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Low-Scale Axion from Large Extra Dimensions

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Abstract

The mass of the axion and its decay rate are known to depend only on the scale of Peccei-Quinn symmetry breaking, which is constrained by astrophysics and cosmology to be between 10^9 and 10^{12} GeV. We propose a new mechanism such that this effective scale is preserved and yet the fundamental breaking scale of $U(1)_{PQ}$ is very small (a kind of inverse seesaw) in the context of large extra dimensions with an anomalous U(1) gauge symmetry in our brane. Unlike any other (invisible) axion model, there are now possible collider signatures in this scenario.

Although CP violation has been observed in weak interactions [1, 2] and it is required for an explanation of the baryon asymmetry of the universe [3], it becomes a problem in strong interactions. The reason is that the multiple vacua of quantum chromodynamics (QCD) connected by instantons [4] require the existence of the CP violating θ term [5]

$$\mathcal{L}_\theta = \theta_{QCD} \frac{g_s^2}{32\pi^2} G^a_{\mu\nu} \widetilde{G}^{a\mu\nu}, \tag{1}$$

where g_s is the strong coupling constant, $G^a_{\mu\nu}$ is the gluonic field strength and $\tilde{G}^a_{\mu\nu}$ is its dual. Nonobservation of the electric dipole moment of the neutron [6] implies that

$$\bar{\theta} = \theta_{QCD} - Arg\ Det\ M_u\ M_d < 10^{-10}, \tag{2}$$

instead of the theoretically expected order of unity. In the above, M_u and M_d are the respective mass matrices of the charge $2/3$ and $-1/3$ quarks of the standard model of particle interactions. This is commonly known as the strong CP problem.

The first and best motivated solution to the strong CP problem was proposed by Peccei and Quinn [7], in which the quarks acquire a dynamical phase from the spontaneous breaking of a new global symmetry $[U(1)_{PQ}]$ and relaxes $\bar{\theta}$ to its natural minimun value of zero. As a result, there appears a Goldstone boson called the axion but it is not strictly massless [8] because it couples to two gluons (like the neutral pion) through the axial triangle anomaly [9].

The scale of $U(1)_{PQ}$ breaking (which is conventionally identified with the axion decay constant f_a) determines the axion coupling to gluons, which is proportional to $1/f_a$. If f_a is the electroweak symmetry breaking scale as originally proposed [7], then the model is already ruled out by laboratory experiments [10]. In fact, f_a is now known to be constrained by astrophysical and cosmological arguments [11] to be between 10^9 and 10^{12} GeV. Hence the axion must be an electroweak singlet or predominantly so. It may couple to the usual quarks and leptons through a suppressed mixing with the standard Higgs doublet [12], or it may couple only to other unknown colored fermions [13], or it may couple to gluinos [14] as well as all other supersymmetric particles.

Because the axion must necessarily mix with the π and η mesons, it must have a two-photon decay mode. This is the basis of all experimental attempts [10] to discover its existence. On the other hand, the accompanying new particles in all viable axion models to date are very heavy, i.e. of order f_a; hence they are completely inaccessible to experimental verification.

In the following we consider instead the possiblily that the $U(1)_{PQ}$ breaking scale is actually very small, but that f_a is large because of a kind of inverse seesaw mechanism. We show how this scenario may be realized in the context of large extra dimensions with an anomalous U(1) gauge symmetry in our brane. The associated new physics now exists at around 1 TeV, with a number of interesting observable consequences at future colliders.

We assume a singlet scalar field χ with a nonzero PQ charge existing in the bulk of large extra dimensions [15]. The *shining* [16] of this field in our brane is the source of spontaneous $U(1)_{PQ}$ breaking in our world (called a 3-brane). The idea is that χ gets a large vacuum expectation value (VEV) in a distant brane, but its effect on our brane is small because we are far away from it. (In the case of lepton number, this mechanism has been used recently to obtain small Majorana neutrino masses [17].) To convert this small $\langle\chi\rangle$ to a large f_a, we need to assume an anomalous U(1) gauge symmetry in our brane at the TeV energy scale, as explained below.

