \pm 0.31 \text{ fb} & 0.99 \pm 1.05 \text{ fb} & 4.8 \pm 2.1 \text{ fb} & 7.7 \pm 4.8 \text{ fb} \\
\text{ATLAS} & 0.21 \pm 0.22 \text{ fb} & 0.88 \pm 0.46 \text{ fb} & 5.5 \pm 1.5 \text{ fb} & 7.6 \pm 1.9 \text{ fb}
\end{array}
$$

ATLAS and CMS do not perform a combined analysis. Naïve combinations of Higgs data gave results close to the official joint combination, so fig. 2 shows the naïve global fit for $\sigma(pp \rightarrow F \rightarrow \gamma\gamma)$ at $\sqrt{s} = 8, 13 \text{ TeV}$. The local excess is about $4\sigma$. The ‘Look Elsewhere Effect’ (LEE) reduces the global statistical significance by about $1\sigma$, assuming that an excess could have materialised in $\sim 10^2$ other places within the same data-set. The trial factor can be increased to $10^3$ by considering other similar data-sets, or to $10^4$ by considering that this search has been repeated about 10 times in the past decades. We don’t need to address such details: new data will decide if $F$ will reach the SM scalar $h$ in the Particle Data Group or if $F$ will instead reach $N$-rays in the cemetery of anomalies.

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Footnote: We leave to the intelligence of the reader to evaluate the possible statistical meaning of this unusual procedure.
2 Widths

The cross section for single production of a boson $F$ with spin $J$ can be written in the narrow-width approximation in terms of its decay widths into partons $\varphi$, $\Gamma_\varphi = \Gamma(F \to \varphi)$, as

$$\sigma(pp \to F) = \frac{2J+1}{s} \sum_\varphi C_\varphi \frac{\Gamma_\varphi}{M_F}.$$  \hspace{1cm} (2)

A resonance with spin $J = 1$ is excluded because it cannot decay into $\gamma\gamma$ (Lee-Yang theorem). The luminosity factors $C_\varphi$ of the main partons are \cite{10}

| $\sqrt{s}$   | $C_{b\bar{b}}$ | $C_{c\bar{c}}$ | $C_{s\bar{s}}$ | $C_{d\bar{d}}$ | $C_{u\bar{u}}$ | $C_{gg}$ |
|--------------|----------------|----------------|----------------|----------------|----------------|---------|
| 8 TeV        | 1.07           | 2.7            | 7.2            | 89             | 158            | 174     |
| 13 TeV       | 15.3           | 36             | 83             | 627            | 1054           | 2137    |

QCD corrections enhance the cross section by $K_{gg} \simeq 1.5$ and $K_{q\bar{q}} \simeq 1.2$ \cite{10,246}. Some authors also consider SM vectors as partons, e.g. $C_{\gamma\gamma} \sim 10$ (60) at $\sqrt{s} = 8$ (13) TeV \cite{141,29,10,178,282}. The diagonal lines in fig. 2 shows the ratio of cross-section $\sigma_{13}/\sigma_8$ predicted by the various partons: data disfavor the partons that give the smallest enhancement (light quarks and SM vectors), favouring $F$ production from heavy quarks or gluons — and an even larger $\sigma_{13}/\sigma_8$ enhancement would give a better fit.$^c$

This can be achieved if $pp$ collisions produce some heavier particle (e.g. a heavy vector with mass $\sim 1.5$ TeV) that decays into $F$ and something invisible \cite{7,10,46,69,62,298,319}, with a phase space almost closed in order to reproduce the lack of extra particles and transverse momentum in the $\gamma\gamma$ excess events. Such kinematics can be used to fake a large $F$ width \cite{107}. A large $F$ width can be faked in other ways: by having two or more nearby narrow resonances (for example the scalar and pseudo-scalar components of a SU(2)$_L$ doublet splitted by $v^2/M_F \sim 40$ GeV, where $v$ is the Higgs vev) \cite{10,174}, or by assuming that $F$ decays into pairs of light particles with mass $m < \sim \text{GeV}$, that decay into two or more photons collimated within an angle $\theta \sim m/M_F$, such that they appear in the detector as a single $\gamma$ \cite{42,56,72,189,222,298}. This latter possibility allows to get a large tree-level...
Fig. 3 shows the $f$ widths that reproduce the $\gamma\gamma$ excess. In the left (right) panel we assumed the production process with the largest (smallest) partonic luminosity, namely $gg$ ($b\bar{b}$). The main lesson is that there is a minimal value of $\Gamma_{\gamma\gamma}$, obtained assuming that $f$ has a small width and is dominantly produced from $gg$: restricting fig. 2a along the $gg$ diagonal line one finds

$$\sigma(pp \rightarrow f \rightarrow \gamma\gamma) = (2.8 \pm 0.7) \text{ fb}$$

such that

$$\frac{(2J+1)}{M_f} \frac{\Gamma_{\gamma\gamma}}{\Gamma_f} = \frac{s}{K_{gg} C_{gg}} \sigma(pp \rightarrow F \rightarrow \gamma\gamma) = (3.8 \pm 0.9) \times 10^{-7}$$

in agreement with the blue region in fig. 3a. A larger $\Gamma_{\gamma\gamma}$ is needed if $f$ has a larger width (yellow regions) and/or is produced from other partons. For example, one needs $\Gamma_{\gamma\gamma}/M_f \gtrsim 10^{-4}$ (green regions) if $\Gamma_f/M_f \sim 0.06$ as favoured by ATLAS. Finally, reproducing $\sigma(pp \rightarrow F \rightarrow \gamma\gamma)$ assuming that production from partonic photons dominates (a possibility disfavoured by data at $\sqrt{s} = 8 \text{ TeV}$) needs the largest $\Gamma_{\gamma\gamma}/M_f \sim 10^{-3}$.

