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Mirror world, supersymmetric axion and gamma ray bursts

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Abstract: A modification of the relation between axion mass and the PQ constant permits a relaxation of the astrophysical constraints, considerably enlarging the allowed axion parameter space. We develop this idea in this paper, discussing a model for an *ultramassive* axion, which essentially represents a supersymmetric Weinberg-Wilczek axion of the mirror world. The experimental and astrophysical limits allow a PQ scale $f_a \sim 10^4 - 10^6$ GeV and a mass $m_a \sim$ MeV, which can be accessible for future experiments.
On a phenomenological ground, such an *ultramassive* axion turns out to be quite interesting. It can be produced during the gravitational collapse or during the merging of two compact objects, and its subsequent decay into e^+e^- provides an efficient mechanism for the transfer of the gravitational energy of the collapsing system to the electron-positron plasma. This could resolve the energy budget problem in the Gamma Ray Bursts and also help in understanding the SN type-II explosion phenomena.

Keywords: Global Symmetries, Cosmology of Theories beyond the SM, CP violation.

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1. Introduction

The strong CP problem is one of the most puzzling points of modern particle physics (for a general reference see e.g. [1]). It resides in the presence of the so-called $\theta-$term in the QCD lagrangian, $\mathcal{L}_\theta = \theta(\alpha_s/8\pi)G_{\mu\nu}\tilde{G}^{\mu\nu}$, where α_s is the fine structure constant of the strong interactions and $G_{\mu\nu}$ is the gluon field strength tensor. The *θ-term*, which receives a contribution also from the complex phases in the quark mass matrices $M_{U,D}$ so that its effective value is $\overline{\theta} = \theta + \arg(\det \mathrm{M_U} \det \mathrm{M_D})$, is CP violating and leads to a neutron electric dipole moment d_n, experimentally not observed. This implies a very strong upper limit on the parameter $\overline{\theta}$, $|\overline{\theta}| < 10^{-9}$, which has no theoretical explanation in the context of the Standard Model.

In the most appealing solution of this problem, the Peccei-Quinn (PQ) mechanism [2], the θ parameter[1] becomes a dynamical field, the axion $a = f_a\theta$, and emerges as the pseudo-Goldstone mode of a spontaneously broken global axial symmetry $\mathrm{U(1)_{PQ}}$. Here f_a is a constant, with dimension of energy, also called *axion decay constant*. We will use the following convention throughout this paper [9]: we indicate with $f_{\rm PQ}$ the VEV of a scalar (or a VEV of a combination of several scalars) responsible for the $\mathrm{U(1)_{PQ}}$ symmetry breaking. The constant f_a, which characterizes axion phenomenology, is defined as $f_{\rm PQ}/N$, where N stands for the color anomaly of $\mathrm{U(1)_{PQ}}$ current.[2]

At quantum level, the PQ symmetry is broken by the chiral anomaly, and since this is a dynamical effect, its strength is measured by non-perturbative QCD contributions. Therefore all the axion properties are essentially related only to the PQ scale $f_{\rm PQ}$ and the QCD scale Λ.

[1]In the following we will use θ instead of $\overline{\theta}$ for simplicity.

[2]The PQ charges are normalized so that each of the standard fermion families contributes as $N = 1$. Therefore, in the Weinberg-Wilczek (WW) model [3] 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 [4]. Other models of the invisible axion, e.g. the hadronic axion [5] or archion [6], generally contain some exotic fermions and so $N \neq N_g$.

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On a phenomenological ground, all the axion characteristics can be roughly estimated from the pion properties. Indeed, axions generically mix with pions so that their mass, as well as their couplings with photons and nucleons, are roughly given by f_π/f_a times those of the π-meson, where $f_\pi \approx 93\,\mathrm{MeV}$ is the pion decay constant. So it is clear that the PQ constant is the relevant scale for axion phenomenology. For example, in the most general context, axion interaction with fermions and photons can be described by the following lagrangian terms:

$$\mathcal{L}_a = ic_{ae}\frac{m_e}{f_a}\, a\,\bar{e}\gamma_5 e + ic_{aN}\frac{m_N}{f_a}\, a\,\bar{N}\gamma_5 N + c_{a\gamma}\frac{\alpha}{8\pi f_a}\, a\, F_{\mu\nu}\tilde{F}^{\mu\nu} + \cdots\,, \tag{1.1}$$

where m_k represents the fermion mass (e.g. $m_e, m_N, \ldots$ for electrons, nucleons, etc.), α is the fine structure constant and c_{ai} are constant coefficients. The factors c_{aN}, which refer to axion-nucleon interaction, $N = p, n$, are generally of order one, while for the axion-electron and axion-photon interaction the related coefficients c_{ae}, $c_{a\gamma}$ are model dependent.

As for the axion mass, in general it can be obtained from the expression [7]:

$$m_a^2 = \frac{1}{f_a^2}\frac{VK}{V + K\mathrm{Tr}M^{-1}}\,, \tag{1.2}$$

where $V = \langle\bar{q}q\rangle \sim \Lambda^3$ is the light quark condensate responsible for the chiral symmetry breaking, $M = \mathrm{diag}(m_u, m_d, \ldots)$ is the mass matrix of light quarks (with $m_q < \Lambda$) and $K \sim \Lambda^4$ accounts for the strength of the instanton induced potential. Then, by taking into acount only two light quarks, u and d, and using the relation $(m_u + m_d)\langle\bar{q}q\rangle = m_\pi^2 f_\pi^2$, from (1.2) one directly arrives to the more familiar formula:

$$m_a = \frac{1}{f_a}\left(\frac{m_u m_d V}{m_u + m_d}\right)^{1/2} = \frac{z^{1/2}}{1+z}\cdot\frac{f_\pi m_\pi}{f_a} \approx \left(\frac{10^6\,\mathrm{GeV}}{f_a}\right) \times 6.2\,\mathrm{eV}\,, \tag{1.3}$$

where $z = m_u/m_d \simeq 0.57$.

In the WW model [3] 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 $v = (v_1^2 + v_2^2)^{1/2} \simeq 247\,\mathrm{GeV}$ is the electroweak scale and $\tan\beta = v_2/v_1$. Therefore, the WW axion is quite heavy, and its mass can vary from a few hundred keV to several MeV:

$$m_a^{\mathrm{WW}} = \frac{2N}{v\sin 2\beta}\left(\frac{m_u V}{1+z}\right)^{1/2} \approx \frac{150\,\mathrm{keV}}{\sin 2\beta}\,. \tag{1.4}$$

However, its couplings with fermions and photons are too strong and for this reason the WW model is completely ruled out for any values of the parameter β by a variety of terrestrial experiments such as the search of the decay $K^+ \to \pi^+ a$, the J/ψ and Υ decays into $a + \gamma$, the nuclear deexcitations via axion emission, the reactor and beam dump experiments, etc. [8].

For any realistic axion model, these experimental constraints generally imply the lower bound on the PQ scale, $f_a \gtrsim 10^4\,\mathrm{GeV}$, which in turn implies $m_a \lesssim 1$ keV. This is what happens in all the so-called invisible axion models. Anyway, such a light axion can easily be produced inside a star at a temperature of a few keV, and can accelerate the cooling process in a dangerous way. So the axion is required to interact only weakly with fermions in order to strongly suppress the energy transport process inside the star. Also, the axion luminosity from the SN core must be constrained in order to not ruin the neutrino signal detected at the time of SN 1987A explosion [9]. These astrophysical considerations exclude all scales f_a up to 10^{10} GeV [1, 9], and so $m_a < 10^{-4}$ eV.[3] 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 $f_a \lesssim 10^{11} - 10^{12}$ GeV [1, 9]. Thus, not much parameter space remains available.

Is it possible to relax the astrophysical constraints? It is quite interesting to note that all astrophysical constraints from stellar evolution could be satisfied for the PQ scales above the laboratory limit, $f_a > 10^4$ GeV, up to values order 10^7 GeV, if the axion mass would remain in the MeV range, as with the mass of the WW axion (1.4). However, the tight relation between the axion mass m_a and the PQ scale f_a (1.3) is very constraining and does not allow such a situation.

If there were, in fact, another source for axion mass, e.g. from Plank scale induced effects, then the axion could change its properties and then no longer be useful for the solution of the strong CP problem. This explains why this relation is universally accepted, and in several papers the astrophysical and cosmological limits are given in terms of axion mass instead of the PQ constant.

