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Is an axizilla possible for di-photon resonance?

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ABSTRACT

Heavy axion-like particles, called axizillas, are simple extensions of the standard model (SM). An axizilla is required not to couple to the quarks, leptons, and Brout–Englert–Higgs doublets of the SM, but couple to the gauge anomalies of the $W^{\pm}$, Z and photon. It is possible to have its branching ratios (BRs) to two photons greater than 10% and to two Z's less than 10%. To have a (production cross section)·(BR to di-photons) at a 10^{-38} cm^2 level, a TeV scale heavy quark Q is required for the gluon–quark fusion process. The decay of Q to axizilla plus quark, and the subsequent decay of the axizilla to two photons can be fitted at the required level of 10^{-38} cm^2.

1. Introduction

The recent report on possible di-photon events at 750 GeV from the LHC Run-II experiments [1–4] triggered a lot of theoretical interest on this issue. The requirement to explain the rate is to require the production rate, $\sigma_{\rm production}\cdot$(branching ratio (BR) to di-photons), of order 10^{-38} cm^2 with the LHC parton distribution at 13 TeV energy. Any model for a diphoton resonance of mass 750 GeV decaying to two photons is better not to give such di-photons at the previous LHC Run-I at 8 TeV.

This invites to search for "Which particle is most economically introduced beyond the standard model (SM)?" A phenomenology on this is summarized in Ref. [4] and the recent papers are listed in [5,6]. Here, we search for a theory motivated particle. The well-known examples are axions [7], majorons [8], ALPs [9], and quintessential pseudoscalars or ultralight axions [10], which are much lighter than electron. We argue that not only these very light pseudoscalars but also TeV scale pseudoscalars are theoretically prospective. Since pseudoscalars are pseudo-Goldstone bosons of some spontaneously broken axial symmetry, we call these heavy pseudoscalars *axizillas*. In this paper, we investigate a possibility of an axizilla for the di-photon resonance of the LHC Run II.

It is known that string theory does not allow any global symmetry below the compactification scale, except the model-independent axion [11,12]. Also from the topological obstruction of global symmetries from gravity [13], it has been argued that a serious fine-tuning problem is present in axion physics [14]. In string compactification with the anomalous U(1) gauge symmetry, a global Peccei–Quinn (PQ) symmetry can survive down to an intermediate scale [15]. The resulting invisible axion is from a phase field of matter scalars instead from the anti-symmetric tensor field $B_{\mu\nu}$ [16]. Except from the anomalous U(1) gauge symmetry, any global symmetry must be approximate. For example, it has been shown numerically that there exist compactification models suppressing the explicit PQ symmetry breaking terms at some level [17]. Generically, however, any global symmetry in consideration must be broken at some leading scale below the compactification scale. Nevertheless some discrete symmetries can be allowed without the gravity obstruction problem [18].

In Fig. 1, we show the superpotential terms allowed by the discrete symmetry in the most left vertical column. If we consider a few lowest order terms, *i.e.* those inside the lavender square, there might appear a global symmetry. The terms allowed by this global symmetry are shown in the horizontal bar including those in the green box. Thus, this global symmetry is broken by the terms in the most left red boxes in the vertical column. In addition, we have shown also the non-Abelian anomaly terms which also break the global symmetry. If the PQ type global symmetry is respected by the superpotential, we neglect the most left column. In this case, the θ angle of the non-Abelian group vacuum chooses $\theta=0$ as

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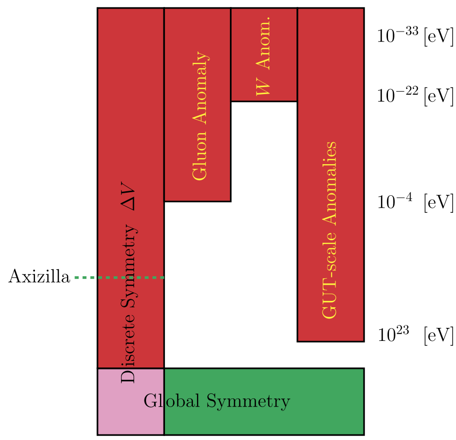

