Strong CP problem and mirror world: the Weinberg-Wilczek axion revisited

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Abstract

A new possibility for solving the strong CP-problem is suggested. It is based on the concept of a mirror world of particles, with the gauge symmetry and Lagrangian completely identical to that of the observable particles. We assume that the ordinary and mirror sectors share the same Peccei-Quinn symmetry realized \textit{à la} Weinberg-Wilczek, so that the $\theta$-terms are simultaneously canceled by the axion VEV in both worlds. This property remains valid even if the symmetry between two sectors is spontaneously broken and the the weak scale of the mirror world is larger than the ordinary weak scale, in which case also the mirror QCD scale becomes larger than the ordinary one. In this situation our axion essentially represents a Weinberg-Wilczek axion of the mirror world with quite a large mass, while it couples the ordinary particles like an invisible axion. The experimental and astrophysical limits are discussed and an allowed parameter window is found with the Peccei-Quinn scale $f_a \sim 10^4 - 10^5$ GeV and the axion mass $m_a \sim 1$ MeV, which can be accessible for future experiments. We also show that our solution to the strong CP-problem is stable against the Planck scale induced effects.
1 Introduction

The strong CP problem remains one of the puzzling points of the modern particle physics (for a review, see e.g. ref. [1]). It is related to the P- and CP-violating term \( \theta G \bar{G} \) in the QCD Lagrangian which is contributed also by the complex phases in the quark mass matrices \( M_U, D \), so that its effective value becomes \( \theta = \theta + \arg(\det M_U \det M_D) \). The CP-violating phenomena in weak interactions indicate that the CP-violating phase in the Cabibbo-Kobayashi-Maskawa (CKM) matrix is large, and hence the quark mass matrices should contain large complex phases. Therefore, one would generally expect that \( \theta \) is order 1. However, the experimental bound on the neutron electric dipole moment yields \( \theta < 10^{-10} \). The strong CP problem simply questions why the effective \( \theta \)-term is so small, or in other words, what is the origin of such a strong fine tuning between the initial value \( \theta \) and the phases of the quark mass matrices?

The most appealing solution to the problem is provided by the Peccei-Quinn (PQ) mechanism [2], based on the concept of spontaneously broken global axial symmetry \( U(1)_{\text{PQ}} \). In the context of this mechanism \( \theta \) essentially becomes a dynamical field, \( \theta = a/f_a \), with an effective potential induced by the non-perturbative QCD effects which fixes vacuum expectation value (VEV) at zero. Here \( a \) stands for the pseudo-Goldstone mode of the spontaneously broken PQ symmetry, an axion, and \( f_a \) is a VEV of scalar (or a VEV combination of several scalars) responsible for the \( U(1)_{\text{PQ}} \) symmetry breaking. This scale is model dependent. In particular, in the original Weinberg-Wilczek (WW) model \( f_a \) is order electroweak scale, but in the invisible axion models [3, 4, 5] it should be much larger.

In the most general context, the axion couplings to fermions and photons are described by the following Lagrangian terms:

\[
L_a = ig_{ak}a\bar{\psi}_k\gamma_5\psi_k + \frac{g_{a\gamma}}{4}aF_{\mu\nu}\tilde{F}^{\mu\nu} + \ldots \tag{1}
\]

where the coupling constants are essentially determined by the PQ scale \( f_a \). Typically, up to model dependent order one coefficients, we have \( g_{ak} \sim f_a^{-1}m_k \), where \( m_k \) is the fermion mass (e.g. \( m_k = m_e, m_N, \ldots \) for electron, nucleons, etc.), and \( g_{a\gamma} \sim (\alpha/2\pi)f_a^{-1} \), with \( \alpha \) being the fine structure constant.

As for the axion mass, it also depends on the QCD scale \( \Lambda \) where strong interactions become nonperturbative. In the theory with no light quarks (with masses below \( \Lambda \)), the axion mass would emerge due to instanton-induced potential which can be computed in the dilute gas approximation, \( P(a) \approx -K\cos(a/f_a) \) with \( K \sim \Lambda^4 \), and so we would have \( m_a^2 \sim \Lambda^4/f_a^2 \). However, in reality the axion state \( a \) gets small mixing with the pseudo-Goldstone bosons \( \pi, \eta, \) etc., related to the chiral symmetry breaking by the light quark condensates \( V = \langle \bar{q}q \rangle \sim \Lambda^3 \). In order to find the true mass eigenvalue corresponding to the axion, the total mass matrix should be diagonalized. In doing so, one finds a general expression [6]:

\[
m_a^2 = \frac{N^2}{f_a^2} \frac{VK}{V + K\text{Tr}M^{-1}} \tag{2}
\]

where \( N \) stands for the color anomaly of \( U(1)_{\text{PQ}} \) current,\(^4\) and \( M = \text{diag}(m_u, m_d, \ldots) \) is a mass matrix of light quarks, with \( m_q < \Lambda \). Thus, the axion mass is given as \( m_a^2 \sim m_q\Lambda^3/f_a^2 \). More precisely, neglecting the \( s \)-quark contribution and using the relation \( (m_u + m_d)\langle \bar{q}q \rangle = m_{\pi}^2f_{\pi}^2 \),

\(^4\) The PQ charges are normalized so that each of the standard fermion families contributes as \( N = 1 \). Hence, in the WW model we have \( N = N_g \), where \( N_g (= 3) \) is the number of fermion families. The same holds in the Dine-Fischler-Srednicki-Zhitnitskii (DFSZ) model [7]. Other models of the invisible axion, e.g. the hadronic axion [8] or archion [9], generally contain some exotic fermions and so \( N \neq N_g \).
from (2) one directly arrives to the more familiar formula:

\[ m_a = \frac{N}{f_a} \left( \frac{m_u m_d V}{m_u + m_d} \right)^{1/2} = \frac{N z^{1/2}}{1 + z} \cdot \frac{f_a m_\pi}{f_a} \approx \left( \frac{10^6 \text{ GeV}}{f_a/N} \right) \times 6.2 \text{ eV} \]  

(3)

where \( z = m_u / m_d \approx 0.57 \), and \( m_\pi, f_\pi \) respectively are the pion mass and decay constant.