But is the relation (1.3) really universal? What would happen, in fact, if the axion could communicate with another, *hidden*, sector of particles and interactions? In general, if the axion potential were to get dominant contribution from the hidden sector, then we would expect it to solve the strong CP problem for that sector rather than for our observable world. However, this is not necessarily the case.

In particular, one can assume that the hidden sector is a mirror world, a parallel sector of "mirror" particles and interactions with the lagrangian completely similar to that of the ordinary particles [10]. In other words, it has the same gauge group and coupling constants as the ordinary sector, so that the lagrangian of the whole theory is invariant with respect to the Mirror parity (M-parity) under the interchange of the two sectors. Several phenomenological and astrophysical implications of the mirror world have been studied in the literature [11]. In particlular, it could provide a new insight for the cosmological dark matter made up by mirror baryons [12]. For a recent review of the mirror matter concept one can refer to [13].

We can further assume that the ordinary and mirror sectors have the common Peccei-Quinn symmetry [14], with the same PQ charges carried by the ordinary Higgs doublets $H_{1,2}$ and their mirror partners $H'_{1,2}$.

If the M-parity is an exact symmetry, then the particle physics should be exactly the same in two worlds, and so the strong CP problems would be simultaneously solved in both sectors. In particular, the axion would couple to both sectors in the same way and their non-perturbative QCD dynamics would produce the same contribution to the axion effective

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

potential. This situation would not bring drastic changes of the axion properties; just the axion mass would increase by factor of $\sqrt{2}$ with respect to the standard expression (1.3).[4]

However, a very interesting situation emerges if the M-parity is spontaneously broken in the Higgs sector, so that the mirror electroweak scale v' is considerably larger than the ordinary one v [16, 17].[5] In this case, as it was shown in ref. [17, 14], also the mirror QCD scale Λ' becomes larger than the ordinary one Λ and thus one expects that the dominant contribution to the axion potential comes from the mirror sector. On the other hand, if the latter contribution is predominant, the axion mass relation with the PQ scale could change in a considerable way. Of course, in this case it is absolutely not evident that the axion can still solve the strong CP problem. Indeed, if we ask for it to be dominantly governed by the mirror QCD, it is natural to expect that it will cancel the *mirror* and not the ordinary θ-term. Fortunately, this is not the case: this axion can still solve our strong CP problem, as far as we ask for the Yukawa structure to be the same in the two sectors [14]. So it appears that the mirror extension of the standard axion model is the only way that allows for a quite heavy axion without spoiling the PQ mechanism.

In particular in this paper we will consider an axion with the mass $m_a > 1\,\mathrm{MeV}$ and with a PQ scale $f_a \sim 10^6\,\mathrm{GeV}$. Certainly such a massive axion cannot ruin stellar evolution process, even though it can still be produced in the hot SN core at a temperature of a few $10\,\mathrm{MeV}$. We expect then that the available parameter range is sizeably increased with respect to the standard DFSZ axion model.

As we will show, such an axion can have many interesting astrophysical implications. In particular, it can help in understanding the Gamma Ray Bursts (GRBs) and the supernova explosion phenomena via the mechanism suggested in ref. [19].[6] In fact, such a non standard axion can be produced during the merging of two compact objects, and then, due to its decay, efficiently transfer the gravitational energy of the collapsing system into an ultrarelativistic e^+e^- plasma, the fireball, far from the impact place. Also such a heavy axion, produced in the core of a SN type-II, can decay into $e^+\ e^-$ before reaching the stellar surface, transferring in this way a huge amount of energy at a distance of about $1000\,\mathrm{km}$ from the stellar core, helping the SN explosion (thermal bomb). On the other hand, because of the different size of SN type-Ib/c, some axions will be able to leave their surface and then decay into photons, explaining the observed events of weak GRBs related to the SN type-I.

This logic was essentially suggested in a previous paper [14], but the allowed mass was not enough to open a new window for the axion of interest in astrophysics. In fact, the parameter space in [14] does not allow for an $f_a >$ few $10^4\,\mathrm{GeV}$. Observe that the QCD scale in the mirror sector (and , as a consequence, the axion mass) is fixed by the renormalization group (RG) evolution of the strong coupling constant, and so it depends on

[4]However, even in this less interesting case there would be significant modification of the primordial oscillations of axion field and their contribution to dark matter of the Universe [15].

[5]Cosmological implications of such a scenario were studied in refs. [17, 18].

[6]We underline that the possible relation between GRBs and supernova was contemplated in the past. In [20], e.g., is considered the possibility that the emission of a light axion from SN can provide the energy necessary for the GRBs production.

the matter content of the theory, as well as on the mirror fermion masses. This observation suggests to us that a supersymmetric version of the model can lead to a quite different phenomenology. The results obtained in [14] are summarized in figure 2 of the cited paper. We observe that the described axion is constrained, from SN data, in the region between the dot-dashed line and the dashed one. Also, from terrestrial experiments, the PQ constant f_a must be bigger than 10^4 GeV if the axion mass stays below $\sim 1\,\mathrm{MeV}$ (twice the electron mass), otherwise the limit on f_a should be increased to $f_a \gtrsim$ few $10^5\,\mathrm{GeV}$. We see that for $m_a > 1\,\mathrm{MeV}$, the raising of the dashed line is strongly suppressed. This is because, for an axion heavier than two electrons, the total axion decay width $\Gamma_{\rm tot}$ is dominated by the axion decay into electrons and positrons. Thus the photon flux is strongly suppressed. This phenomena can be improved if the axion mass is increased for a given PQ constant f_a, since $\Gamma(a \to e^+e^-) \sim m_a/f_a^2$. This, as we will show, is exactly what happens in a supersymmetric model, for a certain range of f_a. We will then consider a supersymmetric mirror axion model and this will allow us to strongly enlarge the parameter space for the axion with relevant phenomenological consequences.

Finally, in this model the hierarchy problem between the mirror and the ordinary e.w. scale is solved naturally via the GIFT (Goldstones Instead of Fine Tuning) mechanism [21]. In fact, the Higgs potential has an additional SU(4) symmetry. This *accidental* symmetry is global, but contains the local one $\mathrm{SU}(2) \times \mathrm{SU}(2) \equiv \mathrm{SO}(4)$. When the M-parity is broken, the mirror Higgses acquire a large VEV, while the ordinary ones remain as Goldstone bosons until the supersymmetry breaking is taken into account. The ordinary Higgs VEVs are then naturally of the order of the supersymmetry soft breaking scale [21, 22].

The paper is organized as follows. In section 2, we present a simple mirror supersymmetric axion model and describe the general features and the experimental and astrophysical bounds, showing that there is an allowed parameter space which can be of interest for the future experimental search. In section 3, we carefully study the interesting relations with the physics of the GRBs and with the dynamics of different supernovae. Finally, in section 4, we summarize our results.

2. The mirror Weinberg-Wilczek axion

Let us consider a model based on the gauge symmetry $G \times G'$ where $G = \mathrm{SU}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)$ stands for the standard model of the ordinary particles: the quark and lepton fields $\psi_i = q_i,\, l_i,\, u^c_i,\, d^c_i,\, e^c_i$ (i is a family index) and two Higgs doublets $H_{1,2}$, while $G' = \mathrm{SU}(3)' \times \mathrm{SU}(2)' \times \mathrm{U}(1)'$ is its mirror gauge counterpart with the analogous particle content: the fermions $\psi'_i = q'_i,\, l'_i,\, u^{c\prime}_i,\, d^{c\prime}_i,\, e^{c\prime}_i$ and the Higgses $H'_{1,2}$ [10]. From now on, all fields and quantities of the mirror sector have an apex to distinguish them from the ones belonging to the ordinary world. All fermion fields ψ, ψ' 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

$$\mathcal{L}_{\rm Yuk} = G_U^{ij} u^c_i q_j H_2 + G_D^{ij} d^c_i q_j H_1 + G_E^{ij} e^c_i l_j H_1 + \text{h.c.}\,,$$

$$\mathcal{L}'_{\rm Yuk} = G'^{ij}_U u^{c\prime}_i q'_j H'_2 + G'^{ij}_D d^{c\prime}_i q'_j H'_1 + G'^{ij}_E e^{c\prime}_i l'_j H'_1 + {\rm h.c.}\,, \tag{2.1}$$

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

We further assume that the two sectors have a common Peccei-Quinn symmetry $\mathrm{U}(1)_{\rm PQ}$ realized *à la* Weiberg-Wilczek model. The essential point is then to have a term in the Higgs potential which mixes the Higgses of different kinds. The simplest possibility is:

$$\mathcal{V}_{\rm mix} = -\kappa (H_1 H_2)^\dagger (H'_1 H'_2) + {\rm h.c.}\,, \tag{2.2}$$

where the coupling constant κ should be real due to M-parity.[7]

As far as the M-parity is an exact symmetry, the particle physics should be exactly the same in the two worlds: in particular, the quark mass matrices are identical, $M_{U,D} = M'_{U,D}$, the QCD scales coincide, $\Lambda = \Lambda'$, and the axion couples with 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 θ-terms both in the mirror and ordinary sectors. In such a realization, however, f_a remains order 100 GeV and thus it is excluded on the same phenomenological grounds as the original WW model.

