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Tabulation of astrophysical constraints on axions and Nambu-Goldstone bosons

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Astrophysical constraints on the couplings of light and weakly coupled pseudoscalar particles (axions, Majorons, familons, . . .) from considerations of various stellar objects are summarized. We tabulate the astrophysical bounds on the mass and the decay constant of Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) axions and Kim-Shifman-Vainshtein-Zakharov (KSVZ) axions, on the triplet Majoron vacuum expectation value v_T, and on the familon breaking scale. The lower bound of the Peccei-Quinn breaking scale in the KSVZ model is generally one order of magnitude weaker than that in the DFSZ model. The most stringent limit on $v_T<2$ keV is obtained from considerations of Majoron emission from the cores of neutron stars. Bounds on the strength of the $1/r$ potential mediated by Gelmini-Roncadelli Majorons are also given.

I. INTRODUCTION

Recently there has been a renewal of interest in the search for Nambu-Goldstone bosons. Suggestive evidence of Majorons in neutrinoless $\beta\beta$ decay was uncovered by Avignone *et al.*[1] The Majoron arises when the global $B-L$ symmetry is spontaneously broken. The question of whether a Majoron in fact exists can only be settled by independent experiments. Other examples of weakly coupled light bosons include axions associated with the Peccei-Quinn symmetry,[2] and familons connected with the spontaneous breaking of a global family symmetry.[3]

It is well known that Goldstone bosons can only have derivative couplings to fermionic matter (see, e.g., Ref. 4). Consequently, the flavor-conserving couplings of Goldstone bosons are pseudoscalar and hence they do not mediate the $1/r$ potential but the spin-dependent $1/r^3$ long-range potential. Writing the pseudoscalar interaction of a Goldstone boson with fermionic matter in the form $(m_f/V)\bar{f}i\gamma_5 fG$, where V is generally the scale of the global-symmetry breaking, then $V\gtrsim 10-100$ GeV will ensure the invisibility of the nonrelativistic potential conveyed by Goldstone bosons in laboratory experiments.[4] Nevertheless, Goldstone bosons could play a potentially important role in astrophysics and cosmology. In particular, the interactions of neutral particles with matter can be severely constrained from astrophysical considerations since any light and weakly interacting particles could provide an important stellar energy-loss mechanism.

The purpose of this paper is first to summarize all available astrophysical constraints in the literature. These astrophysical limits apply to any light pseudoscalar bosons. Then we derive the astrophysical bounds on the mass and the decay constant of two different types of invisible axions, on the triplet Majoron vacuum expectation value, and on the familon breaking scale. Constraints on Kim-type axions are in general either not discussed or not treated right in the literature, for which we try to correct in this paper. Owing to the QCD anomaly, the Goldstone boson can in fact mediate the strong CP-violating $1/r$ long-range potential, as pointed out by Chang, Mohapatra, and Nussinov.[5] Bounds on this new force are also discussed.

II. ASTROPHYSICAL BOUNDS

If the weakly coupled Nambu-Goldstone bosons or any light pseudoscalar particles (denoted by ϕ henceforth) exist, they could carry away a large amount of energy from stellar interiors due to their enormous mean free path compared to a typical stellar radius. In order not to destroy the standard scenario of stellar evolution, the couplings of ϕ defined in

$$\mathcal{L}=(g_{\phi ee}\bar{e}i\gamma_5 e+g_{\phi NN}\bar{N}i\gamma_5 N)\phi+C_{\phi\gamma\gamma}\frac{\alpha}{m_e}\phi F\tilde{F} \qquad (1)$$

must be bounded, where α is a finite-structure constant and N denotes the nucleon doublet $\binom{p}{n}$. Since a true Goldstone boson in general does not have anomalous $\phi G\tilde{G}$, $\phi F\tilde{F}$ couplings (G and F are gluonic and electromagnetic fields, respectively; the arion, a Goldstone boson proposed in Ref. 6, however, does have an anomalous electromagnetic coupling) to lowest order in $\bar{\psi}\psi\phi$ interactions, the two-phonon couplings are very suppressed. Therefore, the constraints on the couplings $C_{\phi\gamma\gamma}$ generally apply to axions only.

There are six relevant processes in which light pseudoscalar particles are emitted from the interior of stellar objects (Fig. 1): (1) photoproduction via the Compton-type scattering ($\gamma+e\to\phi+e$) and the Primakoff process ($\gamma+eZ\to\phi+eZ$), (2) electron-nucleus bremsstrahlung, (3) e^+e^- annihilation and bremsstrahlung, (4) neutron-neutron bremsstrahlung, (5) plasma decay, and (6) free-bound ϕ production in $e+Z\to(e,Z)+\phi$, in which a free electron is captured by a heavy ion into an atomic K shell and emits a ϕ (Ref. 7). The coupling $g_{\phi ee}$ is generally determined from the bremsstrahlung process, while the two-photon–axion vertex manifests in the Primakoff amplitude.