The global fits in fig. 3 take into account the experimental bounds on other $\sigma(pp \rightarrow F \rightarrow f)$ with final states $f$, as reported in table 1.

### 3 Effective Lagrangian

So far, $\sigma(pp \rightarrow F)$ and $F$ decays have been described simply in terms of $F$ widths. In order to compute extra related processes we need to make extra theoretical assumptions. However the Lagrangian interactions of $F$ differ — even in their dimensionality — depending on the unknown $\Gamma_f$ ($\gamma\gamma$–like) and can be tested by better studying the $\gamma$ events (multiple $\gamma$ traveling in the material before the electromagnetic calorimeter give more $\gamma \rightarrow e^+e^-$ conversions than a single photon [204], and with different kinematical features [352]); furthermore it can lead to a displaced vertex, and to no decays into other electroweak vectors.
Table 1: Bounds at 95% confidence level on $\sigma(pp \to F \to f)$ cross sections for various final states $f$. We here assumed $\Gamma/M_F \approx 0.06$.

Lorentz and gauge quantum numbers of $F$:

$$ F \equiv \begin{pmatrix} \text{spin 0} \\ \text{spin -1} \\ \text{spin 2} \end{pmatrix} \times \begin{pmatrix} \text{SU}(2)_L \text{ singlet} \\ \text{SU}(2)_L \text{ doublet} \\ \text{SU}(2)_L \text{ triplet} \end{pmatrix} \times \begin{pmatrix} \text{CP-even} \\ \text{CP-odd} \\ \text{CP-violating} \end{pmatrix} \cdots $$

We proceed assuming that $F$ is a neutral singlet with spin 0, either scalar or pseudo-scalar. Then, the renormalisable interactions of $F$ are

$$ \mathcal{L}_4 = \mathcal{L}_{SM} + \frac{(\partial_\mu F)^2}{2} - V(F,H), $$

where

$$ V(F,H) = \frac{m_F^2}{2} F^2 + \kappa_F m_F F^4 + \lambda_F F^4 + \kappa_{FH} m_F F^2 (|H|^2 - v^2) + \lambda_{FH} F^2 (|H|^2 - v^2). $$

This does not give an acceptable $F \to \gamma\gamma$, so we include dimension 5 non-renormalizable interactions, which are a good approximation to a generic unknown more complete theory with extra particles that mediate $F \to \gamma\gamma$, provided that such extra particles are much heavier than $M_F$. For sure they are much heavier than $M_h$, so we write SU(2)$_L$-invariant effective operators. Getting rid of redundant operators, the most generic effective Lagrangian is

$$ \mathcal{L}_5^{\text{even}} = \frac{F}{\Lambda} \left[ c_{gg} g_3^2 G_{\mu\nu}^a G^{a\mu\nu} + c_{WW} g_2^2 W_{\mu\nu}^a W^{a\mu\nu} + c_{BB} g_1^2 B_{\mu\nu} B^{\mu\nu} + c_{\psi} (H \bar{\psi} \psi_R + \text{h.c.}) \right] + c_H |D_\mu H|^2 - c'_H (|H|^4 - v^4) + c_{3F} \frac{F(\partial_\mu F)^2}{\Lambda}, $$

$^{d}$A spin 2 graviton is disfavoured because it couples universally to the conserved energy momentum tensor, such that $\sigma(pp \to F \to e^+ e^- + \mu^+ \mu^-) = \sigma(pp \to F \to \gamma\gamma)$, but no peak is seen in leptons: bounds are reported in table 1 [10]. A spin 2 resonance can be resurrected by assuming that it couples to $\gamma$ more strongly than to leptons; however this zombie has gauge-dependent cross sections enhanced by inverse powers of $M_F$ (in effective theories one can restrict to regions of the parameter space where unphysical terms are small). Angular distributions allow to discriminate spin 0 from spin 2 [48, 182, 258, 262, 196, 321]. Bound states with spin 2 which have nothing to do with gravity can couple differently to different particles, and can have odd parity (while a graviton is even), leading to different angular distributions, which are as motivated as spin 3 bound states [232].
for CP-even $f$, while

$$\mathcal{L}_\text{odd}^\gamma = \frac{f}{\Lambda} \left[ e_{gg} \frac{g_2}{2} G^{\alpha \mu \nu} G_{\alpha \mu \nu} + e_{WW} \frac{g_2}{2} W^{\alpha \mu \nu} \bar{W}^{\alpha \mu \nu} + e_{BB} \frac{g_2}{2} B_{\mu \nu} \bar{B}^{\mu \nu} + e_{\psi} \left( i H \bar{\psi}_L \psi_R + \text{h.c.} \right) \right]. \quad (9)$$

in the CP-odd case. Operators with higher dimension have been listed in [282,286]. Couplings to SM fermions $\bar{\psi}_{L,R}$ are restricted by flavour bounds, which imply that $f$ can have large couplings only to pairs of mass eigenstates $t, b, c, \ldots$