In the WW model [6] the PQ symmetry is broken by two Higgs doublets \( H_{1,2} \) with the VEVs \( v_{1,2} \), namely \( f_a = (v/2) \sin 2\beta \), where \( v = (v_1^2 + v_2^2)^{1/2} \approx 247 \) GeV is the electroweak scale and \( \tan \beta = v_2 / v_1 \). Therefore, the WW axion is quite heavy, and its mass can vary from few hundred keV to several MeVs:

\[ m_{a\,WW} = \frac{2N}{v \sin 2\beta} \left( \frac{m_u V}{1 + z} \right)^{1/2} \approx \frac{150 \text{ keV}}{\sin 2\beta}, \]  

(4)

However, its couplings to fermions and photon are too strong and by this reason the WW model is completely ruled out for any values of the parameter \( \beta \) by a variety of the terrestrial experiments as are the search of the decay \( K^+ \to \pi^+ a, J/\psi \) and \( \Upsilon \) decays into \( a + \gamma \), nuclear deexcitations via axion emission, the reactor and beam dump experiments, etc. [6].

For a generic invisible axion, these experimental data imply a lower bound \( f_a > 10^4 \) GeV or so, which leads to the upper limit on \( m_a \) of about 1 keV. However, for an axion being so light, the stringent bounds emerge from astrophysics and cosmology. In particular, for the DFSZ model the combined limits from the stellar evolution and the supernova neutrino signal exclude all scales \( f_a \) up to \( 10^{10} \) GeV [1, 9], which in turn implies \( m_a < 10^{-4} \) eV [1]. On the other hand, the cosmological limits related to the primordial oscillations of the axion field or to the non-thermal axion production by cosmic strings, demand the upper bound \( f_a < 10^{10-11} \) GeV [3, 9]. Thus, not much parameter space remains available.

In addition, the invisible axion models have a naturality problem related to the Planck scale induced terms. Many arguments suggest that the non-perturbative quantum gravity effects do not respect the global symmetries [10] and thus they can induce the higher order effective operators cutoff by \( M_P \) which explicitly violate the PQ symmetry. For large \( f_a \) these explicit terms can be very dangerous for the stability of the PQ mechanism. For example, in the DFSZ model where \( U(1)_{\text{PQ}} \) is broken by a gauge singlet scalar \( S \), the operator \( \frac{\phi^n}{M_P^6} + \text{h.c.} \) would induce \( \theta \gg 10^{-10} \) unless \( f_a < 10 \) GeV [11]. Therefore, the Planck scale induced effects leave no room for the invisible axion models. On the other hand, the WW model does not have this problem and neither it is sensitive to astrophysical constraints, but it is excluded by the laboratory limits because of too strong couplings with the matter.

In the present paper we suggest a new model for the axion. We assume that there exists a parallel sector of “mirror” particles [13, 14, 15, 16] which is completely analogous to the ordinary particle sector, i.e. it has the same gauge group and all coupling constants are equal in two sectors. In other words, the whole Lagrangian is invariant with respect a discrete mirror parity (M-parity) under the exchange of two sectors. We further assume that the ordinary and mirror worlds have the common PQ symmetry realized à la Weinberg-Wilczek, with the \( U(1)_{\text{PQ}} \) charges carried by the ordinary Higgs doublets \( H_{1,2} \) and their mirror partners \( H_{1,2}' \).

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5 In the case of the hadronic axion [3] or archion [4], a small window around \( f_a \sim 10^6 \) GeV and so the axion mass of few eV can be also permitted.

6 There are also “axionless” models where the solution to the strong CP-problem is due to spontaneously broken P- or CP-invariance [1]. These scenarios, however, have many specifics and do not work in the general context. The Planck scale induced effects for the axionless models have been studied in ref. [2].

7 The possibility to use mirror symmetry for solving the strong CP problem, in different context, was first considered by Rubakov [17]. We will briefly comment on this model later. In principle, two sectors could have also other common symmetries, e.g. the flavour symmetry [4].
If M-parity is an exact symmetry, then the particle physics should be exactly the same in two worlds: in particular, the initial $\theta$-terms are equal, $\theta = \theta'$, the quark mass matrices are identical, $M_{U,D} = M'_{U,D}$, the QCD scales coincide, $\Lambda = \Lambda'$, and the axion couples both QCD sectors in the same way: $f_a^{-1} a (G \tilde{G} + G' \tilde{G}')$, so that their non-perturbative dynamics should produce the same contributions to the axion effective potential. Clearly, in this case the strong CP-problem is simultaneously solved in both worlds – the axion VEV cancels the $\theta$-terms both in the mirror and ordinary sectors. In such a realization, however, $f_a$ remains order $100$ GeV and thus it is excluded by the same phenomenological grounds as the original WW model.

As it was suggested in ref. [13], the M-parity can be spontaneously broken and the electroweak symmetry breaking scale $v'$ in the mirror sector can be larger than the ordinary one $v$. This would lead to somewhat different particle physics in the mirror sector and it is not a priori clear that the strong CP-problem still can be simultaneously fixed in both sectors. However, we show below that this is just the case! The simple reason is that the softly broken M-parity leads to the VEV difference between the ordinary and mirror Higgs doublets, but the structure of the Yukawa couplings remains the same in two sectors.

As far as $U(1)_{PQ}$ is a common symmetry for two sectors, now the PQ scale is determined by the larger VEV $v'$, i.e. $f_a \sim (v'/2) \sin 2\beta'$. The axion state $a$ dominantly comes from the mirror Higgs doublets $H_{1,2}'$, up to small $(\sim v/v')$ admixtures from the ordinary Higgses $H_{1,2}$. Hence, it is a WW-like axion with respect to mirror sector, while it couples the ordinary matter as an invisible DFSZ-like axion. Therefore, the experimental limits on the axion search require $f_a > 10^4$ GeV or so.