The situation is more interesting when the M-parity is spontaneously broken, as it was suggested in ref. [16], and the electroweak symmetry breaking scale v' in the mirror sector becomes larger than the ordinary one v. Since $\mathrm{U}(1)_{\rm PQ}$ is a common symmetry for the two sectors, the PQ scale is determined by the larger VEV v', i.e. $f_{\rm PQ} \sim (v'/2)\sin 2\beta'$. 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 the mirror sector, while it couples with the ordinary matter as an invisible DFSZ-like axion. This would lead to somewhat different particle physics in the mirror sector and it is not *a priori* clear that the strong CP problem can still be simultaneously fixed in both sectors. However, as it was shown in [14], this is just the case as long as the Yukawa structure in the two sectors is the same. It happens then that the mirror quark masses are scaled linearly with respect to the ordinary ones: $m'_{u,c,t} = \zeta_2 m_{u,c,t}$, $m'_{d,s,b} = \zeta_1 m_{d,s,b}$, where $\zeta_2 = v'_2/v_2 = \zeta(\sin\beta'/\sin\beta)$ and $\zeta_1 = v'_1/v_1 = \zeta(\cos\beta'/\cos\beta)$. At very high energies, $\mu \gg v'$, the strong coupling constants $\alpha_s(\mu)$ and $\alpha'_s(\mu)$ should be equal due to M-parity. Under the renormalization group (RG) evolution they both evolve down in parallel ways until the energy reaches the value of mirror-top mass $m'_t \simeq \zeta_2 m_t$. Below it, α'_s will have a different slope than α_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 α_s these thresholds occur at lower scales, $\mu = m_t, m_b$, etc. Then it is very easy to determine the scale Λ' at which α'_s becomes large, once we know that for the ordinary QCD this happens at $\Lambda \simeq 200$ MeV. In other words, Λ' becomes a function of $\zeta_{1,2}$, and for $v' \gg v$ one could obtain a significant difference between the QCD scales: $\Lambda' > \Lambda$ (see figure 1).

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

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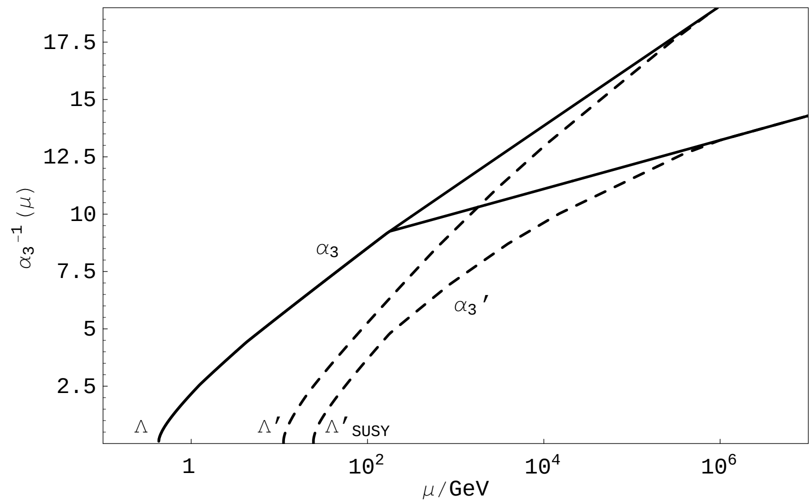

Figure 1: Example of the renormalization group evolution of the strong coupling constants α_s and α'_s (respectively solid and dashed), as a function of the energy scale μ, for $v'/v = 10000$. Supersymmetry is supposed to be broken at the scale $m_s = m_t$.

The value of the QCD constant in the mirror sector depends on the RG evolution of the mirror strong coupling. This, in turn, depends on the matter content of the theory. Therefore in a supersymmetric theory we expect a different result. In general, the relation between the ordinary and mirror QCD scale can be written as

$$\frac{\Lambda'}{\Lambda} = A(\beta, \beta')\,\zeta^{\rho}\,, \tag{2.3}$$

where A is a function of the angles β and β', while ρ is a constant. They both depend on the number of mirror light quarks. We can estimate,[8] for a non supersymmetric theory, $\rho \simeq 0.36$, while in a supersymmetric theory we find the bigger value[9] $\rho \simeq 0.54$.

In the following we will consider the renormalizable lagrangian described by the superpotential:

$$\mathcal{W} = \mathcal{T}(\mathcal{Y}\mathcal{R} - M^2) + \mathcal{S}(\mathcal{H}_1\mathcal{H}_2 + \mathcal{H}_1^{'}\mathcal{H}_2^{'} - \mathcal{R}^2)\,, \tag{2.4}$$

where M is an energy scale $M \gg v$, $\mathcal{H}_1, \mathcal{H}_2, \mathcal{H}_1^{'}, \mathcal{H}_2^{'}$ are the Higgs superfields and $\mathcal{T}, \mathcal{Y}, \mathcal{R}, \mathcal{S}$, are other superfields. $\mathcal{R}$, $\mathcal{Y}$, $\mathcal{S}$ have respectively the PQ charges $\omega_{\mathcal{R}} = -\omega_{\mathcal{Y}} = 1/2$, $\omega_{\mathcal{S}} = -1$, while for the Higgs fields $\omega_1 = \sin^2\beta$, $\omega_2 = \cos^2\beta$ and the same for ω'_1 and ω'_2 with β replaced by β'. $\mathcal{T}$ does not transform under PQ symmetry. For convenience we indicate with $T, Y, R, S, H_i, H_i^{'}$ the scalar components of the superfields $\mathcal{T}, \mathcal{Y}, \mathcal{R}, \mathcal{S}, \mathcal{H}_i, \mathcal{H}_i^{'}$. The second term on the right hand side of (2.4) fixes the VEV pattern of the ordinary and mirror Higgses. The constraint equation is $\langle H_1 H_2\rangle + \langle H'_1 H'_2\rangle \sim \langle R^2\rangle$ where $\langle R^2\rangle$ is fixed in the first term on the right hand side of (2.4): $\langle R\rangle \sim \langle Y\rangle \sim M$. Since the ordinary

[8] We assume, for simplicity, that there are no light quarks in the mirror sector. This is easily verified when $f_a \gtrsim$ few 10^4 GeV.

[9] For this estimation we have assumed that the supersymmetry is broken at the top quark mass. It is also assumed that there are no mirror light quarks, and that the lightest mirror quark is heavier than the ordinary top quark. These assumptions are verified for $f_a \gtrsim$ few 10^6 GeV.

electroweak scale $\langle H\rangle$ is of the same order as the supersymmetry breaking scale m_s, the previous equation fixes $\langle H'\rangle \sim \langle R\rangle \sim M$. Therefore the R VEV fixes the M-parity breaking scale of the theory. Since $M \gg v \sim m_s$ we expect all the mirror fermion masses to be heavier than the ordinary ones and, as a consequence, $\Lambda' > \Lambda$.

The complete M-invariant scalar potential consists in the sum of the F and D terms, plus the supersymmetry breaking contribution:

$$\mathcal{V} = \mathcal{V}_F + \mathcal{V}_D + \mathcal{V}_B\,. \tag{2.5}$$

More explicitly

$$\mathcal{V}_F = \tilde{\mathcal{V}}(H,H^{'}) + |TR|^2 + |YR - M^2|^2 + |TY - 2RS|^2\,, \tag{2.6}$$

where

$$\tilde{\mathcal{V}}(H,H^{'}) = |H_1H_2 + H_1^{'}H_2^{'} - R^2|^2 + S^2(H_1^2 + H_2^2 + H_1^{'2} + H_2^{'2}) + \text{h.c.}\,, \tag{2.7}$$

while:

$$\begin{aligned}\mathcal{V}_D \,=\, & \frac{g^2}{2}\left[(H_1^+H_2)^2 + (H_1^{'+}H_2^{'})^2\right] + \\ & + \frac{g^2+g_y^2}{8}\left[(H_1^+H_1 - H_2^+H_2)^2 + (H_1^{'+}H_1^{'} - H_2^{'+}H_2^{'})^2\right],\end{aligned} \tag{2.8}$$

where g is the coupling constant for isospin and g_y refers to the hypercharge. Finally $\mathcal{V}_B$ contains all the possible soft supersymmetry breaking terms and fixes the scale of the supersymmetry breaking m_s.