Defining $c_{\gamma \gamma} = c_{BB} + c_{WW}$, the above effective Lagrangian leads to the decay widths

$$\Gamma_{\gamma \gamma} = \frac{\pi \alpha^2 M_f^3}{\Lambda^2} (c_{\gamma \gamma}^2 + c_{\gamma g}^2), \quad \Gamma_{gg} = K_{gg} \frac{8\pi \alpha^2 M_f^3}{\Lambda^2} (c_{gg}^2 + c_{gg}^2), \quad (10)$$

($K_{gg} = 1.35$ when the other couplings are renormalised at $\bar{\mu} = M_f$) together with other decays modes with characteristic rates

$$\Gamma_{hh} = \frac{M_f^3}{128\pi \Lambda^2 c_H^2}, \quad \Gamma_{ZZ} = \frac{\pi \alpha^2 M_f^3}{\Lambda^2 s_W^2 c_W^4} (c_{ZZ}^2 + c_{Z\gamma}^2) + \Gamma_{hh},$$

$$\Gamma_{Z\gamma} = \frac{2\pi \alpha^2 M_f^3}{s_W^2 c_W^2 \Lambda^2} (c_{Z\gamma}^2 + c_{Z\gamma}^2), \quad \Gamma_{WW} = \frac{2\pi \alpha^2 M_f^3}{\Lambda^2 s_W^2} (c_{WW}^2 + c_{WW}^2) + 2\Gamma_{hh},$$

where $c_{\gamma Z} = s_W^2 c_{BB} - c_Z^2 c_{WW}$, $c_{ZZ} = s_W^2 c_{BB} + c_W^2 c_{WW}$ and $\hat{c}_H = c_H + 2\kappa_M H A/M_f$. We neglected terms suppressed by $M_{W,Z,h}/M_f$, computed in [282].

Experiments only set upper bounds on these extra decay modes. The bounds are satisfied assuming, for example, that $f$ couples to hypercharge only. A $f$ coupled to SU(2)$_L$ vectors only is instead disfavoured by bounds on $f \rightarrow WW, ZZ$, especially if $f$ is broad. The limiting values are

| operator | $\Gamma_{Z\gamma}/\Gamma_{\gamma \gamma}$ | $\Gamma_{ZZ}/\Gamma_{\gamma \gamma}$ | $\Gamma_{WW}/\Gamma_{\gamma \gamma}$ |
|----------|--------------------------------------|----------------------------------|----------------------------------|
| $c_{WW}$ only | $2/\tan^2 \theta_W \approx 7$ | $1/\tan^4 \theta_W \approx 12$ | $2/\sin^2 \theta_W \approx 40$ |
| $c_{BB}$ only | $2 \tan^2 \theta_W \approx 0.6$ | $\tan^4 \theta_W \approx 0.08$ | $0$ |

Figure 4 – Best-fit regions at 68.90% C.L. (green regions) for the coefficients of the operators that control $f$ decays, assuming that $f$ is narrow (left) or broad (right) and produced as $gg \rightarrow f$. We also show isocurves of $\Gamma(f \rightarrow Z, \gamma Z, WW, hh)/\Gamma(f \rightarrow \gamma \gamma)$ for $f = ZZ$ (red), $f = \gamma Z$ (green dashed), $WW$ (blue dotted), $hh$ (black dot-dashed).
Fig. 4 shows the best-fit region in the \((c_{WW}/c_{BB}, \hat{c}_H/c_{BB})\) plane. Future observations of extra decay modes will over-constrain the fit. In particular, the \(F\) width into \(\gamma Z\) or the width into \(ZZ\) can be fine-tuned to zero, but not both; the diphoton cannot be only a diphoton. The minimal extra effect \(\Gamma_{\text{extra}} > 0.28\Gamma_{\gamma\gamma}\) is obtained for \(c_{WW} \approx 0.04\).

Concerning the decays into Higgs components (the physical \(h\), and the 3 Goldstones eaten by the \(W^\pm\) and the \(Z\)), they can be equivalently described in terms of a mixing angle \(\theta_{hf}\) between the Higgs mass eigenstate \(h\) and the diphoton:

\[
\tan 2\theta_{hf} = \frac{2v(m_{hf} + c_{hf}^2/\Lambda)}{m_h^2 - m_{hf}^2}.
\]

The experimental bound \(\Gamma(F \rightarrow ZZ) < 20\Gamma(F \rightarrow \gamma\gamma)\) implies that such angle is small:

\[
|\sin \theta_{hf}| < 0.015\sqrt{\frac{\Gamma(F \rightarrow \gamma\gamma)}{10^{-6}M_F}}.
\]

### 4 The everybody’s model

Many renormalizable models can realise one or more of the effective operators: for example couplings to fermions can be mediated at tree level by one extra Higgs doublet or by extra vector-like fermions. A different class of models attracted most attention: those where the new particles that mediate \(F \rightarrow \gamma\gamma\) (the only process mandated by data) also mediate \(gg \rightarrow F\). Such renormalizable models are obtained adding to the SM a neutral scalar \(F\) and extra charged fermions or scalars \(Q\) in order to mediate the effective \(F\) couplings to vectors in the following way:

\[
\Gamma_{gg} \sim 7.2 \times 10^{-5} X_{gg}^2, \quad \Gamma_{\gamma\gamma} \sim 5.4 \times 10^{-8} X_{\gamma\gamma}^2
\]

where \(X_{gg} \sim N_q y M_F / M_Q\) and \(X_{\gamma\gamma} \sim N_q^2 y M_F / M_Q\) are loop functions that only contain model-dependent order one factors: the multiplicity \(N\) of states \(Q\), their color or electric charge \(q\), the strength \(y\) of their coupling to \(F\), and their mass. We do not report here the well known expressions for the loop factors, see e.g. [10].