In Table I we summarize the astrophysical constraints on the couplings of the light pseudoscalar bosons obtained from various stellar objects. For details, the reader should consult the original papers cited. Several remarks are in order.

(i) Relying on a realistic model for the Sun, a "laboratory" astrophysical bound on the Yukawa coupling $g_{\phi ee} \leq 5.1 \times 10^{-11}$ was recently set with an ultralow-background germanium spectrometer by Avignone *et al.*[8]

(ii) It has been argued that the Primakoff process dominates in the Sun.[9] Raffelt[10] pointed out that this process is actually suppressed due to the Debye-Hückl screening effects in the solar plasma. For the Sun, the Primakoff emission rate is estimated to be reduced by 2 orders of magnitude. As a consequence, the bremsstrahlung processes from electrons, whose importance was first emphasized by Krauss *et al.*,[11] dominate in the Sun and in white dwarfs except for hadronic axions (i.e., axions which have no tree-order couplings to leptons) which only involve in the Primakoff mechanism. For the Sun, the axiorecombination effect contributes only about 4% to the bremsstrahlung rate.[7] For red giants, the Compton rate dominates, but electron-electron bremsstrahlung is also important.

(iii) Frieman *et al.*[12] argued that the usually quoted solar axion bound, obtained by setting the axion luminosity equal to the photon luminosity of the Sun, is arbitrary and inconsistent. A more careful treatment by them yields $g_{\phi ee} < 1.6 \times 10^{-11}$ which improves the previous limit by a factor of 3.

(iv) Better upper bounds on the couplings of ϕ can be derived from considerations of red giants, super giants, white dwarfs, and neutron stars. Some of them are based on observational data combined with the gross features of the stellar evolution theory (e.g., white dwarf cooling times in Ref. 10 and helium-burning lifetimes in Ref. 17); some depend on the details of models of stars and hence are rather model dependent.[13]

(v) The ϕ-nucleon coupling $g_{\phi NN}$ can be constrained from considerations of the cooling rates of neutron stars by assuming that the dominant energy loss mechanism in the core of the star is the ϕ emission from neutron-neutron collisions. This was first considered by Iwamoto[14] and revamped recently by Pantziris and Kang.[15] (A comment on the latter paper was made by Raffelt.[16]) The uncertainty comes from the lack of knowledge of the neutron-star matter equation of state and from neglecting possible internal and external heat sources, and the possibilities of nonthermal magnetospheric emission for the observed x-ray spectra.[8] Further difficulties and uncertainties are discussed in Ref. 17. In principle, the ϕ-quark Yukawa coupling can be extracted from $g_{\phi NN}$ but, as we shall see later, it is obscured by the inaccurately known coupling parameter S.

In Table I astrophysical constraints with a superscript a are obtained without screening-effect corrections and hence may not be used reliably. The best bound on $g_{\phi ee}$ is 1.4×10^{-13} derived from the requirement that helium ignition occur in red giants,[18] and 4.0×10^{-13} relied on the observational evidence of white-dwarf cooling times.[10] The most severe bound on $C_{\phi\gamma\gamma}$, 1.8×10^{-2}, comes from the observational lifetimes of helium-burning stars.[17]

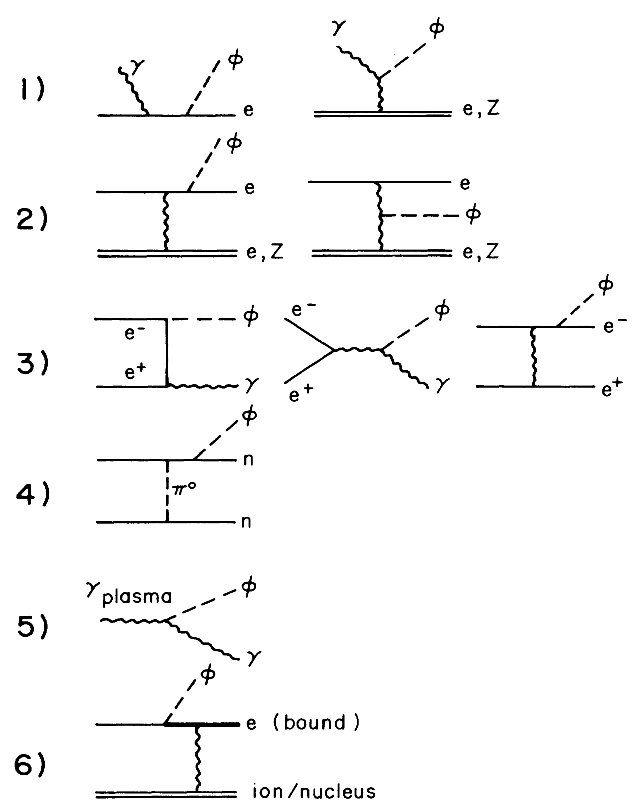

FIG. 1. Six relevant processes in which light pseudoscalar particles emit from the interior of stellar objects.