In this model the hierarchy problem between the mirror and ordinary Higgs VEVs is solved in a rather natural way by the Pseudo-Goldstone or GIFT mechanism [21]. Namely, while the mirror Higgses get VEVs order $M \gg M_W$, the ordinary Higgses can get the masses (and hence the VEVs) of the order m_s.

More precisely, in the supersymmetric limit $\mathcal{V}_B = 0$, the F-term potential (2.7) has an accidental global symmetry SU(4), larger than the local symmetry $\mathrm{SU}(2)\times\mathrm{SU}(2)'$ acting on the Higgses $H_{1,2}$ and $H'_{1,2}$ respectively. Therefore, if the non-zero VEVs are located on the mirror Higgses, then we obtain $v'_1v'_2 = R^2 \sim M^2$, and the D-term (2.8) gives $\tan\beta' = v'_2/v'_1 = 1$, while the ordinary Higgses $H_{1,2}$ remain as Goldstone superfields.

Then, considering the soft terms $\mathcal{V}_B$, one generates the soft mass terms for $H_{1,2}$ and, as a remarkable property of the GIFT mechanism [21, 22], also supersymmertic μ-term, μH_1H_2 with $\mu \sim m_s$, as far the VEV $\langle S\rangle \sim m_s$ is generated by the soft terms $\mathcal{V}_B$. As a result, one can generate non zero VEVs $v_1, v_2 \sim m_s$ (with $v_1^2 + v_2^2 = v^2 = (247\text{ GeV})^2$), and the parameter $\tan\beta = v_2/v_1$ in general can be different from 1. On the other hand, given that the mirror VEVs are generated at a large scale, $v'_1 = v'_2 \sim M$, the soft terms $\mathcal{V}_B$ cannot significantly shift their values and thus we remain with $\tan\beta^{'} \approx 1$.

Let us discuss now in more detail the axion phenomenology in this model. The PQ scale is (see [1]) $f_{\rm PQ} = \sqrt{\sum_j (v_j\omega_j)^2}$ where ω_j and v_j are respectively the PQ charges and VEVs of the scalar fields. We find $f_{\rm PQ} = (f^2 + f'^{\,2} + \langle S\rangle^2 + 1/2M^2)^{1/2}$ where $f = (v/2)\sin 2\beta$

and $f' = (v'/2)\sin 2\beta'$. In the interesting physical limit $f' \sim M \gg f \sim m_s$ we see that $f_{\rm PQ}$ is essentially equal to f'.

For what concerns the axion mass, its square is in general given by the sum of the right hand side of (1.2) and a similar term, but with V, K, M replaced by V', K', M'. The meaning of these parameters is obvious, $M'$is the mass matrix of the mirror light quarks, and the values $K' \sim \Lambda'^4$ and $V' \sim \Lambda'^3$ characterize the mirror gluon and quark condensates. Assuming, for simplicity, that there are no light quarks in the mirror sector and that $\Lambda' \gg \Lambda$, we find $m_a \simeq \Lambda'^2/f_a$. The axion mass is then essentially driven by the mirror QCD. This mass can be several times bigger than the ordinary WW mass:

$$\frac{m_a}{m_a^{WW}} \simeq \frac{\sin 2\beta}{C}\left(\frac{\Lambda'}{\Lambda}\right)^2, \qquad C = 2\left(\frac{m_u}{(1+z)\Lambda}\right)^{1/2} \sim 0.2. \tag{2.9}$$

For the numerical computation we have used the following values of the quark masses: $m_u = 4\,\text{MeV}$, $m_d = 7\,\text{MeV}$, $m_s = 150\,\text{MeV}$ (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 have taken $V = (250\,\text{MeV})^3$ and $K = (230\,\text{MeV})^4$ and we have assumed that the corresponding parameters in the mirror sector scale as $V'/V = (\Lambda'/\Lambda)^3$ and $K'/K = (\Lambda'/\Lambda)^4$. In figure 2 is shown that Λ' can be significantly larger than Λ, allowing for an axion mass $\sim 1\,\text{MeV}$ when $f_a \sim 10^6 - 10^7\,\text{GeV}$.

The axion couples with fermions and photons in the standard way, with a strength inversely proportional to the PQ constant [1]. Referring to (1.1) we find:

$$c_{ae} = \frac{1}{N}\sin^2\beta\,, \qquad c_{ad} = \frac{1}{N}\sin^2\beta - \frac{1\,z}{1+z}\,, \qquad c_{au} = \frac{1}{N}\cos^2\beta - \frac{1}{1+z}\,. \tag{2.10}$$

For the axion-nucleon couplings we consider $c_{aN} \sim 1$, while the axion interaction with photons is measured by:

$$c_{a\gamma} = \frac{8}{3} - \frac{6K\,\text{Tr}(M^{-1}Q^2)}{V + K\,\text{Tr}(M^{-1})} \simeq \frac{2z}{1+z}\,, \tag{2.11}$$

where the trace is taken over the light quark states (u, d), Q are their electric charges $(+2/3, -1/3)$, and $\alpha = 1/137$ is the fine structure constant. In addition, our axion couples with mirror photons, with the constant:

$$c'_{a\gamma} = \frac{8}{3} - \frac{6K'\text{Tr}(M'^{-1}Q'^2)}{V' + K'\text{Tr}(M'^{-1})}\,, \tag{2.12}$$

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

$$\Gamma(a \to \gamma\gamma) = \frac{g_{a\gamma}^2 m_a^3}{64\pi}\,, \qquad \Gamma(a \to \gamma'\gamma') = \frac{g'^2_{a\gamma} m_a^3}{64\pi}\,, \tag{2.13}$$

where $g_{a\gamma} = c_{a\gamma}(\alpha/2\pi f_a)$ and similarly for $g'_{a\gamma}$.

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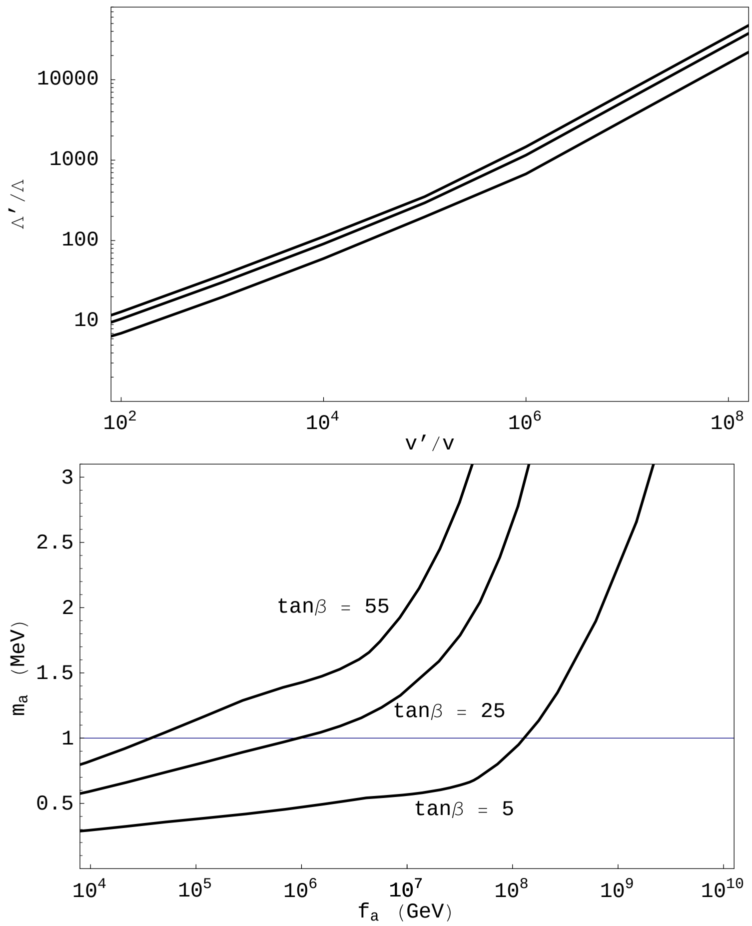

Figure 2: Pictorial representation of the mirror QCD scale and mirror axion mass, in relation to the ordinary value. All the three lines correspond to $\tan\beta' = 1$ and, from the bottom, to $\tan\beta = 5$, $\tan\beta = 25$, $\tan\beta = 55$. Supersymmetry is supposed to be broken at the scale $m_s = m_t$.

