Physics Letters B 757 (2016) 148–153

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Could the width of the diphoton anomaly signal a three-body decay?

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ARTICLE INFO

Article history:
Received 22 December 2015
Received in revised form 8 February 2016
Accepted 24 March 2016
Available online 30 March 2016
Editor: G.F. Giudice

ABSTRACT

The recently observed diphoton anomaly at the LHC appears to suggest the presence of a rather broad resonance. In this note, it is pointed out that this broadness is not called for if the two photons are produced along with an extra state. Specifically, the diphoton invariant mass arising from various $A \to B\gamma\gamma$ processes, with A, B being scalars, fermions, or vectors, though peaked at a rather large value, would naturally be broad and could fit rather well the observed deviations. This interpretation has a number of advantages over the two-photon resonance hypothesis, for example with respect to the compatibility with the 8 TeV diphoton, dilepton or dijet searches, and opens many new routes for New Physics model construction.

1. Introduction and set-up

Recently, a small deviation in the diphoton mass spectrum was announced by both ATLAS [1] and CMS [2] at a mass of around 750 GeV. While the statistical significance of this signal is still low, the simultaneous observation by both experiments lends some credence to the presence of a yet unknown resonance in this channel, and has led to an incredibly intense phenomenological activity (see Refs. [4–41]).

In this note, we want to point out that one feature of this $\gamma\gamma$ signal, namely its width, could be well explained if it arises from a three-body decay $A \to B\gamma\gamma$, with the mass splitting $M_A - M_B$ a bit higher than 750 GeV. The B particle would either be stable and escape undetected, or would be produced on-shell and would subsequently decay into some other invisible states.

Let us recall that the differential rate for the decay $A \to B\gamma\gamma$ depends only on the invariant mass of the two photons, $z \equiv m_{\gamma\gamma}^2/M_A^2$, or equivalently, on the B momentum $P_B \equiv |\mathbf{p}_B|/M_A = \sqrt{\lambda}/2$, with $\lambda \equiv \lambda(1, z, r^2) = 1 + z^2 + r^4 - 2z - 2r^2 - 2zr^2$ the standard kinematical function and $r \equiv M_B/M_A$. Specifically,

$$\Gamma(A \to B\gamma\gamma) = \int_0^{(1-r)^2} dz \frac{d\Gamma}{dz}[z]$$

$$= \int_0^{(1-r^2)/2} \frac{2P_B dP_B}{\sqrt{r^2 + P_B^2}} \frac{d\Gamma}{dz}\left[z(P_B) = 1 + r^2 - 2\sqrt{P_B^2 + r^2}\right]. \quad (1)$$

To match the observed ATLAS spectrum [1], all that is needed is a differential rate falling down sufficiently fast above 750 GeV. Far below the peak, the SM background quickly increases and would wipe out any sensitivity to the $A \to B\gamma\gamma$ process. Still, slightly below the peak, at around 600 GeV, the event rate matches the background. Even if this corresponds only to a few data point, for which the uncertainty is still rather large, the differential rate should preferably fall down not too slowly as $m_{\gamma\gamma}^2$ decreases.

2. Effective four-point interactions

To check whether a peaked behavior for the diphoton invariant mass spectrum is realistic, and since the nature of the decaying state is no longer constrained, we can consider various assignments for A and B. Our basic assumption is that A and B are neutral under the SM gauge group, but nevertheless share some conserved charge χ. If $\chi(A) = -\chi(B)$, the effective interactions involving a pair of photons can derive from either

Scalar case :

$$\mathcal{L}_{eff} = \frac{1}{\Lambda^2}(S_A S_B F_{\mu\nu} F^{\mu\nu} + S_A S_B F_{\mu\nu} \tilde{F}^{\mu\nu}), \quad (2a)$$

Fermion case :

$$\mathcal{L}_{eff} = \frac{1}{\Lambda^3}(\bar{\psi}_A^C \psi_B F_{\mu\nu} F^{\mu\nu} + \bar{\psi}_A^C \gamma_5 \psi_B F_{\mu\nu} \tilde{F}^{\mu\nu} + h.c.), \quad (2b)$$