The resulting decay widths are

\[
\Gamma_{gg} = \frac{\alpha^2}{2\pi^3} X_{gg}^2 \sim 7.2 \times 10^{-5} X_{gg}^2, \quad \Gamma_{\gamma\gamma} = \frac{\alpha^2}{16\pi^3} X_{\gamma\gamma}^2 \sim 5.4 \times 10^{-8} X_{\gamma\gamma}^2
\]

\[\text{where} \quad X_{gg} \sim N_q y M_F / M_Q \text{ and } X_{\gamma\gamma} \sim N_q^2 y M_F / M_Q \text{ are loop functions that only contain model-dependent order one factors: the multiplicity } N \text{ of states } Q, \text{ their color or electric charge } q, \text{ the strength } y \text{ of their coupling to } F, \text{ and their mass.} \]

We do not report here the well known expressions for the loop factors, see e.g. [10]. If \(F\) is a pseudo-scalar, the fermion loop is resonantly enhanced when \(M_Q = M_F / 2\) [233, 251, 60]. A useful general result holds if \(F\) is a scalar much lighter than \(Q\); the \(F\) coupling to vectors are determined by the contribution \(\Delta b^Q\) of particles \(Q\) to the gauge beta functions as

\[
\mathcal{L}_{\text{eff}} = \sum_{i,Q} \Delta b^Q \frac{\alpha_i}{8\pi} (F_{\mu\nu}^i)^2 \ln \frac{M_Q(F)}{M_Q}
\]

\footnote{A resonance that apparently decays only to \(\gamma\gamma\) is possible if one or both photons actually are collimated jets of photons: models that realise this have been presented in footnote c.}
where $M_Q(F)$ is the $Q$ mass for a generic vev of $F$. It can be computed in any given model, and expanding it at first order in $F$, in $M_Q(F)/M_Q \approx F/v_Q$, gives the desired $F$ coupling to vectors in terms of the model-dependent constant $v_Q$. The main message is that

order one charges, multiplicities and couplings in eq. (14) can reproduce the value of $\Gamma_{\gamma\gamma}/M_F \sim 10^{-6}$ suggested by the measured $\sigma(pp \to F \to \gamma\gamma)$ assuming that the diphoton is narrow.

The conclusion drastically changes if instead the diphoton is broad. ATLAS gives a $\sim 1\sigma$ hint in favour of $\Gamma/M_F \sim 0.06$. Such a large width, by itself, would not be a problem: it can be obtained as a tree level two-body decay with a order one coupling, analogous to the top Yukawa coupling. The bounds on $F$ decays of table 1 allow for a large $F$ decay width into jets and/or invisible channels, such as neutrinos or Dark Matter. However, if $F$ is broad, larger values of $\Gamma_{\gamma\gamma}/M_F \sim 10^{-3-4}$ are needed to reproduce $\sigma(pp \to F \to \gamma\gamma)$. This is the problem: according to eq. (14), in order to reproduce such a large $\Gamma_{\gamma\gamma}$, $y$ and/or $q$ and/or $N$ need to be large. Apart from plausibility issues, all these possibilities lead to some coupling becoming non-perturbative and hitting a Landau pole at energies not much above the diphoton mass [90,105,116,10,156,193,200,202]. A large electric charge $q \sim 3$ or $N \gg 1$ states with $q \sim 1$ implies a fast running of the hypercharge gauge coupling and of the strong coupling (if they are colored); a large Yukawa or a large scalar quartic renormalize themselves to larger values at higher energy; a large scalar cubic has the same problem, once vacuum stability bounds are taken into account [193].

The Landau pole can be delayed or avoided if some extra comparably strong interaction is present. For example, in the Standard Model the top Yukawa coupling $y_t \approx 1$ does not hit a nearby Landau pole because the strong gauge coupling $g_3 \sim y_t$ keeps the running $y_t$ under control, providing a RGE flow with an infra-red fixed point. Something similar can allow a larger $F$ width, in models where $N$ states with $q \sim 1$ lie in the fundamental of a new gauge group such as $SU(N)$ [193].

While it’s premature to build models based on a $\sim 1\sigma$ hint in favour of a large width, models with a new strong interaction have been explored because of their own interest.

5 Composite diphoton

The above situation prompted many authors to consider strongly-coupled models, where the diphoton is a composite resonance.