In a theory of large extra dimensions with quantum gravity at the TeV scale, there is no large scale available for the axion. Since the behavior of Goldstone bosons depends not on the coupling but only on the scale of symmetry breaking in general, it is a problem which is not easily resolved [18]. Here we find a new and novel solution to this apparent contradiction in the case where there is an anomalous U(1) gauge symmetry, which is of course well studied

[19] as a possible manifestation of string theory near the string scale (now considered also at around a few TeV) and has well-known applications in quark and lepton Yukawa textures and supersymmetry breaking.

We extend the standard model of particle interactions to include an extra $U(1)_A$ gauge symmetry and an extra $U(1)_{PQ}$ global symmetry. All standard-model particles are trivial under these two new symmetries. We then introduce a new heavy quark singlet ψ and two scalar singlets σ and η with $U(1)_A$ and $U(1)_{PQ}$ charges as shown in Table 1. All fields except χ are confined to our brane.

Table 1: Peccei-Quinn charges of the fermions and scalars

Fields	Transformation under		
	$SU(3)_C \times SU(2)_L \times U(1)_Y$	$U(1)_A$	$U(1)_{PQ}$
$(u_i, d_i)_L$	(3,2,1/6)	0	0
u_{iR}	(3,1,2/3)	0	0
d_{iR}	(3,1,−1/3)	0	0
$(\nu_i, e_i)_L$	(1,2,−1/2)	0	0
e_{iR}	(1,1,−1)	0	0
ψ_L	(3,1,–1/3)	1	k
ψ_R	(3,1,–1/3)	–1	$-k$
(ϕ^+, ϕ^0)	(1,2,1/2)	0	0
σ	(1,1,0)	2	$2k$
η	(1,1,0)	2	$2k-2$
χ	(1,1,0)	0	2

Because of our chosen charge assignments, only the field σ couples to the colored fermion

ψ, i.e.

$$\mathcal{L}_Y = f\sigma\bar{\psi}_L\psi_R + h.c. \tag{3}$$

Hence it also couples to two gluons through the usual triangular loop. As σ acquires a VEV, say u, of order 1 TeV, both $U(1)_A$ and $U(1)_{PQ}$ are broken, whereas the latter is broken by $\langle\chi\rangle = z$, and it induces a VEV also for η, i.e. $\langle\eta\rangle = w$. We will show in the following that given z is small from its origin in the bulk, w is also small. Now the longitudinal component of the Z_A boson is mostly given by Imσ, so the axion is excluded to be mostly a linear combination of Imη and Imχ, but the latter two fields do not couple to the colored fermion ψ. As a result, the axion's coupling to two gluons is now effectively

$$\frac{1}{f_a} = \frac{w^2}{u^2\sqrt{w^2+z^2}}, \tag{4}$$

which can be thought of as a kind of inverse seesaw, i.e. the largeness of f_a is explained by the smallness of w. Details will be given later.

Our brane $\mathcal{P}$ is located at a point $y = 0$ in the bulk. Peccei-Quinn symmetry is broken maximally in a distant brane $\mathcal{P}'$, located at a point $y = y_*$ in the bulk. We assume for simplicity that the separation of the two branes is of order the radius of compactification of the extra space dimensions, i.e. $|y_*| = r$, which is only a few μm in magnitude. The fundamental scale M_* in this theory is then related to the reduced Planck scale $M_P = 2.4\times 10^{18}$ GeV by the relation

$$r^n M_*^{n+2} \sim M_P^2. \tag{5}$$

The $U(1)_{PQ}$ symmetry breaking in the distant brane acts as a point source J, which induces an effective VEV, i.e. z, to the singlet bulk field χ. Other effects which may perturb the *shined* value of $\langle\chi\rangle$ in our world are all included as boundary conditions to the source J, so that the effect of the field χ in our brane always appears in the combination $z(y=0)e^{i\varphi}$, where $\varphi(x)$ is a dynamical phase which transforms under $U(1)_{PQ}$ to preserve its invariance.

This formulation has also been used for the spontaneous breaking of lepton number in the case of neutrinos [17].