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http://dx.doi.org/10.1016/j.physletb.2016.03.068
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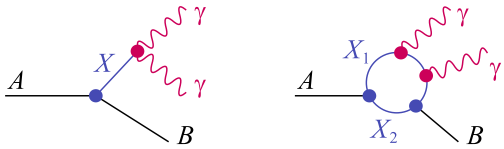

Fig. 1. Example of short-distance processes leading to the effective interactions in Eq. (2). For the tree-level diagram, X is a scalar or tensor state, whose coupling to two photons must involve yet another state. For the loop diagram, there must be a pair of states circulating the loop to ensure χ conservation.

Vector case :

$$\mathcal{L}_{eff} = \frac{1}{\Lambda^4}(A_{\alpha\beta}B^{\alpha\beta}F_{\mu\nu}F^{\mu\nu} + A_{\alpha\beta}\tilde{B}^{\alpha\beta}F_{\mu\nu}\tilde{F}^{\mu\nu} + \ldots), \tag{2c}$$

where $\tilde{F}^{\mu\nu} = \frac{1}{2}\varepsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$ and possible Wilson coefficients dressing each operator can be thought as being absorbed in the scale Λ for notation clarity. These effective operators are all independent, and assumed valid above the electroweak scale. In this respect, they should thus actually be written in terms of the $SU(2)_L$ and/or $U(1)_Y$ field strengths. For instance, replacing

$$F^{\mu\nu} \to B^{\mu\nu} = \cos\theta_W F^{\mu\nu} - \sin\theta_W Z^{\mu\nu}, \tag{3}$$

the $\gamma\gamma$, $Z\gamma$, and ZZ modes would be produced in the ratio $1 : 2\tan^2\theta_W : \tan^4\theta_W$, up to kinematical effects (in exactly the same way as for the two-body interpretation of the diphoton anomaly, see *e.g.* Ref. [13]). Finally, the CP symmetry can be enforced without loss of generality, since it is always possible to set the two photons in the adequate CP state ($F_{\mu\nu}F^{\mu\nu}$ and $F_{\mu\nu}\tilde{F}^{\mu\nu}$ have opposite parity).

At this stage, the main issue is whether simpler interactions, as for instance those involving a single photon, are possible. Though a full answer to this question would require constructing full-fledged UV completions, which is well beyond our current scope, we can nevertheless draw a number of observations. These effective interactions could either arise at tree level or at loop level, see Fig. 1, and in general require more than one extra state. For instance, in the former case, the additional resonance X would be a scalar or tensor state coupled to two photons. We only allow it to be off-shell, since otherwise the three-body signature would be lost. The X would simply be a true diphoton resonance with a mass of 750 GeV. Still, even if off-shell, this state can couple to two photons only through additional new degrees of freedom, for example a vector fermion loop. The main interest of this scenario is that the single-photon processes are automatically absent.

If generated at loop level, two new states are also required in general to ensure the conservation of χ and prevent $A, B \to \gamma\gamma$. Both of them could be fermions when A and B are scalars or vectors, but at least one new scalar or vector is needed to induce $\psi_A \to \psi_B\gamma\gamma$. The only exception is the charged scalar loop when A, B are themselves also scalars, with a renormalizable ABX^+X^- vertex. Anyway, looking at Fig. 1, it seems obvious that such loops induce also single photon modes (along with potentially large mixings between the two states, which we assumed have been dealt with properly so that states occurring in the effective interactions are true mass eigenstates). Whether such processes truly occur, and in case they do, the relative strength of the one and two photon modes, depends on the nature of A and B, so we now discuss the various assignments separately.