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

$$\Gamma(a \to e^+e^-) = \frac{g_{ae}^2 m_a}{8\pi}\sqrt{1 - \frac{4m_e^2}{m_a^2}}\,, \tag{2.14}$$

where $g_{ae} = c_{ae}(m_e/f_a)$.

We present in figure 3 the axion lifetime in the present supersymmetric model. Notice that, as soon as the axion mass turns over 1 MeV its lifetime is strongly suppressed. In fact in this last case its decay width is completely dominated by the $a \to e^+e^-$ channel.

For details of the experimental bounds we refer to ref. [14]. We remark here that the parameter region of interest for us is $f_a \sim 10^5 - 10^7$ GeV, a value safe from terrestrial limits. In this region, the axion mass is about 1 MeV so it is not constrained by the standard astrophysical considerations applicable for the DFSZ axion. In fact, since it is quite heavy, its production rate in the stellar cores, with typical temperatures T up to 10 keV, is suppressed by the exponential factor $\exp(-m_a/T)$. On the other hand, this argument is not applicable to the SN, whose core temperature is several 10 keV. If the

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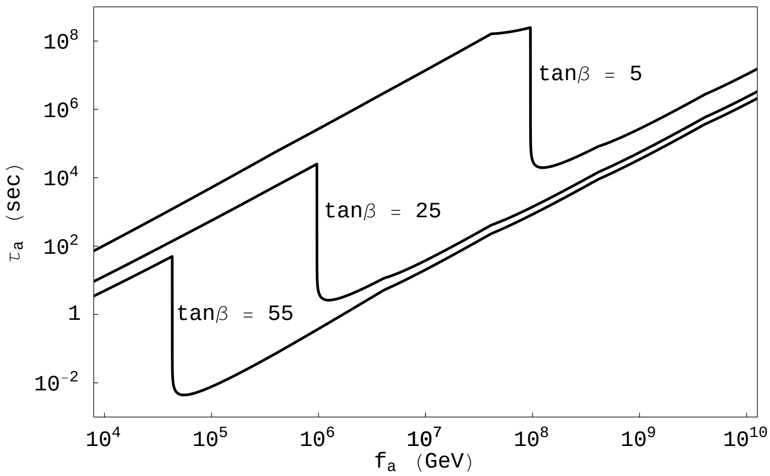

Figure 3: Axion lifetime as a function of f_a for different values of $\tan\beta = 5, 25, 55$. Supersymmetry is supposed to be broken at the scale $m_s = m_t$.

axion-nucleon couplings are large enough, $g_{ap}, g_{an} > 10^{-7}$ (then $f_a < 10^7$ GeV), the axions are strongly trapped in the SN, inside a core of radius $R \simeq 10\,\mathrm{km}$, and they have a thermal distribution. In this case the energy luminosity at $t = 1\mathrm{s}$ can be estimated as $L_a \simeq f_a^{16/11} \times 3 \cdot 10^{50}$ erg/s [23, 24]. Hence, for $f_a <$ a few 10^6 GeV the axion luminosity L_a becomes smaller than a few 10^{51} erg/s, and the total energy and duration of the SN 1987A neutrino burst should not be affected.

The constraint related to the axion decay into ordinary photons is not relevant for an axion with $m_a > 2m_e$, since then, the total axion decay width is almost completely driven by $a \to e^+e^-$ (see figure 3). On the other hand, another constraint emerges due to axion decay into mirror photons. Since 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 decay rate for an axion with energy E into mirror photons is $(m_a/E)\Gamma'$, where $\Gamma' = \Gamma(a \to \gamma'\gamma')$ is given by eq. (2.13). 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 $L'_\gamma \simeq \Gamma' m_a (4\pi R^3/3)(1.2T^3/\pi^2) \sim 10^{-2} g'^2_{a\gamma} m_a^4$. Hence, the condition $L'_\gamma < 10^{51}$ erg/s implies that the decay width Γ', for $m_a \sim 1$ MeV, should not exceed a few s^{-1}. In other terms, using eq. (2.12), we roughly obtain the bound $m_a^2/f_a < 10^{-7}$ MeV.

Finally, a possible cosmological problem must be considered. If the axion contribution to the energy density at the time of the Big Bang Nucleosynthesis (BBN) is very large, this model can disagree with the strongly verified prediction on the primordial Helium abundance. This energy contribution is expressed in terms of effective number of extra neutrinos δN_ν in figure 4.

At the present, deuterium (D) and ^4H_e data seem to indicate that a large number of neutrino species is disfavored. There are several recent papers on this problem (see, e.g. [25] and reference therein), which give different bounds $\delta N_\nu \leq 0.3$, $\delta N_\nu \leq 0.5$ or even $\delta N_\nu \leq 1$ [25]. These limits translate to different lower bounds on the axion mass.

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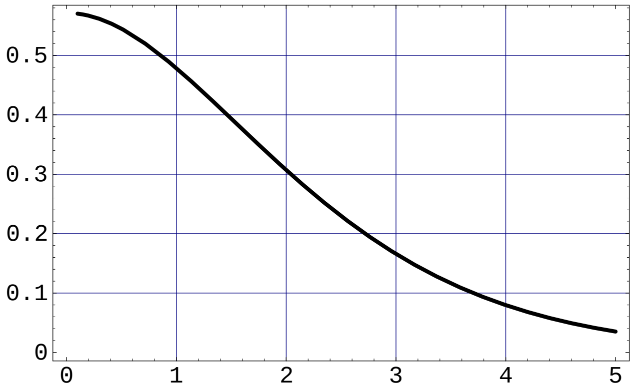

Figure 4: Effective extra neutrino added due to mirror axion as function of axion mass in MeV.

3. Mirror axion: supernovae and gamma ray bursts

The Gamma-Ray Bursts (GRBs) puzzle theorists from many points of view [26]. The most striking feature is that an enormous energy, up to 10^{53-54} erg, is released in a few seconds, in terms of photons with typical energies of several hundred keV. The time-structure of the prompt emission and the afterglow observations well agree with the fireball model [27] in which the GRB originates from the e^+e^- plasma that expands at ultrarelativistic velocities undergoing internal and external shocks. The Lorentz factor of the plasma has to be very large, $\Gamma \sim 10^2$, which requires a very efficient acceleration mechanism. Namely, the e^+e^- plasma should not be contaminated by more massive matter (baryons), and hence the fireball has to be formed in a region of low baryonic density.

In particular, the fireball could be powered via annihilation $\nu\bar{\nu} \to e^+e^-$ of the neutrinos emitted from the dense and hot medium in the accretion disk around a central black hole (BH), which can be formed at the merger of two neutron stars (NS), or a NS and a black hole (BH) [28]. One can consider also the merger of a BH and a white dwarf (WD). In addition, the accretion disk can be formed by the collapse of a rotating massive star, so called failed supernova or collapsar [29]. These objects could potentially provide the necessary energy budget for the GRB. Namely, the typical values of the mass M accreted through a disk, the radial size of a disk R and the accretion time t, can be estimated as:

$$\begin{array}{rllll} \mathrm{NS}+\mathrm{NS}: & M \sim 0.1 M_\odot & R \sim 50\,\mathrm{km}\,, & t \sim 0.1\,\mathrm{s} & \\ \mathrm{NS}+\mathrm{BH}: & M \sim 0.5 M_\odot & R \sim 50\,\mathrm{km}\,, & t \sim 0.1\,\mathrm{s} & \\ \mathrm{Collapsar}: & M \sim 2 M_\odot & R \sim 200\,\mathrm{km}\,, & t \sim 20\,\mathrm{s} & \\ \mathrm{WD}+\mathrm{BH}: & M \sim 1 M_\odot & R \sim 10^4\,\mathrm{km}\,, & t \sim 100\,\mathrm{s}\,. & (3.1) \end{array}$$

However, the problem remains how to transform efficiently enough the available energy into the powerful GRBs. Due to the low efficiency of the $\nu\bar{\nu} \to e^+e^-$ reaction, the models invoking it as a source for the GRBs have serious difficulties in reaching such large photon luminosities. Even though during the collapse of compact objects an energy of $\sim 10^{53}$ erg is normally emitted in terms of neutrinos, they deposit only a small percent of their energy to fireball, and take the rest away. In addition, neutrino annihilation is effective only at small distances, less than 100 km, which are still contaminated by baryon load, and cannot provide a sufficiently large Lorentz factor [28].