However, there are good news concernin the axion mass – for a given scale $f_a$ it appears to be much larger than the mass of the ordinary DFSZ axion $a$. The reason can be briefly explained as follows. Due to the larger weak scale, $v' \gg v$, the mirror quark masses are scaled up by a factor $v'/v$ with respect to the ordinary quarks, $m'_{u,d,\ldots} \sim (v'/v) m_{u,d,\ldots}$. By this reason, also the mirror QCD scale becomes somewhat larger than the ordinary one: $\Lambda' > \Lambda [15, 16]$. Therefore, the axion potential dominantly emerges from the mirror non-perturbative dynamics and we have $m^2_a \sim m'_{q} \Lambda^3 / f_a^2$ or $m^2_a \sim \Lambda'^3 / f_a^2$ (depending on whether the mass of lightest mirror quark $m_{q}'$ is smaller or larger than $\Lambda'$). Therefore, our axion can be much heavier than the ordinary invisible axion with $m^2_a \sim m_{q} \Lambda^3 / f_a^2$. In particular, if with increasing $v'$ also $\Lambda'$ grows fastly enough, our axion can maintain the weight category of the WW axion [4] and its mass can remain in the MeV range even for $f_a > 10^4$ GeV.

The paper is organized as follows. In sect. 2 we present our model and show how it solves the strong CP-problem. We also compute the axion mass and its couplings to the fermions and photons. In sect. 3 we confront our model with the experimental and astrophysical bounds and demonstrate that there is an allowed parameter space which can be of interest for the future experimental search. The stability of our model against the Planck scale induced effects is discussed in sect. 4. Finally, in sect. 5, we summarize our results.

## 2 Mirror World and the Peccei-Quinn symmetry

It was suggested a long time ago that there can exist a parallel (mirror) gauge sector of particles and interactions which is the exact copy of the visible one and it communicates the latter through the gravity and perhaps also via some other interactions. Various phenomenological and astrophysical implications of this idea have been studied in refs. [13, 14, 15, 16].

In particular, one can consider a model based on the gauge symmetry $G \times G'$ where $G = SU(3) \times SU(2) \times U(1)$ stands for the standard model of the ordinary particles: the quark and lepton fields $\psi_i = q_i, \ l_i, \ a^c_i, \ d^c_i, \ e^c_i$ (i is a family index) and two Higgs doublets $H_{1,2},$
while $G' = SU(3) \times SU(2) \times U(1)'$ is its mirror gauge counterpart with the analogous particle content: the fermions $\psi_i' = q_i', l_i', u_i', d_i', e_i'$ and the Higgses $H_{1,2}'$. From now on all fields and quantities of the mirror sector have an apex to distinguish from the ones belonging to the ordinary one. All fermion fields $\psi, \psi'$ are taken in a left-chiral basis.

Let us assume that the theory is invariant under the mirror parity $M$: $G \leftrightarrow G'$, which interchanges all corresponding representations of $G$ and $G'$. Therefore, the two sectors are described by identical Lagrangians and all coupling constants (gauge, Yukawa, Higgs) have the same pattern in both of them. In particular, for the Yukawa couplings

$$L_{\text{Yuk}} = G_{ij}^0 u_i^c q_j + G_{DD}^0 d_i^c q_j + G_{EE}^0 e_i^c l_j + \text{h.c.,}$$

$$L'_{\text{Yuk}} = G_{ij}^0 u_i'^c q_j H_2 + G_{DD}^0 d_i'^c q_j H_1 + G_{EE}^0 e_i'^c l_j H_1' + \text{h.c.},$$

we have $G_{ij,U,D,E}^0 = G_{ij,U,D,E}^0$. In addition, also the initial $\theta$-terms are equal, $\theta = \theta'$.

We further assume that two sectors have a common Peccei-Quinn symmetry $U(1)_{\text{PQ}}$ under which fermions $\psi_i, \psi_i'$ change their phases by a factor $\exp(-i\omega/2)$ and the Higgses $H_{1,2}, H_{1,2}'$ change phases as $\exp(i\omega)$. Then the renormalizable Higgs Potential has a general form $V_{\text{tot}} = V + V' + V_{\text{mix}}$, where for the ordinary Higgses we have:

$$V = -\mu^2_H H_1 H_1 + \mu^2_H H_2 H_2 + \lambda_1(H_1^2 H_1)^2 + \lambda_2(H_2^2 H_2)^2 + \lambda(H_1 H_2)^2 (H_1 H_2),$$

the mirror Higgs potential $V'$ has exactly the same pattern by $M$-parity: $H_{1,2} \rightarrow H_{1,2}'$, and the mixing terms are:

$$V_{\text{mix}} = -\kappa(H_1 H_2)' (H_1' H_2') + \text{h.c.}$$

where the coupling constant $\kappa$ should be real due to $M$-parity.

The mirror parity can be spontaneously broken so that the weak interaction scales are different in two sectors. The easiest way is to introduce a real singlet scalar $\eta$ which is odd under $M$-parity: $\eta \rightarrow -\eta$. Then the following renormalizable terms can be added to the Higgs potential:

$$\Delta V = h(\eta^2 - \mu^2) + \sigma_k \eta^2 (H_k^* H_k + H_k'^* H_k') + \rho_k \mu \eta (H_k^* H_k - H_k'^* H_k'), \quad k = 1, 2$$

A non-zero VEV $\langle \eta \rangle = \mu$ spontaneously breaks the $M$-parity and induces the different effective mass terms for the ordinary and mirror Higgses: $-\mu^2_{\text{eff}} = -\mu^2_k + \sigma_k \mu^2 + \rho_k \mu^2$ and $-\mu^2_{\text{eff}} = -\mu^2_k + \sigma_k \mu^2 - \rho_k \mu^2$. Therefore, the mirror Yukawa's $v'_{1,2}$ are different from the ordinary ones $v_{1,2}$, and in general also the VEV ratios $v_2/v_1 = \tan \beta \equiv x$ and $v_2'/v_1' = \tan \beta' \equiv x'$ must be different in two sectors. In particular, we will be interested in a situation when $v_{1,2} \gg v_{1,2}'$, i.e. the weak interaction scale $v' = (v_1^2 + v_2^2)^{1/2}$ in the mirror sector is much larger than the ordinary weak scale $v = (v_1^2 + v_2^2)^{1/2} = 247$ GeV, $v'/v \ll \zeta \gg 1$. Clearly, the large hierarchy between $v'$ and $v$ requires the fine tuning of the parameters in the Higgs potential. We do not discuss here this question and refer the reader e.g. to ref. [13].