Ref.s [55,219,103,220,293] explored the simplest possibility that $F$ is a QCD bound state of heavy quarks $\bar{Q}Q$ with mass $M_Q \approx \frac{1}{4}750$ GeV. At this energy $\alpha_3 \approx 0.100$ is relatively small, such that the QCD binding energy is small: only a small fraction of the produced $\bar{Q}Q$ pairs manifests as a 750 GeV resonance, as well as inducing extra features in the $\gamma\gamma$ spectrum [103,318]. The $\gamma\gamma$ excess can be reproduced if uncertain QCD factors are favourable, if $Q$ decays with a life-time longer than the life-time of the bound state, and into a final state not subject to strong bounds.

These difficulties can be alleviated adding a new strong interaction that confines just below $M_Q [10,93,254,288]$. In the limit where the $\bar{Q}Q$ potential can be approximated as $V(r) = -\alpha_{eff}/r$ where $\alpha_{eff}$ is some effective constant larger than $\alpha_3$, the quarkonium-like resonances have mass

\footnote{This can be experimentally tested trough high-energy tails of $pp \to \ell^+\ell^-$ distributions at LHC [201,205], which probe the SM electroweak couplings renormalised at $m_{\ell\ell} \sim 2$ TeV, and trough precision measurements around the $Z$-peak at a future $e^+e^-$ collider.}
\[ M_n = 2M_Q(1 - \alpha^2_{\text{eff}}/8n^2) \]. The lightest resonance with \( n = 1 \) is identified with \( F \) and its decay widths are

\[ \Gamma_{\gamma\gamma} = \frac{q^4 N}{4} \frac{\alpha^2}{\alpha_{\text{em}}} \frac{1}{\alpha_{\text{eff}}} = 10^{-6} N q^4 \left( \frac{\alpha_{\text{eff}}}{0.4} \right)^3, \quad \Gamma_{gg} = \frac{2\alpha^2}{9q^4} \alpha^3_{\text{em}} - \Gamma_{\gamma\gamma} \]

having assumed that the particles \( Q \) are \( N \) color triplets with charge \( q \). The slightly heavier resonances with \( n > 1 \) have smaller widths, \( \Gamma_{\gamma\gamma}^{(n)} \propto 1/n^3 \), giving a characteristic pattern. However, since \( 3 \otimes \bar{3} = 1 \oplus 8 \), models based on a new strong interaction predict that each neutral resonance is accompanied by a quasi-degenerate color octet resonance. QCD repulsion reduces its binding energy, making its production cross sections less problematic.

Many authors explored the possibility that \( F \) is a bound state of a new strong interaction. There are three main classes of models:

1. **Models where \( H \) and \( F \) are composite.** [212, 304] This can be realised in simple fundamental models, where a new TechniColor (TC) gauge interaction (for example with gauge group \( SU(N) \)) becomes strong around the weak scale, and TC quarks are chiral under \( SU(2)_L \). While the diphoton is totally natural, these models have big problems in reproducing higgs, electro-weak and flavor data.

2. **Scenarios where \( H \) and \( F \) are partially composite** [10, 80, 26, 172, 322] postulate chiral effective Lagrangians with the needed properties that allow to bypass the TC problems, ignoring the issue of finding a fundamental dynamics that realises them. The lightness of the Higgs is interpreted assuming that it is the pseudo-Goldstone boson of an accidental global symmetry broken by unknown dynamics, and often the pattern of symmetry breaking gives extra light singlets, that can be identified with \( F \). Such models tend to give \( F \to t\bar{t} \) decays.

3. **Models where \( F \) is composite** [2, 6, 10, 19, 31, 38, 70, 158, 191, 214, 245, 257, 280, 287, 312] This can be realised in simple fundamental models where a new TC gauge interaction becomes strong around the weak scale, and the TC particles are not chiral under the SM gauge group, that is thereby left unbroken by the TC dynamics. These models use an elementary Higgs doublet, like the SM, and are thereby equally compatible with electro-weak, higgs and flavor data.\(^h\)

Roughly speaking, the first class of models are dead, the second class are never born, so we focus on the third class. In order to obtain both \( F \to \gamma\gamma \) and \( gg \to F \) the TC particles \( Q \) must be both colored and charged. Then \( \bar{Q}Q \) necessarily contains SM singlets, that can be identified with \( F \), but also color octets, subject to strong LHC bound [289]. If \( M_Q \gtrsim \Lambda_{\text{TC}} \) such models realise the quarkonium-like scenario of eq. (16), which should be accompanied by a quasi-degenerate color octet. More plausible realisations thereby identify \( F \) with a bound state that is much lighter than the others. The main possibilities are:

- \( F \) as TC\( \eta \). If the TC dynamics breaks an accidental global symmetry, a set of \( \bar{Q}Q \) pseudo-scalar bound states remains light, being the TC-pions. Some of them are neutral TC\( \eta \).

\(^h\)Different authors use different names such as ‘dark’, ‘hyper’, ‘hidden’, ‘big’ to distinguish the new gauge interaction from TechniColor. Since there is no unique name, we use the old-fashioned name Techni Color.