In [19], it was proposed that the emission of the light pseudoscalar particles like axions— which can be effectively produced inside the accretion disks and then decay into e^+e^- outside the system— can provide an extremely efficient mechanism for transferring the gravitational energy of the collapsing system into the ultrarelativistic e^+e^- plasma. The advantage of using the decay $a \to e^+e^-$ instead of $\nu\bar{\nu} \to e^+e^-$ annihilation is obvious. First, it is 100 percent efficient, since the decaying axions deposit their energy and momentum entirely into the e^+e^- plasma. And second, the decay can take place in baryon free zones, at distances of 1000 km or more, and so the plasma can get a Lorentz factor $\Gamma \sim 10^2$.

Here we suggest that the described mirror axions can be the required pseudoscalar particles. In the dense and hot medium of the accretion disk it is mainly produced by its bremsstrahlung in the nucleon-nucleon scattering [9]. In order to have efficient production, its mass should not over-exceed the characteristic temperature of the matter, typically a few MeV. Assuming for simplicity the non-degenerate and symmetric baryonic matter, the energy-loss rate per unit mass due to emission of axions is:

$$\epsilon \simeq g_N^2 \rho_{11} T_{\rm MeV}^{7/2} \times 2 \cdot 10^{31}\,{\rm erg\,g^{-1}s^{-1}}\,, \tag{3.2}$$

where ρ_{11} is the disk density in units of $10^{11}\,{\rm g/cm^3}$ and $T_{\rm MeV}$ is the matter temperature in MeV.

Axions are in the free streaming regime, if their mean free path is larger than the accretion disk size ($R_{100} = R/100\,{\rm km}$), which yields [19]:

$$g_N < g_N^{tr}(\rho, T) = 2 \times 10^{-6} \rho_{11}^{-1} T_{\rm MeV}^{1/4} R_{100}^{-1/2}\,. \tag{3.3}$$

Then the total axion luminosity from the accretion disk with a mass M can be roughly estimated as:

$$L \simeq \epsilon M \simeq (10^6 g_N)^2 \rho_{11} T_{\rm MeV}^{7/2} (M/M_\odot) \times 4 \cdot 10^{54}\,{\rm erg\,s^{-1}}\,. \tag{3.4}$$

This value can be so big that the axions can extract all the available energy from the collapsing system with very high efficiency.

We need the emitted axions to decay into e^+e^- outside the disk, in the regions of low baryon density, which correspond to distances of several hundred or thousand km.

For the sake of simplicity, let us fix the parameters as $m_a = 1.5\,{\rm MeV}$, $g_{ae} = 10^{-9}$ and $g_{a\gamma} \sim 10^{-12}\,{\rm MeV}^{-1}$. In this range of parameters the axion lifetime can be well approximated by $\tau \sim \tau(a \to e^+e^-) = 8\pi g_{ae}^{-2}(m_a^2 - 4m_e^2)^{-1/2}$, since the axion decay width into electrons is much larger than the one into photons (cfr. (2.13), (2.14)). Therefore, since the mean decay length of the axion is $D = c\tau E/m_a$, where $E \sim 2T$ is the average energy of the emitted axions, we find:

$$D \simeq (10^9 g_e)^{-2}\, m_{\rm MeV}^{-2}\, E_{\rm MeV} \times 5 \cdot 10^3\,{\rm km}\,, \tag{3.5}$$

where m_{MeV} represents $m_a/{\rm MeV}$ and $E_{\rm MeV} = E/{\rm MeV}$.

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In the view of our mechanism, the short GRBs (duration $\sim 10^{-1}$ s), can be naturally explained by the NS-NS merger, with typical values $M \sim 0.1 M_\odot$ and $t \sim 0.1$ s, the typical density $\rho_{11} \simeq 1$ and temperature $T \simeq 4\,\mathrm{MeV}$. Then the total energy emitted in axions from the accretion torus can be roughly estimated as $\mathcal{E} \simeq Lt \simeq 10^{53}\,\mathrm{erg}$, while the mean decay length is $D \simeq 4 \cdot 10^3\,\mathrm{km}$, much larger than the size of the system ($R \sim 50\,\mathrm{km}$). Thus, the axions decay in the baryon clean zones and deposit their energy entirely to the e^+e^- plasma, which can get a large Lorentz factor and give rise to rather isotropic photon emission with total energies up to $10^{53}\,\mathrm{erg}$. The relative hardness of the photon spectrum in the short GRBs well agrees with this situation. Somewhat more energetic short bursts can be obtained in the case of the NS-BH merger, with $M \sim 0.5 M_\odot$.

On the other hand, the long GRBs can be related to the collapsar. In this case, we need to estimate the fraction of energy deposited from the accretion disk with say $M \sim 2M_\odot$, $R \sim 200\,\mathrm{km}$ and $t \sim 20$, within a cone of the solid angle Ω along the polar axis, where the baryon density is lower and the funneling of the plasma in this direction can take place producing a jet expanding outwards. By taking $\rho \sim 10^{10}\,\mathrm{g/cm^3}$ and $T \sim 2\,\mathrm{MeV}$, we obtain the beamed GRB with the energy release $\mathcal{E}/\Omega \sim Lt \sim 10^{54}\,\mathrm{erg}$. As far as axions decay at large distances, about 1000 km, the Lorentz factor can approach large values, $\Gamma \sim 10^2$. This analysis is supported by the result of the simulation in [30], which shows that if the energy would be transferred to the plasma at distances $\sim 600\,\mathrm{km}$, a successful burst could be obtained with $\Gamma \sim 40$.

The axions can be produced also at the supernova explosion. For $g_N \sim 10^{-6}$ they are in the trapping regime in the collapsing core and are emitted from the axiosphere having a thermal spectrum with a temperature T of a few MeV [24]. Therefore, in total, an energy of a few $\times 10^{51}\,\mathrm{erg}$ can be emitted during the collapse period and subsequent cooling of the proto-neutron star, in terms of axions with the mean energy $E \sim 3T$. The latter undergo the decay into e^+e^- at the distance $D \sim 10^3\,\mathrm{km}$. In this case, the impact of the axion emission crucially depends on the geometrical size of the collapsing star.

In particular, supernovae type-Ib/c result from the core collapse of relatively small stars, where the hydrogen and perhaps also the helium shells are missing. Their radius can be as small as $R \sim 10^4\,\mathrm{km}$, comparable to the axion decay length D. This in turn implies that $\exp(-R/D_a)$ is not very small, and it can be of order 10^{-3} to 1, in which case a reasonable amount of axions can decay outside the mantle producing a fireball. So the weaker GRBs associated with a supernova type-Ib/c could take place, having typical energies up to a few 10^{51} erg.

The SN type-II are associated with large stars, having an extended hydrogen shell ($R > 10^7\,\mathrm{km}$). Thus, the axion decay essentially takes place completely inside the mantle — the fraction of axions decaying outside the star, $\exp(-R/D)$, is essentially zero and thus no GRB can be observed. Indeed, the SN 1987A event did not show any γ signal. On the other hand, the energy of a few $10^{51}\,\mathrm{erg}$ released by axion decay at distances $\sim 1000\,\mathrm{km}$ can help to solve the painful problem of mantle ejection (in the prompt mechanism, shock usually stalls at a distance of a few hundreds km).

Concluding, the axion emission from the collapsing systems and their subsequent conversion into the relativistic plasma via the decay $a \to e^+e^-$ outside these systems could naturally explain a variety of the GRBs. This mechanism suffers no energy deficit and it makes more natural the possibility of the plasma acceleration. In particular, the short GRBs, with timescale $\sim 0.1\,\mathrm{s}$ and total energies up to a few $\times 10^{53}\,\mathrm{erg}$ can originate from the NS-NS or NS-BH mergers, while the collapsar could give rise to the longer GRBs, with $t \sim 10 - 30\,\mathrm{s}$ and $\mathcal{E}/\Omega$ up to a few $10^{54}\,\mathrm{erg}$. The events with $t \sim 100\,\mathrm{s}$ could also be initiated by the BH-WD merger. In later cases, one can expect more baryon dirty fireball. All this well agrees with the observational features of the GRBs. As explained, an interesting possibility is the association of some weak GRBs (with total energy up to $10^{51}\,\mathrm{erg}$ or so) with the supernovae type Ib,c. Considering that the axion emission could also help the supernovae type-II explosion, we see that this model can provide a unified theoretical base for the GRB and SN phenomena. Interestingly, the emission of these axions can also be important for explaining the observed GRB's preceded by supernova explosions, via collapse of usual neutron stars to quark or hybrid stars [31].