As far as the ordinary and mirror Yukawa constants in (1) are the same, the quark and lepton mass matrices essentially have the same form in both sectors: $M_U = G_U(H_2)$, $M'_U = G_U(H_2')$, $M_D = G_D(H_1)$, $M'_D = G_D(H_1')$, etc. Thus, the fermion mass and mixing structure in the mirror sector should be the same as in the ordinary one, with the mirror up quark masses scaled as $m_{u,e,c,t}' = \xi m_{u,e,c,t}$ and the down quark and charged lepton masses scaled as $m_{d,e} = \xi m_{d,e}$.

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*Notice, that in the limit $\kappa = 0$ there emerge two separate axial global symmetries, $U(1)_A$ for ordinary sector under which $\psi_i \rightarrow \exp(-i\omega/2)\psi_i$ and $H_{1,2} \rightarrow \exp(i\omega)H_{1,2}$, and $U(1)'_A$ for mirror sector: $\psi_i' \rightarrow \exp(-i\omega'/2)\psi_i'$ and $H_{1,2}' \rightarrow \exp(i\omega')H_{1,2}'$. Therefore, the term $V_{\text{mix}}$ demands that $\omega' = \omega$ and thus it reduces $U(1)_A \times U(1)'_A$ to its diagonal subgroup $U(1)_{\text{PQ}}$.  

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once we know that for the ordinary QCD this happens at $\Lambda' \approx \mu = \alpha_q$ quark thresholds $m'_{u,d,s} = \zeta_2 v'_2/v_2 = \zeta_1(v'_1/v_1 = \zeta_1(v'_1/v_1 = \zeta_2 m_t)$. Below it $\alpha'_s$ will have a different slope than $\alpha_s$, and this slope will change every time below the mirror quark thresholds $m'_b \sim \zeta_1 m_b$, etc. In the evolution of $\alpha_s$ these thresholds occur at lower scales, $\mu = m_t, m_b$, etc. Then it is very easy to determine the scale $\Lambda'$ at which $\alpha'_s$ becomes large, once we know that for the ordinary QCD this happens at $\Lambda \sim 200$ MeV. In other words, $\Lambda'$ becomes a function of $\zeta_1,2$, and for $v' \gg v$ one could obtain a significant difference between the QCD scales: $\Lambda' > \Lambda$.

Let us consider for simplicity the case $x = x'$, which yields $\zeta_1 = \zeta_2 = v'/v$. Then for the values of $v'$ up to $10^4$ GeV or so one obtains an approximate scaling law:

$$\frac{\Lambda'}{\Lambda} \approx \left(\frac{v'}{v}\right)^{\frac{b_2-b_n}{2}} \approx \left(\frac{v'}{v}\right)^{0.28} \tag{9}$$

where $b_n = 11 - \frac{2}{3}n$ are the RG coefficients for the case of $n$ light quarks. For $v' > 10^4$ GeV the mirror light quarks masses become larger than $\Lambda'$, since they grow faster with $v'$: $m'_{u,d}/m_{u,d} = v'/v$, and the approximation (9) is not valid – for $\mu < m'_{u,d}$ the mirror QCD becomes the pure gluodynamics and hence $\Lambda'/\Lambda$ starts to increase faster with $v'$ (see Fig. 1).

Let us discuss now the axion physics in our model. Since we are interested only in Goldstone modes, let us express the neutral components of the ordinary and mirror Higgs fields as:

$$H^0_k = \frac{v_k}{\sqrt{2}} \exp(i\phi_k/v_k), \quad H'^0_k = \frac{v'_k}{\sqrt{2}} \exp(i\phi'_k/v'_k); \quad k = 1,2 \tag{10}$$

and identify the combination of the phases $\phi_{1,2}$ and $\phi'_{1,2}$ which corresponds to the axion field. We immediately observe that the states $\phi_Z = v^{-1}(v_2\phi_2 - v_1\phi_1)$ and $\phi'_Z = v'^{-1}(v'_2\phi'_2 - v'_1\phi'_1)$ are the true Goldstone modes eaten up respectively by the standard and mirror gauge bosons $Z$ and $Z'$ which become massive particles by the Higgs mechanism.

In the limit $\kappa = 0$, two remaining combinations $\phi = v^{-1}(v_1\phi_2 + v_2\phi_1)$ and $\phi' = v'^{-1}(v'_1\phi'_2 + v'_2\phi'_1)$ would represent the Goldstone modes of larger global symmetry $U(1)_A \times U(1)'_A$, respectively with the decay constants $f = v_1v_2/v = (v/2)\sin 2\beta$ and $f' = v'_1v'_2/v' = (v'/2)\sin 2\beta'$.
Thus, we would have two axion states separately for the ordinary and mirror sectors, and their masses would emerge respectively due to quantum anomalies of the $U(1)_A$ and $U(1)'_A$ currents related to ordinary and mirror QCD. In particular, $\phi$ would represent the familiar WW axion with order $\sim 1$ MeV mass given by eq. (4).

However, the mixing term \(\mathcal{L}_{\text{mix}}\) reduces the larger global symmetry $U(1)_A \times U(1)'_A$ to a common PQ symmetry $U(1)_{\text{PQ}}$. Substituting \(\mathcal{L}_{\text{mix}}\) in $V_{\text{mix}}$, we obtain:

$$V_{\text{mix}} = -\frac{\kappa}{2} v_1 v_2 v'_1 v'_2 \cos \left( \frac{\phi}{f} - \frac{\phi'}{f'} \right) = \frac{\kappa v v'}{4 f f'} (f' \phi - f \phi')^2 + \ldots$$

(11)

Thus, the combination $A = f_a^{-1} (f' \phi - f \phi')$, where $f_a = (f^2 + f'^2)^{1/2}$, receives a mass$^2$ term:

$$M_A^2 = \frac{\kappa f_a^2}{2} (x + \frac{1}{x}) (x' + \frac{1}{x'}) = \frac{2 \kappa f_a^2}{\sin 2\beta \sin 2\beta'}$$

(12)

which is positive for $\kappa > 0$. The present experimental limits on the Higgs search imply the lower bound on $M_A$ about 100 GeV [8]. As for the other combination $a = f_a^{-1} (f \phi + f' \phi')$, it remains as a Goldstone mode of the $U(1)_{\text{PQ}}$ symmetry to be identified with the axion field.