Fig. 1. A cartoon of classifying symmetries at low energy. The terms in the vertical column are allowed by discrete symmetries in string compactification. The terms in the superpotential belong to the most left column. The other three columns show anomalous terms. If one considers a few leading effective terms, *i.e.* corresponding to the lavender square, the terms there define an effective global symmetry. However, this global symmetry is broken by the terms in the red boxes. Some mass scales needed in cosmology are also shown with the axion mass scales [19–22]. (For interpretation of the references to color in this figure, the reader is referred to the web version of this article.)

the minimum, which is used in extremely light axions, quintessential axion [19], ultra light axion [20], QCD axion [21], and axionic inflation [22]. The respective axion mass scales are shown at the far right. On the left-hand side (LHS), the mass scale of axizilla is shown. The breaking scale is not known in any known non-Abelian gauge symmetry, and its mass derives from the global symmetry breaking potential, ΔV. Let the global symmetry be $\mathrm{U}(1)_\Sigma$. Suppose S carries the discrete quantum number but is neutral under $\mathrm{U}(1)_\Sigma$ and σ carries both the discrete and $\mathrm{U}(1)_\Sigma$ quantum numbers,

$$S,\quad \text{with a GUT scale VEV } M_{\rm GUT},\ \text{but not breaking } \mathrm{U}(1)_{\rm global},$$
$$\sigma,\quad \text{with a VEV } f/\sqrt{2},\ \text{breaking } \mathrm{U}(1)_{\rm global}. \tag{1}$$

For some discrete symmetry $\mathbf{Z}_N$, assign the global quantum numbers Σ as those of $\mathbf{Z}_N$ with $-N<\Sigma<N$. For a discrete symmetry $\mathbf{Z}_4$ for example, let the discrete quantum numbers of σ be 1. Then, the following global symmetry breaking term, allowed by $\mathbf{Z}_4$, as shown in Table 1 is present

$$\Delta V\sim\sigma^4+\text{h.c.}\to f^4\cos\left(4\frac{P}{f}\right)$$

where $|\sigma|=f/\sqrt{2}$. Without S fields, if f is of order TeV scale, which can be determined by the $\mathrm{U}(1)_\Sigma$ preserving terms, then the pseudoscalar mass m_P is at the TeV scale. This is an axizilla. With S fields included, more complicated discrete symmetries can produce TeV scale axizillas. A cosmological effect of $\mathbf{Z}_N$ symmetry is the appearance of domain walls [23] which however is difficult to be observed after inflation.

2. Axizillas with discrete symmetry $\mathbf{Z}_N$

With a discrete symmetry, we try to introduce a global anomaly but without a color anomaly. If it has a color anomaly, it is necessarily related to the QCD axion and the symmetry breaking scale must be larger than 10^{10} GeV. As the simplest example, let us introduce a vectorlike doublet with the hypercharge $Y=-\frac12$,

$$\ell_L=\begin{pmatrix}N\\L\end{pmatrix}_L,\quad \ell_R=\begin{pmatrix}N\\L\end{pmatrix}_R. \tag{2}$$

Table 1
The $\mathrm{U}(1)_\Sigma$ quantum numbers. ΔV can contain $\sigma^4(S^*+S^2)$.

	ℓ_L	ℓ_R	σ	S	E_L	E_R
$\mathbf{Z}_{12}$	$-\frac12$	$+\frac12$	1	+4	$+\frac12$	$-\frac12$
Σ	$-\frac12$	$+\frac12$	1	+4	$+\frac12$	$-\frac12$