4. Conclusions

We have presented a new model of axion for the solution of the strong CP problem. The main feature of this model is the modification of the relation between the axion mass and the PQ constant, with quite interesting phenomenological consequences.

We have hypothesized the existence of another sector of particles and interactions, the *mirror world*, which is an exact copy of our world but with a larger electroweak scale $v' \gg v$. This difference also implies a different dynamics in the QCD sector. The fermion masses are in fact driven by the Higgs VEVs, and the different thresholds for the mirror masses lead to a pole Λ' (in the evolution of the mirror strong coupling constant), that can be significantly higher than the ordinary one Λ. Since the value of Λ in the two sectors depends on the RG evolution equations, and consequently on the matter content of the theory, this behavior can be improved in a supersymmetric model.

The result is that the relation of axion mass to the PQ constant is relaxed, and our particle can be quite heavy, maintaining the weak coupling with matter and photons typical of the invisible axions. This behavior has great benefit on a phenomenological and astrophysical ground. Because of its large mass our axion has in fact no influence on stellar evolution. All astrophysical constraints come from SN explosion and allow its mass to be $m_a \gtrsim \mathrm{MeV}$ with a PQ constant of order 10^5–10^6 GeV.

Such a heavy axion happens to have quite an interesting phenomenology, in particular in relation to the GRB and SN physics. The axion emission from the collapsing systems, and their subsequent conversion into the relativistic plasma via the decay $a \to e^+e^-$ outside these systems, could naturally explain the fireball formation and, consequently, a variety of the GRBs. This mechanism was first proposed using the reaction $\nu\bar{\nu} \to e^+e^-$, instead of the axion decay $a \to e^+e^-$, but the advantage of using the last decay, instead of the neutrino annihilation, is clear. First, it is 100 percent efficient, since the decaying axions deposit their energy and momentum entirely into the e^+e^- plasma. And second, the decay can take place in baryon-free zones, at distances of 1000 km or more, so the plasma can get a Lorentz factor $\Gamma \sim 10^2$.

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Another interesting possibility is the association of some weak GRBs (with total energy up to 10^{51} erg or so) with the supernovae type-Ib,c. The axion mean-free path is a few 10^3 km, so a few of them can reach the supernovae type-I surface and then decay into photons, giving rise to the observed weak GRBs. Finally, the axion emission could also help the supernovae type-II explosion, thus providing a unified theoretical base for the GRB and SN phenomena.

As a final note, the parameter window allowed for our axion is accessible to the axion search in the future reactor and beam dump experiments.

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References

[1] J.E. Kim, *Light pseudoscalars, particle physics and cosmology*, *Phys. Rept.* **150** (1987) 1;
H.-Y. Cheng, *The strong CP problem revisited*, *Phys. Rept.* **158** (1988) 1.

[2] R.D. Peccei and H.R. Quinn, *Constraints imposed by CP conservation in the presence of instantons*, *Phys. Rev.* **D 16** (1977) 1791.

[3] S. Weinberg, *A new light boson?*, *Phys. Rev. Lett.* **40** (1978) 223;
F. Wilczek, *Problem of strong p and t invariance in the presence of instantons*, *Phys. Rev. Lett.* **40** (1978) 279.

[4] A.R. Zhitnitsky, *On possible suppression of the axion hadron interactions. (in russian)*, *Sov. J. Nucl. Phys.* **31** (1980) 260;
M. Dine, W. Fischler and M. Srednicki, *A simple solution to the strong CP problem with a harmless axion*, *Phys. Lett.* **B 104** (1981) 199.

[5] J.E. Kim, *Weak interaction singlet and strong CP invariance*, *Phys. Rev. Lett.* **43** (1979) 103;
M.A. Shifman, A.I. Vainshtein and V.I. Zakharov, *Can confinement ensure natural CP invariance of strong interactions?*, *Nucl. Phys.* **B 166** (1980) 493.

[6] Z.G. Berezhiani, *The weak mixing angles in gauge models with horizontal symmetry — a new approach to quark and lepton masses*, *Phys. Lett.* **B 129** (1983) 99;
Z.G. Berezhiani, *Horizontal symmetry and quark-lepton mass spectrum: the* $\mathrm{SU}(5) \times \mathrm{SU}(3)$*-h model*, *Phys. Lett.* **B 150** (1985) 177;
A.A. Anselm and Z.G. Berezhiani, *Could neutrinos with masses of a few kev be shortlived?*, *Phys. Lett.* **B 162** (1985) 349;
Z.G. Berezhiani and M.Y. Khlopov, *Cosmology of spontaneously broken gauge family symmetry*, *Z. Physik* **C 49** (1991) 73; *The theory of broken gauge symmetry of generations*, *Sov. J. Nucl. Phys.* **51** (1990) 739; *Physical and astrophysical consequences of a violation of the symmetry of generations*, *Sov. J. Nucl. Phys.* **51** (1990) 935;
Z.G. Berezhiani, M.Y. Khlopov and R.R. Khomeriki, *Cosmic nonthermal electromagnetic background from axion decays in the models with low scale of family symmetry breaking*, *Sov. J. Nucl. Phys.* **52** (1990) 65; *On the possible test of quantum flavor dynamics in the searches for rare decays of heavy particles*, *Sov. J. Nucl. Phys.* **52** (1990) 344.

[7] K. Choi, K. Kang and J.E. Kim, *Effects of η' in low-energy axion physics*, *Phys. Lett.* **B 181** (1986) 145.

[8] Particle Data Group collaboration, D.E. Groom et al., *Review of particle physics*, *Eur. Phys. J.* **C 15** (2000) 1.

[9] G.G. Raffelt, *Astrophysical methods to constrain axions and other novel particle phenomena*, *Phys. Rept.* **198** (1990) 1.

[10] T.D. Li, C.N. Yang, *Question of parity conservation in weak interactions*, *Phys. Rev.* **104** (1956) 254;
Y. Kobzarev, L. Okun, I. Pomeranchuk, *On the possibility of experimental observation of mirror particles*, *Yad. Fiz.* **3** (1966) 1154 [*Sov. J. Nucl. Phys.* **5** (1986) 83];
M. Pavsic, *External inversion, internal inversion and reflection invariance*, *Int. J. Theor. Phys.* **9** (1974) 229 [hep-ph/0105344];
R. Foot, H. Lew and R.R. Volkas, *A model with fundamental improper space-time symmetries*, *Phys. Lett.* **B 272** (1991) 67.

[11] B. Holdom, *Two* U(1)*'s and epsilon charge shifts*, *Phys. Lett.* **B 166** (1986) 196;
S.L. Glashow, *Positronium versus the mirror universe*, *Phys. Lett.* **B 167** (1986) 35;
E.D. Carlson and S.L. Glashow, *Nucleosynthesis versus the mirror universe*, *Phys. Lett.* **B 193** (1987) 168;
R. Foot and R.R. Volkas, *Neutrino physics and the mirror world: how exact parity symmetry explains the solar neutrino deficit, the atmospheric neutrino anomaly and the lsnd experiment*, *Phys. Rev.* **D 52** (1995) 6595 [hep-ph/9505359];
Z. Silagadze, *Neutrino mass and the mirror universe*, *Phys. Atom. Nucl.* **60** (1997) 272;
S.I. Blinnikov, *A quest for weak objects and for invisible stars*, astro-ph/9801015;
Gamma-ray bursts produced by mirror stars, astro-ph/9902305;
R. Foot and Z.K. Silagadze, *Supernova explosions, 511 KeV photons, gamma ray bursts and mirror matter*, astro-ph/0404515;
Z. Berezhiani, *Unified picture of the particle and sparticle masses in SUSY GUT*, *Phys. Lett.* **B 417** (1998) 287 [hep-ph/9609342].