Let us consider now the effective $\theta$-terms in two sectors. Recalling that $\theta = \theta'$ and $G_{U,D} = G'_{U,D}$ due to M-parity, we obtain:

$$\bar{\theta} = \theta + \arg(\det M_U \det M_D) = \theta + \arg(\det G_U \det G_D) + N \langle \phi / f \rangle$$

$$\bar{\theta}' = \theta' + \arg(\det M'_U \det M'_D) = \theta' + \arg(\det G'_U \det G'_D) + N \langle \phi' / f' \rangle$$

(13)

where the last terms describe contributions of the phases in (11), and they can be immediately rewritten as $\phi / f = f_a^{-1} [a - (f / f') A]$ and $\phi / f = f_a^{-1} [a + (f / f') A]$. We see that the axion field $a$ contributes both eqs. (13) in the same way, whereas the heavy state $A$ is irrelevant since it has a vanishing VEV. Therefore, $\bar{\theta}$ and $\bar{\theta}'$ will be simultaneously cancelled by the axion VEV $f_a^{-1} (a)$ and so in our model the strong CP-problem is solved in both sectors.

In addition, for the case $f' \gg f$ we obtain a quite interesting situation. The axion state $a$ dominantly emerges from the phases of the mirror Higgses $H'_{1,2}$, while it contains only a small admixture from the ordinary Higgses $H_{1,2}$: $a \approx \phi' - (f / f') \phi$. Therefore, it appears to be a WW-like axion with respect to the mirror sector, while from the view of the ordinary sector it is just a DFSZ-like invisible axion with a decay constant $f_a \approx f' = (v'/2) \sin 2\beta'$.

However, there is a dramatic difference from the DFSZ case. Since the mass of our axion receives contributions from both ordinary and mirror QCD, in a full analogy to eq. (3) one can write:

$$m_a^2 = \frac{N^2}{f_a^2} \left( \frac{VK}{V + K \text{Tr} M^{-1}} + \frac{V'K'}{V' + K' \text{Tr} M'^{-1}} \right)$$

(14)

where $M = \text{diag}(m''_u, m''_d)$ is the mass matrix of the mirror light quarks, and the values $K' \sim \Lambda'^4$ and $V' \sim \Lambda'^4$ characterize the mirror gluon and quark condensates, with $\Lambda'$ being the QCD scale in the mirror sector. The first term in this expression gives an ordinary small contribution and it is negligible in comparison to the second term which can give much larger contribution.

\* In the following, we maintain parametrically the number of fermion families, $N = N_g$, in general formulas, but for the numerical calculations we always take $N = 3$.

\*\* Actually $\bar{\theta}$ can get radiative corrections due to the fact that the electroweak scales are different in two sectors. However, it is well-known that these corrections are negligibly small.

\*\*\* This is very easy to understand from the following observation. The combination $H'_1 H'_2 \equiv S$ in (5) is a standard model singlet carrying the double PQ charge, and it can be considered as a composite operator which communicates the standard Higgses $H_{1,2}$ via the term $V_{\text{mix}} = (H_1 H_2) S + h.c.$, just as in the DFSZ model.
as far as \( m'_{u,d} \gg m_{u,d} \) and \( \Lambda' > \Lambda \). Certainly, the exact form of the latter term depends on how many light quarks (with masses less than \( \Lambda' \)) are contained in mirror sector. In particular, for the cases of two \((u' \text{ and } d')\), one \( (u') \), or zero light quarks, the second term in brackets of eq. (14) becomes respectively \((1 + z')^{-1} m'_{u'} v', m'_{d'} v' \) or \( K' \), where \( z' = m'_{u}/m'_{d} = (\zeta_{2}/\zeta_{1})z = (x'/x)z \).

Let us take for example the case \( x = x' \), and consider the regime when at least \( u' \) is light, \( m'_{u} = (v'/v)m_{u} < \Lambda' \) (recall that this regime holds up to the values \( v' \sim 10^{4} \text{ GeV or so} \)). Then for the axion mass holds an approximate scaling low:

\[
m_{a} \approx \frac{2N}{v' \sin 2\beta'} \left( \frac{m'_{u} V}{1 + z'} \right)^{1/2} \approx \left( \frac{N}{\Lambda} \right)^{3/2} \left( \frac{v}{v'} \right)^{1/2} m_{a}^{WW} \tag{15}\]

with respect the ordinary WW axion mass \([4]\). Using also eq. (14), we thus obtain:

\[
m_{a} \approx \left( \frac{v}{v'} \right)^{0.08} \times 150 \text{ keV} \sin 2\beta', \tag{16}\]

In general case, for very large \( v' \) and arbitrary \( x' \), the simple scaling low \([13]\) is no more valid, since for \( x \neq x' \) the masses of \( u' \) and \( d' \) scale differently and in addition, with increasing \( v' \), only one or none of them may remain light, with mass less than \( \Lambda' \). In this situation one has to use the general formula \([14]\). Nevertheless, \( m_{a} \) still decreases very slowly with increasing \( v' \) and for enough large \( x, x' \), its value can be of the order of MeV.

The results of numerical computations from the general formula \([14]\) for some interesting cases are presented in Fig. 2. The relevant parameters are taken as \( x = \tan \beta, x' = \tan \beta' \) in the range \( 1 - 100 \), and \( f_{a} = v'/(x' + x'^{-1}) \). We use one-loop RG evolution for the strong coupling constants with \( \Lambda = 200 \text{ MeV} \). The following values of the quark masses are used: \( m_{u} = 4 \text{ MeV}, m_{d} = 7 \text{ MeV}, m_{s} = 150 \text{ MeV (at } \mu = 1 \text{ GeV), and } m_{c}(m_{c}) = 1.3 \text{ GeV}, m_{b} = 4.3 \text{ GeV and } m_{t} = 170 \text{ GeV (respectively at } \mu = m_{c}, m_{b}, m_{t}) \). For the parameters \( V \) and \( K \) related to the quark and gluon condensates we take \( V = (250 \text{ MeV})^{2} \) and \( K = (230 \text{ MeV})^{4} \), and assume that the similar parameters in the mirror sector scale as \( V' / V = (\Lambda'/\Lambda)^{3} \) and \( K' / K = (\Lambda'/\Lambda)^{4} \). We see that for enough large values of \( x, x' \) the axion mass can be order 1 MeV even for \( f_{a} = 10^{4} - 10^{6} \text{ GeV} \).