In our brane, the profile of χ is given by the Yukawa potential in the transverse dimensions

$$\langle\chi(y=0)\rangle = J(y=y_*)\Delta_n(r), \tag{6}$$

where

$$\Delta_n(r) = \frac{1}{(2\pi)^{\frac{n}{2}} M_*^{n-3}} \left(\frac{m_\chi}{r}\right)^{\frac{n-2}{2}} K_{\frac{n-2}{2}}\left(m_\chi r\right), \tag{7}$$

K being the modified Bessel function. We consider the source to be dimensionless, which we take to be $J=1$. For the interesting case of $n>2$ and $m_\chi r \ll 1$, the *shined* value of χ is given by

$$\langle\chi\rangle \approx \frac{\Gamma(\frac{n-2}{2})}{4\pi^{\frac{n}{2}}} \frac{M_*}{(M_* r)^{n-2}} = \frac{\Gamma(\frac{n-2}{2})}{4\pi^{\frac{n}{2}}} M_* \left(\frac{M_*}{M_P}\right)^{2-(4/n)}. \tag{8}$$

For $n=3$ and $M_*=10$ TeV, we get $\langle\chi\rangle \sim 0.2$ keV. This is the smallest value possible with our assumptions. However, if the distant brane is located at y_* less than r, larger values of $\langle\chi\rangle$ may be obtained. As we will show, the range 1 keV to 1 MeV corresponds nicely to the axion decay constant of 10^{12} to 10^9 GeV.

We express the bulk field as

$$\chi = \frac{1}{\sqrt{2}}(\rho + z\sqrt{2})e^{i\varphi}. \tag{9}$$

Its self-interaction terms are now given by

$$V(\chi) = \lambda z(y)^2\rho(x,y)^2 + \frac{1}{\sqrt{2}}\lambda z(y)\rho(x,y)^3 + \frac{1}{8}\lambda\rho(x,y)^4. \tag{10}$$

This Lagrangian has the virtue of universality, i.e., λ is unchanged, but z can change depending on where our brane is from the distant brane. The invariance under $U(1)_{PQ}$, i.e. $\rho \rightarrow \rho$ and $\varphi \rightarrow \varphi + 2\theta$, is also maintained in the other interactions, as described below. The parameters in the potential of χ are thus guaranteed to be independent of the parameters of our brane.

The scalar potential in our brane excluding $V(\chi)$ is now given by

$$\begin{aligned} V &= m_1^2\Phi^\dagger\Phi + m_2^2\sigma^\dagger\sigma + m_3^2\eta^\dagger\eta + \frac{1}{2}\lambda_1(\Phi^\dagger\Phi)^2 + \frac{1}{2}\lambda_2(\sigma^\dagger\sigma)^2 + \frac{1}{2}\lambda_3(\eta^\dagger\eta)^2 \\ &+\lambda_4(\Phi^\dagger\Phi)(\sigma^\dagger\sigma) + \lambda_5(\Phi^\dagger\Phi)(\eta^\dagger\eta) + \lambda_6(\sigma^\dagger\sigma)(\eta^\dagger\eta) + (\mu z e^{i\varphi}\sigma^\dagger\eta + h.c.), \end{aligned} \tag{11}$$

where μ has the dimension of mass and we assume that all mass parameters are of the same order of magnitude, i.e. 1 TeV.

The minimum of V satisfies the following conditions:

$$m_1^2 + \lambda_1 v^2 + \lambda_4 u^2 + \lambda_5 w^2 = 0, \tag{12}$$

$$u(m_2^2 + \lambda_2 u^2 + \lambda_4 v^2 + \lambda_6 w^2) + \mu z w = 0, \tag{13}$$

$$w(m_3^2 + \lambda_3 w^2 + \lambda_5 v^2 + \lambda_6 u^2) + \mu z u = 0, \tag{14}$$

where $\langle\phi^0\rangle = v$. Hence

$$v^2 \simeq \frac{-\lambda_2 m_1^2 + \lambda_4 m_2^2}{\lambda_1\lambda_2 - \lambda_4^2}, \tag{15}$$

$$u^2 \simeq \frac{-\lambda_1 m_2^2 + \lambda_4 m_1^2}{\lambda_1\lambda_2 - \lambda_4^2}, \tag{16}$$

and

$$w \simeq \frac{-\mu z u}{m_3^2 + \lambda_5 v^2 + \lambda_6 u^2}, \tag{17}$$

which is indeed of order z as mentioned earlier.