In the SM, the pseudoscalar (P) coupling to color singlet gauge bosons are

$$\mathcal{L}=\frac{P}{f}\frac{g_2^2}{32\pi^2}W^a_{\mu\nu}\tilde{W}^{b\,\mu\nu}\,(\mathrm{Tr}\ T_aT_b)+\frac{P}{f}\frac{g'^2}{32\pi^2}Y_{\mu\nu}\tilde{Y}^{\mu\nu}\left(\mathrm{Tr}\,Y_L^2+\mathrm{Tr}\,Y_R^2\right) \tag{3}$$

where $W^a_{\mu\nu}$ is the non-Abelian field strength of SU(2) gauge fields A^a_μ and $Y_{\mu\nu}$ is the field strength of $\mathrm{U}(1)'$ gauge fields Y_μ. For a vectorlike fundamental representation in SU(N), like a heavy quark axion model [21] or Eq. (2), $\mathrm{Tr}\,T_aT_b=\frac12\delta_{ab}$. Thus, Eq. (3) becomes

$$\begin{aligned}\mathcal{L}&=\frac{P}{f}\frac{g_2^2}{32\pi^2}W^a_{\mu\nu}\tilde{W}^{a\,\mu\nu}+\frac{P}{f}\frac{g'^2}{32\pi^2}Y_{\mu\nu}\tilde{Y}^{\mu\nu}\\&=\frac{P}{32\pi^2 f}\Big(2g_2^2W^+_{\mu\nu}\tilde{W}^{-\,\mu\nu}+2e^2F^{\rm em}_{\mu\nu}\tilde{F}^{{\rm em}\,\mu\nu}\\&\quad+g_2^2(1/c_W^2-2s_W^2)Z_{\mu\nu}\tilde{Z}^{\mu\nu}+2\frac{eg_2}{c_W}F^{\rm em}_{\mu\nu}\tilde{Z}^{\mu\nu}\Big)\end{aligned} \tag{4}$$

where

$$\begin{aligned}W^3_\mu&=\cos\theta_W Z_\mu+\sin\theta_W A_\mu,\\Y_\mu&=-\sin\theta_W Z_\mu+\cos\theta_W A_\mu,\\c_W&=\cos\theta_W=\frac{g_2}{\sqrt{g_2^2+g'^2}},\\s_W&=\sin\theta_W=\frac{g'}{\sqrt{g_2^2+g'^2}}.\end{aligned} \tag{5}$$

The massless combination to the photon coupling is parametrized by $c_{P\gamma\gamma}$,

$$\mathcal{L}_{P\gamma\gamma}=\frac{c_{P\gamma\gamma}P}{f}\frac{e^2}{32\pi^2}F^{\rm em}_{\mu\nu}\tilde{F}^{{\rm em}\,\mu\nu} \tag{6}$$

where $c_{P\gamma\gamma}$ turns out to be 2 (see Table 1). Treating the W and Z bosons as massless, we estimate the branching ratios (BRs), to the decay modes to W^+W^-, 2γ, $2Z$, and $Z\gamma$ of Eq. (2), for which the BRs are shown in the first row of Table 2, where we used $\sin^2\theta_W\simeq 0.23$. The effect of $\mathbf{Z}_N$ symmetry is to constrain possible interactions such that the leading term is suppressed by one power of f. Without the $\mathbf{Z}_N$ symmetry, some terms in the potential dominate this anomaly term [14].

Let us now introduce n_1 pairs of $Q_{\rm em}=-1$ vectorlike SU(2) singlet E and n_2 pairs of (2),

$$n_2\left\{\ell_L=\begin{pmatrix}N\\L\end{pmatrix}_L,\ \ell_R=\begin{pmatrix}N\\L\end{pmatrix}_R\right\},\quad n_1\left\{E_L,\ E_R\right\}. \tag{7}$$

Now, we have

$$\begin{aligned}\mathcal{L}_{P\text{-decay}}&=n_2\frac{P}{f}\frac{g_2^2}{32\pi^2}W^a_{\mu\nu}\tilde{W}^{a\,\mu\nu}-(2n_1-n_2)\frac{P}{f}\frac{g'^2}{32\pi^2}Y_{\mu\nu}\tilde{Y}^{\mu\nu}\\&=\frac{P}{32\pi^2 f}\Big(g_2^2\,2n_2W^+_{\mu\nu}\tilde{W}^{-\,\mu\nu}\end{aligned}$$