Let us discuss now the interaction terms \([1]\). Clearly, our axion couples the ordinary fermions and photons in exactly the same way as the DFSZ axion with the PQ scale \( f_{a} \). Hence, for the Yukawa coupling constants to the electron, down quark and up quark we can use the familiar expressions \([1]\):

\[
g_{ae} = \frac{m_{e}}{f_{a}} \sin^{2} \beta, \quad g_{ad} = \frac{m_{d}}{f_{a}} \left[ \sin^{2} \beta - \frac{N}{1 + z} \right], \quad g_{au} = \frac{m_{u}}{f_{a}} \left[ \cos^{2} \beta - \frac{Nz}{1 + z} \right], \tag{17}\]

while for the axion-photon coupling we have:

\[
g_{a\gamma} = \frac{\alpha}{2\pi} \frac{NC}{f_{a}}, \quad C = \frac{8}{3} - \frac{6KTr(M^{-1}Q^{2})}{V + KTr(M^{-1})} \approx \frac{2z}{1 + z} \tag{18}\]

where the trace is taken over the light quark states \((u,d)\), \( Q \) are their electric charges \((+4/3, -1/3)\), and \( \alpha = 1/137 \) is the fine structure constant. In addition, our axion couples to mirror photons, with the constant:

\[
g'_{a\gamma} = \frac{\alpha}{2\pi} \frac{NC'}{f_{a}}, \quad C' = \frac{8}{3} - \frac{6K'\text{Tr}(M'^{-1}Q'^{2})}{V' + K'\text{Tr}(M'^{-1})} \tag{19}\]

where the factor \( C' \) for the case of two \((u',d')\), one \( (u') \) or no light quarks respectively takes the values \( 2z'/(1 + z'), 0 \) and \( 8/3 \).
Figure 2: Contour plots for the axion mass (solid) as a function of $f_a$ and $x'$, for the cases $x = x'$ and $x = 60$. The regions left from the dash-dots curves are excluded by the observed neutrino signal from SN 1987A while the regions right from the dash curves are excluded by the SMM bounds on the gamma ray flux from SN 1987A. The vertical dotted line indicates a conservative bound on the PQ scale from the terrestrial experiments.

Hence, the axion decay widths into the visible and mirror photons respectively are:

$$\Gamma(a \rightarrow \gamma\gamma) = \frac{g_{ae}^2 m_a^3}{64\pi}, \quad \Gamma(a \rightarrow \gamma'\gamma') = \frac{g_{a\gamma'}^2 m_a^3}{64\pi}, \quad (20)$$

In addition, if $m_a > 2m_e$, the axion can decay also into electron-positron pair:

$$\Gamma(a \rightarrow e^+e^-) = \frac{g_{ae}^2 m_a}{8\pi} \sqrt{1 - \frac{4m_e^2}{m_a^2}} \quad (21)$$

3 Experimental and astrophysical bounds

Let us discuss now the constraints which should be imposed on the axion mass $m_a$ and its coupling constants for fulfilling the existing experimental and astrophysical bounds. Anticipating our results, we remark that the consistent parameter region corresponds to large values of $x = \tan \beta$ and so $\sin \beta \approx 1$. Hence, the axion coupling constants to the electron is $g_{ae} = 5 \times 10^{-10} f_a^{-1}$, and for the axion-nucleon coupling constants we have the following estimates $g_{an} = -1.6 \times 10^{-6} f_a^{-1}$ and $g_{aN} = 0.6 \times 10^{-6} f_a^{-1}$ (see e.g. in ref. [8]), where $f_a = (f_a/10^6 \text{ GeV})$.

The conservative interpretation of the experimental constraints from the axion search [8] imply the lower limit on the axion-nucleon coupling constants $g_{aN} < 10^{-4}$ or so, which translates as $f_a > 10^{-2}$. For $m_a > 2m_e$, the somewhat stronger limit emerges from the reactor search of $a \rightarrow e^+e^-$ decay. Namely, for $m_a = 1.4 - 2$ MeV one has $f_6 > 0.1$ [20]. The latter limit, however, is subject of various theoretical uncertainties, and in the following we take the most conservative lower limit $f_a > 10^4 \text{ GeV}$. The existing limits on the axion decay into two photons are also fulfilled in this case [21].

For the standard DFSZ axion, since it is very very light, the severe limits on the scale $f_a$ emerge from the stellar evolution. In particular, the astrophysical bounds on the white dwarf luminosity function, helium burning lifetime of the HB stars or helium ignition time in low mass red giants imply $g_{ae} < (2.5 - 5) \times 10^{-13}$, and hence $f_a > (1 - 2) \times 10^9 \text{ GeV}$ [8]. However,
these limits do not apply to our axion which is quite heavy and so its production rate in the stellar cores, with typical temperatures $T$ up to 10 keV, is suppressed by the exponential factor $\exp(-m_a/T)$. For $m_a > 300$ keV this factor becomes enough small to render our model safe as far as the stellar evolution limits are concerned.

However, stringent bounds on $f_a$ emerge due to the observed neutrino signal from the SN 1987A which indicates that the supernova core cooling rate due to axion emission at the typical time $t = 1$ s from the collapse should not exceed $10^{51}$ erg/s or so. If the axion-nucleon couplings are enough large, $g_{\text{ap}}, g_{\text{an}} > 10^{-7}$, the axions are trapped in the collapsing core and are emitted with the thermal spectrum from an axisphere with a radius $R \sim 10$ km, and energy luminosity at $t = 1$ s can be estimated as $L_a \approx f_6^{16/11} \times 3 \cdot 10^{50}$ erg/s \[22, 23\]. Hence, for $f_6 < 3$ the axion luminosity $L_a$ becomes smaller than $10^{51}$ erg/s, and the total energy and duration of the SN 1987A neutrino burst should not be affected. This consideration is certainly applicable for an axion with mass of order MeV, and so we impose the upper limit $f_a < 3 \cdot 10^6$ GeV or so.\[12\]

In addition, the supernova bounds provide strong limits on the axion decay rates into the ordinary as well as mirror photons. For $m_a \sim 1$ MeV these decays are rather fast and strongly restrict the allowed parameter space.