\(^h\)Before the announcement of the diphoton excess, models of automatically stable composite Dark Matter based on such dynamics were explored, even in papers that mentioned \( pp \to F \to \gamma\gamma \) signals.
Their couplings to SM vectors $V, V'$ are given in terms of the TC-pion decay constant $f_{TC}$ and of the gauge quantum numbers of the TC particles $Q$ as

$$\frac{c_{VV'}}{\Lambda} = \frac{\kappa_{VV'}}{8\pi^2 f_{TC}}, \quad \kappa_{VV'} = N \text{Tr} (T_f T^V T^{V'}).$$

(17)

where $T^V$ are the generators of the SM gauge group and $T_f$ is the chiral symmetry generator associated to $F$ (for example $T_f = 1/\sqrt{2N}$ corresponds to the TC-$\eta'$ state that gets a mass of order $\Lambda_{TC}$ from TC anomalies). So

$$\frac{\Gamma_{\gamma\gamma}}{M_F} = \frac{\alpha_{em}^2 \kappa_{\gamma\gamma} M_F^2}{64\pi^3 f_{TC}^2} = 10^{-6} \left( \frac{\kappa_{\gamma\gamma}}{f_{TC}} \right)^2\left( \frac{120 \text{ GeV}}{f_{TC}} \right)^2$$

(18)

where $\kappa_{\gamma\gamma} = \kappa_{BB} + \kappa_{WW}$. Measuring other $F \rightarrow VV'$ decays would allow to infer the techni-particle content $Q$.

- $F$ as TC$\sigma$ or dilaton. Another state that can be especially light is the scalar pseudo-Goldstone boson of scale invariance [10, 16, 24, 44, 159, 243, 278]. Scale invariance is a good symmetry if two conditions are satisfied. a) no TC particles have masses around $\Lambda_{TC}$, unlike in QCD. b) the TC-strong dynamics is in a ‘walking’ regime, unlike the QCD dynamics where $\alpha_3$ ‘runs’ to non-perturbative values. These conditions imply an appropriate content of light TC particles $Q$.

From a low-energy perspective, the light TC-dilaton is the $\sigma$ field sometimes explicitly included in effective chiral Lagrangians [10, 243]. Its coupling to SM vectors is dictated by eq. (15) such that

$$\frac{\Gamma_{\gamma\gamma}}{M_F} = 10^{-6} \left( \frac{\Delta b_{em}}{f_{TC}} \frac{120 \text{ GeV}}{f_{TC}} \right)^2$$

(19)

where $\Delta b_{em}$ is the contribution to the running of the electromagnetic coupling from techni-particles $Q$.

Coming back to the issue of a large $F$ total width, $\Gamma \sim 0.06 M_F$ can be realised adding extra $F$ decays to SM particles or into other techni-pions, which can include Dark Matter candidates. However, a large $\Gamma$ needs a large $\Gamma_{\gamma\gamma} \sim 10^{-3-4} M_F$: a look at all expressions for $\Gamma_{\gamma\gamma}$ in composite $F$ models shows achieving such a large $\Gamma_{\gamma\gamma}$ remains difficult.

6 What next?

1) Does $F$ have spin 0, 2, or more?
2) Is $F$ a SU(2)$_L$ singlet or doublet or something else?
3) Is $F$ produced through $gg$, $q\bar{q}$ or weak vector collisions?
4) Is $F$ narrow or broad? How large are its couplings and to which particles does $F$ couple?
5) Is $F$ CP-even or CP-odd or its couplings violate CP?
6) Is $F$ elementary or composite?
7) Is $F$ a cousin of the SM scalar?
8) Does $F$ exist?
In the case of the SM scalar, these kind of questions had ‘a similar potential for surprise as a football game between Brazil and Tonga’. In the case of $Ϝ$ they are as open as a match between Brazil and Germany. Many works explored how to ask these questions to nature and get answers [43, 73, 93, 164, 209, 215, 267, 275, 282, 283, 310, 296, 311, 313, 326, 328, 329, 347, 349]. The main ideas are summarised below.

1) The **spin** can be identified in the following ways:

1a) spin 1 and half-integer spin are already excluded by the observation of $Ϝ \rightarrow \gamma\gamma$.

1b) from angular distributions, as well known.

1c) a particle with spin 2 or higher can only be a bound state, that should come with other ones.

We will focus on spin 0.

2) The **weak representation** of $Ϝ$ can be identified in the following ways:

2a) if $Ϝ$ is not a neutral singlet, its extra charged components must be around 750 GeV. Find them.

2b) identifying the production mode: a singlet can couple at dimension 5 to all SM particles, while a doublet is more likely to be produced from quarks, to which it may have renormalisable couplings.

2c) measuring the $p_T$ spectra in $Ϝ$ associated production, in view of the different dimensionalities of the effective couplings to quarks (dimension 5 for singlet $Ϝ$ and 4 for doublet) and gauge bosons (dimension 5 for singlet and 6 for doublet).