2.1. Scalar transitions

The single photon production $S_A \to S_B\gamma$ is forbidden by Lorentz and gauge invariance (for the same reason as, *e.g.*, $K^+ \nrightarrow \pi^+\gamma$ or $\eta \nrightarrow \pi^0\gamma$). At the renormalizable level, trivially, a direct coupling of the photon field A^μ to the scalar current $S_A\partial_\mu S_B - S_B\partial_\mu S_A$ is not gauge invariant since the current is not conserved when $m_A \neq m_B$. Beyond leading order, effective operators involving a single photon field can be constructed, for instance

$$\frac{1}{\Lambda^2}S_A\partial^\nu S_B\partial^\mu F_{\mu\nu}, \tag{4}$$

but the amplitude necessarily vanishes for an on-shell photon. There is no corresponding operator involving $\tilde{F}^{\mu\nu}$, as can be easily understood at the Feynman rule level since there are only three independent four-vectors to be contracted with $\varepsilon^{\mu\nu\rho\sigma}$. This implies that if S_A and S_B are real fields with the same parity, then $S_A \to S_B\gamma^*$ is CP-violating (as is *e.g.* $\eta \to \pi^0\ell^+\ell^-$).

Interestingly, this could suffice to reduce the $S_A \to S_B\ell^+\ell^-$ or $S_A \to S_Bq\bar{q}$ signals, even when CP conserving. Since the effective interaction is of the same dimension as the two-photon ones, producing the fermion pair through $A \to B[\gamma^* \to f\bar{f}]$ is at best comparable to the $\gamma\gamma$ mode, and could actually end up very suppressed if the situation for $K^0 \to \pi^0\gamma\gamma$ compared to $K_S \to \pi^0\ell^+\ell^-$ is of any guide [42].

Coming back to the vector fermion loop, it is easy to see that it never induces the operator Eq. (4). If both scalars couple as $S_{A,B}\bar{\psi}_F\psi_F$ or $S_{A,B}\bar{\psi}_F\gamma_5\psi_F$ with ψ_F the electrically charged heavy vector fermion circulating in the loop, then the process is CP-violating and the sum of the two diagrams where ψ_F circles clockwise and anticlockwise cancel each other. If one scalar couples through $\bar{\psi}_F\psi_F$ and the other though $\bar{\psi}_F\gamma_5\psi_F$, then both amplitudes are proportional to $\varepsilon_{\mu\nu\rho\sigma}(\varepsilon_\gamma^*)^\mu p_A^\nu p_B^\rho p_\gamma^\sigma = 0$ since $p_A = p_B + p_\gamma$. At this level, single photon processes cannot be induced.

2.2. Vector transitions

For the vector case, first remark that we do not consider all possible index contractions among the four field strengths in Eq. (2), but only some representative examples from the point of view of the differential rate. More importantly, we have not included dimension-six operators like $A_\alpha B^\alpha F_{\mu\nu}F^{\mu\nu}$ for three reasons. First, those would lead to differential rates very similar to the scalar case. Second, they may be quite complicated to generate from some UV completion. Finally, nothing would prevent a renormalizable coupling to a single photon like $A^\mu B^\nu F_{\mu\nu}$. The Landau–Yang theorem does not apply without gauge invariance or with two different vector bosons in the final state.

Even if we insist on constructing only operators involving field strengths, the $V_A \to V_B\gamma$ process is not manifestly forbidden because $m_{A,B} \neq 0$, as can be seen starting with

$$\frac{1}{\Lambda^2}(A_{\nu\alpha}B^{\alpha\mu}F_\mu^\nu + A_{\nu\alpha}B^{\alpha\mu}\tilde{F}_\mu^\nu + \ldots). \tag{5}$$

Nevertheless, the $V_A \to V_B\gamma$ process along with $V_A \to V_B[\gamma^* \to \ell^+\ell^-, q\bar{q}]$ could be very suppressed. Taking again the vector fermion loop, and assuming V_A and V_B have both either vector or axial-vector couplings to ψ_F, charge conjugation ensures the cancellation of all the diagrams to which an odd number of vector fields are attached. So, instead of the Landau–Yang theorem, what really matters in this case is the Furry theorem of QED. Note that axial-vector couplings seem more tenable to prevent the kinetic mixing $V_{A,B} \leftrightarrow \gamma$, though we have not analyzed the vector coupling scenario further.

2.3. Fermion transitions

For the fermion case, operators involving a single field strength are not forbidden. Gauge invariance prevents the direct coupling to