As far as the ordinary matter is transparent for the mirror photons, the emission of the latter can lead to the unacceptably fast cooling of the supernova core. The energy lose rate $L'_\gamma$ related to the mirror photon emission from the core volume can be estimated as follows. For $f_6 < 10^7$ GeV the axions are strongly trapped inside a core of the radius $R \sim 10$ km and have a thermal distribution. The temperature $T$ is a function of the radial distance $r$ from the center, and it also slowly changes with the time $t$ elapsed from the collapse. The decay rate for an axion with the energy $E$ into mirror photons is $(m_a/E)\Gamma'$, where $\Gamma' = \Gamma(a \rightarrow \gamma'\gamma')$ is given by eq. (20). Therefore, we have $L'_\gamma(t) = 4\pi m_a \Gamma' \int_0^R r^2 n_a(r, t) dr$, where $n_a = 1.2T^3/\pi^2$ is the axion number density. Taking a core temperature $T$ of about $20 - 30$ MeV, one can roughly estimate that $L'_\gamma \simeq \Gamma' m_a (4\pi R^3/3)(1.2T^3/\pi^2) \sim 10^{-2} g_{\gamma'\gamma}^2 m_a^4$. Hence, the condition $L'_\gamma < 10^{51}$ erg/s implies that the decay width $\Gamma'$, for $m_a \sim 1$ MeV, should not exceed few s$^{-1}$. In other terms, using eq. (19), we obtain roughly the bound $m_a^2/f_a < 10^{-7}$ MeV.

By taking the typical temperature profiles in the supernova core at different time moments, one can compute the value $L'_\gamma(t) = 4\pi \Gamma' m_a \int_0^R r^2 n_a(r, t) dr$ more precisely. In addition, $m_a$ and $g_{\gamma'\gamma}$ can be expressed as functions of $x, x'$ and $f_a$ according to eqs. (14) and (19). In Fig. 2 we show isocontours (dash-dott) corresponding to $L'_\gamma = 10^{51}$ erg/s at $t = 1$ s from the collapse, which in fact indicate a lower limit on the PQ scale $f_a$ for given $x$ and $x'$.

Another constraint on our model emerges due to the axion decay into ordinary photons. For $f_a > 10^4$ GeV the flux of axions emitted from the collapsing core is quite large, and thus if the axions would decay into the visible photons in their travel to the earth, one would observe the gamma ray burst associated with the SN 1987A. However, the SMM satellite data \[24\] sets the stringent upper limit on the $\gamma$ flux from the SN 1987A. In particular, these data indicate that photon fluence for for the photon energy band $E_\gamma = 4.1 - 6.4$ MeV integrated over first 10 s from the collapse, did not exceed one photon per cm$^2$. Therefore, the only way to to suppress the gamma ray signal from is to assume that the dominant part of the axions have undergone the decay inside the star envelope, during the flight time $t_s = c R_s \sim 100$ s, where $R_s \sim 3 \cdot 10^7$ km is a radius of the SN 1987A progenitor star, the blue giant Sanduleak. Hence, for the relevant axion energies $E \simeq 2E_\gamma = 8 - 13$ MeV, the exponential factor $\exp(-m_a \Gamma_{\text{tot}} t_s/E)$ should be very small, where the total decay width $\Gamma_{\text{tot}}$ incorporates the decays into the ordinary and mirror photons \[21\], and if $m_a > 2m_e$, also into electron-positron \[21\]. Therefore, $\Gamma_{\text{tot}}$ should

\[12\] The region $f_a > 10^{10}$ is also allowed by the SN 1987A \[1\], however it is not of our interest in this paper.
not exceed few seconds.

More precisely, by taking that the axions emitted from the axiosphere have a thermal spectrum with a temperature which at the time moment $t$ from the collapse estimated as $T_a \approx t_0^{4/11} (t/10 \text{ ms})^{-1/4} \times 6 \text{ MeV}$ \cite{23}, one can compute the expected photon fluence in the energy band $E_\gamma = 4.1 - 6.4 \text{ MeV}$ and confront it to the SMM bound. The results for various values of the parameter $x$ are shown in Fig. 2. Here the dash curves correspond to the SMM limit on the $\gamma$ fluence and so they set an upper limit on $f_a$ for given values of $x$ and $x'$.

Thus, we see that there is a parameter space compatible with the existing experimental and astrophysical limits, which yields $f_a$ around $10^4 - 10^5 \text{ GeV}$ and $m_a$ about 1 $\text{ MeV}$. This parameter range can be accessible for the experimental testing at the future reactor or beam dump experiments.

4 Planck scale corrections

A potential problem for the invisible axion models arises from the Planck scale effects. Many arguments suggest that the non-perturbative quantum gravity effects can induce the higher order effective operators cutoff by the Planck scale $M_P$. Many arguments suggest that the non-perturbative quantum gravity effects can induce the higher order effective operators cutoff by the Planck scale $M_P$. Let us take the DFSZ model \cite{4} in which $U(1)_{\text{PQ}}$ symmetry is spontaneously broken by the gauge singlet scalar $S$, $\langle S \rangle = f_a/\sqrt{2}$ and consider a dimension 5 operator ($\lambda S(S^\dagger S)^2/M_P + \text{h.c.}$, with $\lambda$ being some complex coupling constant. This term explicitly breaks the PQ symmetry and induces the mass term to the axion field:

$$M_a^2 = \frac{\lambda f_a^3}{2\sqrt{2} M_P} \approx \lambda f_a^3 \times 0.17 \text{ GeV}^2$$

which should be compared with the dynamical mass term $m_a^2$ \cite{3}. The potential for $\bar{\theta} = a/f_a$ then becomes $f_a^{-2}P(a) = m_a^2 \cos \bar{\theta} + M_a^2 \cos(\bar{\theta} + \delta)$, where $\delta$ is related to the complex phase of the coupling constant $\lambda$ and it is generically order 1. Although the first ”genuine” term in this potential would fix the axion VEV at $\langle \bar{\theta} \rangle = 0$, the second term does not care about the strong CP problem and tends to minimize the potential along $\langle \bar{\theta} \rangle = -\delta$. Therefore, $\theta < 10^{-10}$ demands that $M_a < 10^{-5} m_a$, which in turn translates into the upper limit $f_a < \lambda^{-1/5} \times 10 \text{ GeV}$. This is in sharp contradiction with the lower limit $f_a > 10^{10} \text{ GeV}$ unless the constant $\lambda$ is extremely small, $\lambda < 10^{-45}$. Therefore, if one takes the Planck scale induced effects seriously, there is no room left for the invisible axion models.