3) The **initial state** that produces $Ϝ$ can be identified in the following ways:

3a) measuring how $\sigma(pp \rightarrow F)$ depends on energy; data at $\sqrt{s} = 8$ vs 13 TeV already disfavour production from light quarks or photons.

| $\sqrt{s} = 13$ TeV | $b\bar{b}$ | $c\bar{c}$ | $s\bar{s}$ | $u\bar{u}$ | $d\bar{d}$ | $GG$ |
|---------------------|--------|--------|--------|--------|--------|------|
| $\sigma_f / \sigma_f$ | 9.2%   | 7.6%   | 6.8%   | 6.7%   | 6.2%   | 27%  |
| $\sigma_{f\bar{b}} / \sigma_f$ | 6.2%   | 0      | 0      | 0      | 0      | 0.32%|
| $\sigma_{f\bar{c}} / \sigma_f$ | 1.4%   | 1.0%   | 0.95%  | 1.2%   | 1.0%   | 4.7%  |
| $\sigma_{f\bar{s}} / \sigma_f$ | 1.2%   | 0.18%  | 0.19%  | 0.34%  | 0.31%  | 0.096%|
| $\sigma_{f\bar{u}} / \sigma_f$ | 0.31%  | 0.17%  | 0.18%  | 0.34%  | 0.31%  | 0.024%|
| $\sigma_{f\bar{d}} / \sigma_f$ | 0.37%  | 1.5%   | 0.38%  | 1.6%   | 0.41%  | $< 10^{-6}$|
| $\sigma_{fZ} / \sigma_f$ | 1.1%   | 1.1%   | 1.3%   | 2.0%   | 1.9%   | 3 $10^{-6}$ |
| $\sigma_{fW^+} / \sigma_f$ | 5 $10^{-5}$ | 1.7% | 2.4%   | 2.6%   | 4.1%   | $< 10^{-6}$ |
| $\sigma_{fW^-} / \sigma_f$ | 3 $10^{-5}$ | 2.3% | 1.2%   | 1.0%   | 1.7%   | $< 10^{-6}$ |
| $\sigma_{fh} / \sigma_f$ | 1.0%   | 1.1%   | 1.2%   | 1.9%   | 1.8%   | 1 $10^{-6}$ |

Table 2: *Predictions for the associated production of the resonance $Ϝ$, assuming the effective couplings of section 3, whose validity is far from guaranteed. We assumed the standard cuts $\eta_j < 5$, $p_T \eta > 150$ GeV, $\Delta R_{jj} > 0.4$ on jets, and $\eta_{\gamma} < 2.5$, $p_T, \gamma > 10$ GeV on photons.*
3b) any partonic $\varphi \rightarrow F$ production process implies a corresponding $F \rightarrow \varphi$ decay. Find it.

3c) the rapidity distribution and the transverse momentum spectrum of the diphoton system retain features of the initial parton state \[113\].

3d) from the amount of extra jets from initial-state radiation in $pp \rightarrow F$, see table 2.

3e) production from $b$ quarks implies $\sigma(pp \rightarrow F b) \approx 6\% \sigma(pp \rightarrow F)$ within the effective theory.

3f) $F$ production in association with a gauge or Higgs boson is a useful discriminator, see table 2. In particular, no vector bosons accompanying $F$ are expected from gluon initial states, and no $W$ from $b$ initial states. Ratios of $F W$, $F Z$, and $F h$ provide additional handles to identify the production process.

3g) if $F$ is a singlet produced from quarks, the $F \bar{q}qH$ operator implies a sizeable three-body decay width, $\Gamma(F \rightarrow q\bar{q}H) \sim 1\% \times \Gamma(F \rightarrow q\bar{q})$ where $H = \{h, Z, W^\pm\}$.

4) The $F$ couplings can be measured in the following ways:

4a) If $F$ is broad enough that its total width can be measured, such that the couplings can be reconstructed from the branching ratios.

4b) If $F$ is broad, it might even be possible to measure interference with SM cross sections \[133, 282, 296, 307, 328\].

4c) SU(2)$_L$-invariance relates different decay widths, allowing to disentangle the effective operators.

4d) $F$ couplings to DM can be accessed from missing energy signals in the usual ways.

4e) Associated processes such as $pp \rightarrow F j$ or $pp \rightarrow F V$ probe the energy dependence of the couplings.

4f) Observation of $pp \rightarrow FF$, would imply relatively large couplings: either $F^3$ cubics, or of $F$ to the particles $Q$ that mediate $F \rightarrow \gamma\gamma$ (this effect can be computed by expanding eq. (15) to order $F^2$):

5) The CP parity of $F$ can be identified in the following ways:

5a) From angular distributions of $F \rightarrow \gamma^+\gamma^- \rightarrow \ell^+\ell^-\ell'^+\ell'^-$; however its rate is 3 orders of magnitude below the $\gamma\gamma$ rate.

5b) From $F \rightarrow \gamma\gamma$ using $\gamma \rightarrow e^+e^-$ conversions in the detector matter; however the angle between the $e^\pm$ pairs is very small.

5c) From the angular distribution of $pp \rightarrow jjF$ events.

5d) If $F \rightarrow ZZ$ exists, the CP parity of $F$ can be measured from angular distributions in leptonic $Z$ decays and in $pp \rightarrow ZF$.

5e) If $F \rightarrow Z\gamma$ exists, something can be done combining 5a) with 5d).
5f) If $F \rightarrow hh$ exists, it implies that $F$ is a scalar.

5g) If $F \rightarrow hZ$ exists, it implies that $F$ is a pseudo-scalar, and that the effective theory approximation fails, given that low-dimension operators do not induce such decay.

7 Connection with Dark Matter, axions, vacuum stability, baryogenesis...

Possible connections of $F$ with other open issues have been investigated.