Whereas $\mathrm{Im}\phi^0$ becomes the longitudinal component of the usual Z boson, $(u\mathrm{Im}\sigma + w\mathrm{Im}\eta)/\sqrt{u^2+w^2}$ becomes that of the new Z_A boson. Since the 3×3 mass matrix in the basis $[\mathrm{Im}\sigma,\ \mathrm{Im}\eta,\ z\varphi]$ is given by

$$\begin{pmatrix} -\mu z w/u & \mu z & \mu w \\ \mu z & -\mu z u/w & -\mu u \\ \mu w & -\mu u & -\mu u w/z \end{pmatrix}, \tag{18}$$

the axion a is identified as the following:

$$\begin{aligned}\frac{a}{\sqrt{2}} &= \frac{1}{N}\left[uw^2\mathrm{Im}\sigma - wu^2\mathrm{Im}\eta + z(u^2+w^2)z\varphi\right] \\ &\simeq \frac{w^2}{u\,(w^2+z^2)^{1/2}}\mathrm{Im}\sigma - \frac{w}{(w^2+z^2)^{1/2}}\mathrm{Im}\eta + \frac{z}{(w^2+z^2)^{1/2}}z\varphi, \end{aligned} \tag{19}$$

where $N = \{w^2u^2(w^2+u^2) + z^2(w^2+u^2)^2\}^{1/2}$ is the normalization. Since only σ couples to the colored fermion ψ and the component of $\mathrm{Im}\sigma$ in the axion is u times a phase, the axion coupling to the gluons through ψ is effectively as given by Eq. (4) as mentioned earlier. Using $u \sim 1$ TeV and $w \sim z \sim 1$ keV to 1 MeV, we see that f_a is indeed in the range 10^{12} to 10^9 GeV.

In Table 1, we have not specified the value of k for the PQ charge of ψ. This is intentional because our model is independent of it. This ambiguity also helps us to understand its pattern of symmetry breaking. For example, if $\langle\chi\rangle = 0$, then $\langle\eta\rangle = 0$ also. In that case, there is no axion and the Peccei-Quinn symmetry disappears, i.e. $k = 0$. Hence the true scale of $U(1)_{PQ}$ breaking is indeed small, i.e. z from the bulk, as asserted.

To understand why we have an exception to the general rule that the axion coupling is inversely proportional to the scale of $U(1)_{PQ}$ breaking, we point out that the anomalous nature of $U(1)_A$ is crucial. If we attempt to make it free of the axial triangle anomaly, we need to add colored fermions with opposite $U(1)_A$ charges to $\psi_{L,R}$. They must then acquire mass through a new scalar field with opposite $U(1)_A$ charge to σ. The longitudinal component of Z_A takes up a linear combination of the two imaginary parts, leaving free the other to be the axion, which now couples to the colored fermions with the same scale as $U(1)_A$ symmetry breaking. The above is of course the analog of what happens in the well-known original Peccei-Quinn proposal [7].

All axion models to date have no accompanying verifiable new physics other than the $a \to \gamma\gamma$ decay, and that depends on the axion being a component of dark matter. In our scenario, the possibility exists for this new physics to be at the TeV scale and be observable at future colliders.

(1) The stable heavy colored fermion ψ may be produced in pairs, i.e. $gg \to \psi\bar{\psi}$. Both ψ and $\bar{\psi}$ carry light quarks and gluons with them and appear as jets, but when these jets hit the hadron calorimeter in a typical detector, a large part (i.e. $2m_\psi$) of the initial collision energy is "frozen" in the mass and appears "lost".

(2) There is mixing between the standard-model Higgs boson Reϕ^0 with the new scalar Reσ of order v/u, i.e. 0.1 or so. This means that the lighter (call it h) of the two physical scalar bosons has a small component of Reσ, but that only modifies its (small) gg and $\gamma\gamma$ decay amplitudes through the ψ loop. Hence h behaves almost exactly like the standard-model Higgs boson.