As for our model, it appears to be stable against the Planck scale corrections by threefold reason: (i) the PQ scale is lower, $f_a = 10^4 - 10^5 \text{ GeV}$, (ii) the axion is heavier, $m_a \sim 1 \text{ MeV}$, and (iii) the minimal possible explicit Planck scale operators in the Higgs potential are of dimension 6 and so they are suppressed $M_P^2$. Let us consider, e.g. the terms:

$$\lambda \frac{(H_1 H_2)^3}{M_P^6} + \lambda \frac{(H_1' H_2')^3}{M_P^6} + \text{h.c.}$$

They induce the axion explicit mass term:

$$M_a^2 = \frac{9 \lambda (x' + x'^{-1}) f_a^4}{4 M_P^2} \approx \lambda x' f_a^4 \times 0.5 \cdot \text{keV}^2$$

and the constraint $M_a < 10^{-5} m_a$, for a typical value $m_a \sim 1 \text{ MeV}$, implies $f_a < 1.2 (\lambda x')^{-1/4} 10^5 \text{ GeV}$, which is consistent with the parameter range found in previous section.
The Planck scale effects could explicitly break the PQ symmetry also in the Yukawa like operators \([\mathcal{L}^2]\), e.g.
\[
a^{ij} \frac{H_1 H_2}{M_P^2} d^*_i q_j H_1 + a^{ij} \frac{H'_1 H'_2}{M'_P^2} d^*_i q'_j H'_1
\]
(25)
which would interfere with the contribution to the Yukawa terms (5) and could deviate \(\bar{\theta}\) from zero. However, for \(f_a < 10^6\) GeV these corrections are also negligible and do not give problems.

Let us now comment about the model of Rubakov \([17]\) which was based on the grand unified picture \(SU(5) \times SU(5)\). It was assumed that the mirror GUT symmetry \(SU(5)\) breaks down to the mirror standard model gauge group \(SU(3)' \times SU(2)' \times U(1)'\) at lower scale than the ordinary \(SU(5)\) breaks to standard gauge group. In other words, the VEV of the Higgs 24-plet \(\Phi'\) of \(SU(5)\) is somewhat lower than the VEV of the 24-plet Higgs of \(SU(5)\), \(\langle \Phi \rangle > \langle \Phi' \rangle\). Then the difference between the infrared poles for the strong coupling constants \(\Lambda'\) and \(\Lambda\) can emerge obtained due to the RG evolution from the corresponding scales down in energy. The gauge constant of \(SU(5)\) runs faster then the one of \(SU(3)\), and thus it is possible to have the bigger QCD scale in mirror sector, \(\Lambda' \gg \Lambda\). In this way one could obtain obtain a very massive axion (e.g. \(m_a \sim 100\) GeV or more).

However, in this case the M-parity is broken at GUT scales and therefore there is no strong reason to expect that it will be valid in the low energy Yukawa sector. For example, the Planck scale operators of dimension 5 are dangerous:
\[
G 510 \bar{H} + G' 5'10' \bar{H}' + k \frac{\Phi}{M_P} 510 \bar{H} + k' \frac{\Phi'}{M_P} 5'10' \bar{H}'...
\]
(26)
where \(H\) is the 5-plet Higgs. Due to M-parity we have \(G = G'\) and \(k = k'\) but the effective Yukawa coupling \(\overline{G} = G + k\langle \Phi \rangle / M_P\) is not equal to the mirror effective Yukawa coupling \(\overline{G'} = G + k'\langle \Phi' \rangle / M_P\). This should cause the big phase difference between them and so that \(\overline{G} \neq \overline{G'}\). Therefore, the model \([17]\) is not stable against the Planck scale corrections.

A way to overcome this problem is to break M-parity only at lower energies. E.g., one can take \(SU(5) \times SU(5)\) model with \(\langle \Phi \rangle = \langle \Phi' \rangle\), and break M-parity only due to different electroweak scales in two sectors, so that \(\overline{G} = \overline{G'}\). Clealy, this is nothing but embedding of our model considered in this paper in the grand unified picture.

### 5 Conclusions

We have suggested a new possibility for solving the strong CP-problem. It assumes that in parallel to the observable particle world described by the standard model, there exists a mirror world of particles with the identical Lagrangian, and two sectors share the same Peccei-Quinn symmetry realized à la Weinberg-Wilczek model. We assume that the mirror symmetry between two worlds is spontaneously broken so that the electroweak scale in the mirror sector \(v'\) is considerably larger than the ordinary electroweak scale \(v = 247\) GeV. This in turn implies the infrared scale \(\Lambda'\) of the mirror strong interactions has to be somewhat larger than the ordinary QCD scale \(\Lambda \simeq 200\) MeV. In this way, the axion mass and interaction constants are actually determined by mirror sector scales \(v'\) and \(\Lambda'\), while the \(\theta\) terms are simultaneously cancelled in both sectors due to mirror symmetry.

This means that the axion state essentially emerges from the mirror Higgs doublets, up to small (\(\sim v/v'\)) admixture from the ordinary Higgses. Hence, our axion is a WW-type axion for the mirror sector while from the point of view of ordinary sector it is DFSZ-type invisible axion with the symmetry breaking scale \(f_a \sim v' \gg v\).
We have discussed the experimental and astrophysical limits on such an axion and have found an interesting parameter window where the PQ scale is \( f_a \sim 10^4 - 10^5 \) GeV and the axion mass is \( m_a \sim 1 \) MeV, which is accessible for the axion search in the future reactor and beam dump experiments. Interestingly, for such a parameter range our model is stable against the Planck scale corrections as far as the strong CP-problem is concerned.

In our model, which in fact is the simplest mirror extension of the standard model, we cannot approach parameter space which could be applicable for explanation of the Gamma Ray Bursts via axion emission from the collapsing compact objects \[25\]. This, however, could be easily achieved in the supersymmetric version of our model, due to stronger effects in the renormalization group evolution of the strong coupling constants in two sectors.

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