Various authors explored the possibility that the diphoton is the mediator that couples Dark Matter to SM particles [3, 4, 49, 66, 81, 101, 124, 149, 171, 234, 242, 279, 305, 309, 320, 323]. The freeze-out DM relic abundance can reproduce the observed cosmological DM abundance for natural values of the parameters, as it is customary for weak-scale particles. In particular, if the diphoton decays into Dark Matter, one needs a DM mass around 100-300 GeV, depending on the diphoton width. DM direct detection is somewhat below present bounds if the diphoton is a scalar, or suppressed by non-relativistic factors if the diphoton is a pseudo-scalar.

Adding to the SM field an axion allows to understand the smallness of the CP-violating $\theta$ angle of QCD. However the axion must be ultra-light and coupled ultra-weakly: it cannot be identified with $F$. Nevertheless, some authors tried to identify the $F$ resonance with an axion-like state, by building models where it can be heavy and significantly coupled [9, 112, 120, 189, 218, 265, 330]. For example [265] tries to realise a Nelson-Barr-like model at the weak scale.

The addition of $F$ can eliminate the instability of the SM vacuum [65, 97, 102, 316] in two different ways: a) by providing a tree-level threshold corrections that increases the Higgs quartic $\lambda_H$; b) the extra charged particles that mediate $F \rightarrow \gamma\gamma$ modify the RGEs, such that $g_{2,Y}$ and consequently $\lambda_H$ become larger at large energy.

The electro-weak phase transition, extended including $F$, could become of first order leading to gravitational waves and baryogenesis [166, 168, 235, 271, 294, 316, 317]. Other works discussed connections with neutrino masses [135, 161, 171, 253, 334], flavor [71, 119, 128, 274, 340], inflation [65, 123, 131, 338], extra dimensions [33, 48, 110, 57, 92, 177, 196, 236, 249, 252, 262, 350], and string/$F$-theory [64, 89, 117, 120, 142, 134, 163, 221, 248, 255, 325, 327, 357].

8 Who ordered that?

After that experiments will answer the questions of section 6, clarifying what $F$ really is, it will be possible to understand which role $F$ plays in particle physics.

In the meantime, various authors started to explore the possibility that $F$ has something to do with the origin of the electro-weak scale, trying to identify $F$ with one or another supersymmetric particle: sneutrino [10, 52, 91], extra scalar or pseudo-scalar Higgs [10, 109, 129, 297, 333], extra NMSSM singlet [10, 224, 106, 129, 145, 180, 197, 273], sgoldstino [16, 18, 27, 98, 190, 240, 244] (its rate in $\gamma\gamma$ implied by gaugino masses seems too small), stopponia [293, 345], sbino [45, 308] or else [58, 99, 140, 217]. In this context, extra full SU(5) multiples around the weak scale can enhance the $\gamma\gamma$ rate [84, 99, 145, 273, 300, 344]. Connections with other solutions to the hierarchy problem have been also explored, such as composite or partially composite Higgs or extra dimensions.

A related theoretical issue is the naturalness of the extra charged particles introduced to mediate $F \rightarrow \gamma\gamma$. Even if they are fermions, they have no chirality reason to be around the
weak scale, unlike the SM fermions. Supersymmetry and other solutions to the hierarchy problem imply extra charged particles at the weak scale. Indeed, such extensions of the SM tame quadratically divergent corrections to the Higgs mass at the price of introducing a lot of new physics at the weak scale.

However, such new physics has not been observed, and bounds relegate solutions to the hierarchy problem to fine-tuned corners of their parameter space. Some authors, following the point of view that quadratic divergences give no physical effects, explored models where no physical correction is unnaturally large and tried to build extensions of the SM where a hierarchically small weak scale is induced by some new dynamics. Refs. [59, 78, 131, 316] explored the possibility that $\mathcal{F}$ is (a manifestation of) such dynamics: the smoking gun of this scenario would be observing that $\mathcal{F}$ couples to all particles proportionally to their masses. In this context, broken scale invariance can justify extra charged particles, which cannot be much above the weak scale in order to avoid unnaturally large physical corrections to the Higgs mass.

9 Conclusions

The $\mathcal{F} \to \gamma\gamma$ excess can be a statistical fluctuation, or a the first sign of new physics. In the latter case various reasonable models can reproduce it, at least if the $\mathcal{F}$ width is narrow. A generic prediction is that we expect to see also $\mathcal{F} \to \gamma Z$ and/or $\mathcal{F} \to ZZ$ and extra charged particles. All the rest is model-dependent. Today the diphoton excess could be everything, including nothing.

Note Added The above discussion is a disproportionate amplification of a quantum fluctuation: no 750 GeV $\gamma\gamma$ excess is present in the first (12.2 + 12.9) fb$^{-1}$ of new 2016 LHC data [371], which confirm the Standardissimo Model and the bad reputation of the digamma symbol $\mathcal{F}$.

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*Unless an enlarged gauge symmetry broken around the weak scale, such as SU(3)$_L \otimes$ U(1) $\otimes$ SU(3)$_c$, [74, 76, 111, 331] or SU(3)$_c \otimes$ SU(3)$_R \otimes$ SU(3)$_c$, [78] or SU(6) $\to$ G$_{SM}$ $\otimes$ U(1) [290], provides an extended set of charged chiral fermions. Furthermore, SU(3 + $N$) $\to$ SU(3), provides extra charged scalars [336].
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