(3) The $U(1)_A$ gauge boson Z_A may be produced by gg^* fusion through the ψ loop. If kinematically allowed, it will decay into Reη + Imη. Since Imη is partly $(w/\sqrt{w^2+z^2})$ the axion a which will escape detection, this event has a lot of possible missing transverse momentum. The subsequent decay of Reη is into a and a virtual Z_A which turns into gg. This adds more missing transverse momentum. The end result of the production and subsequent decay of Z_A is thus two gluon jets and two axions. This is a distinctive signature of our scenario [20]. It predicts collider events with large missing energy without the existence of supersymmetry.

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References

[1] J. H. Christenson *et al.*, Phys. Rev. Lett. **13**, 138 (1964).

[2] A. Alavi-Harat *et al.*, Phys. Rev. Lett. **83**, 22 (1999).

[3] A. Sakharov, Pis'ma Zh. Eksp. Teor. Fiz. **5**, 32 (1967).

[4] G. 't Hooft, Phys. Rev. Lett. **37**, 8 (1976).

[5] C. G. Callan, R. Dashen, and D. J. Gross, Phys. Lett. **63**, 334 (1976); R. Jackiw and C. Rebbi, Phys. Rev. Lett. **37**, 172 (1976).

[6] P. G. Harris *et al.*, Phys. Rev. Lett. **82**, 904 (1999).

[7] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. **38**, 1440 (1977).

[8] S. Weinberg, Phys. Rev. Lett. **40**, 223 (1978); F. Wilczek, Phys. Rev. Lett. **40**, 279 (1978).

[9] S. L. Adler, Phys. Rev. **177**, 2426 (1969); J. S. Bell and R. Jackiw, Nuovo Cim. **60A**, 47 (1969); S. L. Adler and W. A. Bardeen, Phys. Rev. **182**, 1517 (1969).

[10] For a recent review, see L. J. Rosenberg and K. A. van Bibber, Phys. Rep. **325**, 1 (2000).

[11] For a recent review, see G. G. Raffelt, Ann. Rev. Nucl. Part. Sci. **49**, 163 (1999).

[12] A. P. Zhitnitskii, Sov. J. Nucl. Phys. **31**, 260 (1980); M. Dine, W. Fischler, and M. Srednicki, Phys. Lett. **B104**, 199 (1981).

[13] J. E. Kim, Phys. Rev. Lett. **43**, 103 (1979); M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Nucl. Phys. **B166**, 493 (1980).

[14] D. A. Demir and E. Ma, hep-ph/0004148; D. A. Demir, E. Ma, and U. Sarkar, hep-ph/0005288.

[15] N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. **B429**, 263 (1998); I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Lett. **B436**, 257 (1998); N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, Phys. Rev. **D59**, 086004 (1999).

[16] N. Arkani-Hamed and S. Dimopoulos, hep-ph/9811353; N. Arkani-Hamed, L. Hall, D. Smith, and N. Weiner, Phys. Rev. **D61**, 116003 (2000); Y. Sakamura, hep-ph/9912511.

[17] E. Ma, M. Raidal, and U. Sarkar, hep-ph/0006046; E. Ma, G. Rajasekaran, and U. Sarkar, hep-ph/0006340.

[18] Recent axion proposals in the context of extra dimensions include S. Chang, S. Tazawa and M. Yamaguchi, Phys. Rev. **D 61**, 084005 (2000); K. R. Dienes, E. Dudas, and T. Gherghetta, hep-ph/9912455; L. Di Lella, A. Pilaftsis, G. Raffelt, and K. Zioutas, hep-ph/0006327.

[19] For a recent review, see H.-P. Nilles, hep-ph/0003102.

[20] Reη will also decay (through its mixing with Reϕ^0) into standard-model final states such as ZZ, WW, etc. However, this mixing is very small, i.e. of order $wv/\mu u$. Details of the phenomenology of this model will be discussed in a forthcoming paper.