[12] Z. Berezhiani, D. Comelli and F.L. Villante, *The early mirror universe: inflation, baryogenesis, nucleosynthesis and dark matter*, *Phys. Lett.* **B 503** (2001) 362 [hep-ph/0008105];
L. Bento and Z. Berezhiani, *Leptogenesis via collisions: the lepton number leaking to the hidden sector*, *Phys. Rev. Lett.* **87** (2001) 231304 [hep-ph/0107281]; *Baryon asymmetry, dark matter and the hidden sector*, *Fortschr. Phys.* **50** (2002) 489; *Baryogenesis: the lepton leaking mechanism*, hep-ph/0111116;
A.Y. Ignatiev and R.R. Volkas, *Mirror dark matter and large scale structure*, *Phys. Rev.* **D 68** (2003) 023518 [hep-ph/0304260];
R. Foot and R.R. Volkas, *Explaining omega(baryon) approx. 0.2 omega(dark) through the synthesis of ordinary matter from mirror matter: a more general analysis*, *Phys. Rev.* **D 69** (2004) 123510 [hep-ph/0402267];
Z. Berezhiani, P. Ciarcelluti, D. Comelli and F.L. Villante, *Structure formation with mirror dark matter: CMB and LSS*, astro-ph/0312605;
P. Ciarcelluti, *Cosmology of the mirror universe*, astro-ph/0312607;
For earlier works see S. Blinnikov and M. Khlopov, *Possible astronomical effects of mirror particles*, *Astron. Zh.* 60 (1983) 632;
M. Khlopov et al., *Observational physics of a mirror world*, *Astron. Zh.* **68** (1991) 42.

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[13] Z. Berezhiani, *Mirror world and its cosmological consequences*, *Int. J. Mod. Phys.* **A 19** (2004) 3775 [hep-ph/0312335];
R. Foot, *Mirror matter-type dark matter*, astro-ph/0407623.

[14] Z. Berezhiani, L. Gianfagna and M. Giannotti, *Strong CP problem and mirror world: the Weinberg-wilczek axion revisited*, *Phys. Lett.* **B 500** (2001) 286 [hep-ph/0009290];
see also V.A. Rubakov, *Sov. Phys. JETP Lett.* **65** (1997) 621.

[15] Z. Berezhiani, L. Gianfagna and M. Giannotti, in preparation.

[16] Z.G. Berezhiani and R.N. Mohapatra, *Reconciling present neutrino puzzles: sterile neutrinos as mirror neutrinos*, *Phys. Rev.* **D 52** (1995) 6607 [hep-ph/9505385];
E.K. Akhmedov, Z.G. Berezhiani and G. Senjanovic, *Planck scale physics and neutrino masses*, *Phys. Rev. Lett.* **69** (1992) 3013 [hep-ph/9205230].

[17] Z.G. Berezhiani, A.D. Dolgov and R.N. Mohapatra, *Asymmetric inflationary reheating and the nature of mirror universe*, *Phys. Lett.* **B 375** (1996) 26 [hep-ph/9511221].

[18] Z. Berezhiani, *Acta Phys. Polon.* **B27** (1996) 1503;
R.N. Mohapatra and V.L. Teplitz, *Structures in the mirror universe*, *Astrophys. J.* **478** (1997) 29 [astro-ph/9603049];
V.S. Berezinsky and A. Vilenkin, *Ultra high energy neutrinos from hidden-sector topological defects*, *Phys. Rev.* **D 62** (2000) 083512 [hep-ph/9908257];
R.N. Mohapatra, S. Nussinov and V.L. Teplitz, *Mirror matter as self interacting dark matter*, *Phys. Rev.* **D 66** (2002) 063002 [hep-ph/0111381];
A.Y. Ignatiev and R.R. Volkas, *Mirror dark matter and large scale structure*, *Phys. Rev.* **D 68** (2003) 023518 [hep-ph/0304260].

[19] Z. Berezhiani and A. Drago, *Gamma ray bursts via emission of axion-like particles*, *Phys. Lett.* **B 473** (2000) 281 [hep-ph/9911333].

[20] A. Loeb, *Are gamma-ray bursts at cosmological distances optically thin?*, *Phys. Rev.* **D 48** (1993) 3419 [astro-ph/9308048];
O. Bertolami, *Gamma-ray bursts, axion emission and string theory dilaton*, *Astropart. Phys.* **11** (1999) 357 [astro-ph/9901184].

[21] K. Inoue, A. Kakuto and H. Takano, *Higgs as (pseudo)goldstone particles*, *Prog. Theor. Phys.* **75** (1986) 664;
A.A. Anselm and A.A. Johansen, *SUSY GUT with automatic doublet-triplet hierarchy*, *Phys. Lett.* **B 200** (1988) 331;
A.A. Anselm, *A supersymmetric theory of grand unification with automatic hierarchy and low-energy physics*, *Sov. Phys. JETP* **67** (1988) 663;
Z. Berezhiani and G. Dvali, *Possible solution of the hierarchy problem in supersymmetrical grand unification theories*, *Bull. Lebedev Phys. Inst.* 5 (1989) 55.

[22] R. Barbieri, G.R. Dvali and M. Moretti, *Back to the doublet — triplet splitting problem*, *Phys. Lett.* **B 312** (1993) 137;
R. Barbieri, G.R. Dvali, A. Strumia, Z. Berezhiani and L.J. Hall, *Flavor in supersymmetric grand unification: a democratic approach*, *Nucl. Phys.* **B 432** (1994) 49 [hep-ph/9405428];
G. Dvali and Q. Shafi, *Gauge hierarchy in* $\mathrm{SU}(3)_c \times \mathrm{SU}(3)_L \times \mathrm{SU}(3)_R$ *AND LOW-ENERGY IMPLICATIONS*, *Phys. Lett.* **B 326** (1994) 258. Z. Berezhiani, C. Csáki and L. Randall, *Could the supersymmetric Higgs particles naturally be pseudogoldstone bosons?*, *Nucl. Phys.* **B 444** (1995) 61 [hep-ph/9501336];

Z. Berezhiani, *SUSY* SU(6) *gift for doublet-triplet splitting and fermion masses*, *Phys. Lett.* **B 355** (1995) 481 [hep-ph/9503366]; *Solving SUSY GUT problems: gauge hierarchy and fermion masses*, hep-ph/9412372; *2/3 splitting in SUSY GUT: higgs as goldstone boson*, hep-ph/9703426;
G.R. Dvali and S. Pokorski, *Role of the anomalous* U(1)*a for the solution of the doublet-triplet splitting problem via the pseudo-goldstone mechanism*, *Phys. Rev. Lett.* **78** (1997) 807 [hep-ph/9610431].

[23] M.S. Turner, *Axions from sn1987a*, *Phys. Rev. Lett.* **60** (1988) 1797.

[24] A. Burrows, M.T. Ressell and M.S. Turner, *Axions and SN1987A: axion trapping*, *Phys. Rev.* **D 42** (1990) 3297.

[25] F.L. Villante and A.D. Dolgov, *Big bang nucleosynthesis and neutrinos*, hep-ph/0310138.

[26] T. Piran, *Gamma-ray bursts and the fireball model*, *Phys. Rept.* **314** (1999) 575 [astro-ph/9810256].

[27] P. Meszaros and M.J. Rees, *Astrophys. J.* **476** (1997) 261;
M. Vietri, *Astrophys. J.* **488** (1997) L105;
E. Waxmann, *GRB afterglow: supporting the cosmological fireball model, constraining parameters, and making predictions*, *Astrophys. J.* **485** (1997) L5.

[28] M. Ruffert and H.T. Janka, *Colliding neutron stars: gravitational waves, neutrino emission and gamma-ray bursts*, astro-ph/9804132.

[29] B. Paczyǹski, *Are gamma-ray bursts in star forming regions?*, *Astrophys. J.* **494** (1998) L45.

[30] M.A. Aloy, E. Mueller, J.M. Ibáñez, J.M. Marti and A. MacFadyen, *Title: relativistic jets from collapsars*, astro-ph/9911098.

[31] Z. Berezhiani, I. Bombaci, A. Drago, F. Frontera and A. Lavagno, *Gamma ray bursts from delayed collapse of neutron stars to quark matter stars*, *Astrophys. J.* **586** (2003) 1250 [astro-ph/0209257];
Z. Berezhiani et al., *Gamma-ray bursts from delayed quark deconfinement phase transition in neutron stars*, *Nucl. Phys.* **113** *(Proc. Suppl.)* (2002) 268.