III. CONSTRAINTS ON INVISIBLE AXIONS

When weak CP violations is "hard" (i.e., CP symmetry is broken by dimensional-four operators), a natural (and the only known) solution to the strong CP problem is to impose a Peccei-Quinn (PQ) invariance[2] (for a review of the strong CP problem, see Refs. 19 and 20). The standard axion[21] associated with the PQ symmetry which is spontaneously broken at the electroweak scale $v = 246$ GeV is not established experimentally. One possibility of accounting for the nonobservation of the standard axion is to bring up the PQ breaking scale, so that the coupling of the axion to fermionic matter is suppressed. In this section we consider two types of invisible axions, namely, the Dine-Fischler-Srednicki-Zhitnitsky[22] (DFSZ) and Kim-Shifman-Vainshtein-Zakharov[23] (KSVZ) axions. An $SU(2) \times U(1)$-singlet scalar field, which develops an arbitrary large vacuum expectation value (VEV), is introduced in the DFSZ model, while weak interactions in the KSVZ model are as usual and the PQ symmetry is implemented by invoking gauge-singlet exotic quarks Q. The KSVZ axion is a type of hadronic axion; namely, it does not couple to leptons at the tree level. In the following, we first sum-

TABLE I. Astrophysical constraints on the couplings of light pseudoscalar bosons to electrons, nucleons, and photons set from various stellar objects.

Sun	$g_{\phi ee} < 5.1\times 10^{-11}$ (Ref. 8)	
	$< 8.5\times 10^{-11}$ (Ref. 9)[a]	$C_{\phi\gamma\gamma} < 7.8\times 10^{-12}$ (Ref. 9)[a]
	$< 4.6\times 10^{-11}$ (Ref. 10)	$< 4.2\times 10^{-11}$ (Ref. 10)
	$< 1.6\times 10^{-11}$ (Ref. 12)	
Red giant	$g_{\phi ee} < 9.0\times 10^{-13}$ (Ref. 9)	
	$< 8.0\times 10^{-13}$ (Ref. 10)	$C_{\phi\gamma\gamma} < 1.9\times 10^{-12}$ (Ref. 10)
	$< 1.4\times 10^{-13}$ (Ref. 18)	$< 2.4\times 10^{-13}$ (Ref. 18)[a]
		$< 1.8\times 10^{-12}$ (Ref. 17)
Super giant	$g_{\phi ee} < 1.5\times 10^{-11}$ (Ref. 9)[a]	$C_{\phi\gamma\gamma} < 8.7\times 10^{-12}$ (Ref. 9)[a]
White dwarf	$g_{\phi ee} < 1.9\times 10^{-11}$ (Ref. 9)[a]	$C_{\phi\gamma\gamma} < 2.3\times 10^{-11}$ (Ref. 9)[a]
	$< 4.0\times 10^{-13}$ (Ref. 10)	
Neutron-star crust	$g_{\phi ee} < 6\text{–}9\times 10^{-13}$ (Ref. 14)	
	$< 5\text{–}6\times 10^{-13}$ (Ref. 15)	
Neutron-star core	$g_{\phi nn} < 4\text{–}6\times 10^{-10}$ (Ref. 14)	
	$< 6\text{–}9\times 10^{-11}$ (Ref. 15)	

[a]This constraint is unduly restrictive since screening effects are not taken into consideration. It should be stressed that the definition of $C_{\phi\gamma\gamma}$ by Fukugita *et al.* (Ref. 9) is different from ours by a factor of 2.

marize the relevant results for both axions (for details, see Secs. 3.4 and 3.5 of Ref. 19), then we discuss the astrophysical bounds on axions.

The DFSZ axion

The couplings g_{aee} and $C_{a\gamma\gamma}$ for DFSZ-type axions are given by

$$g_{aee}=\frac{m_a m_e}{f_\pi m_\pi}\frac{1+z}{N\sqrt{z}}\frac{1}{x^2+1}\ ,$$
$$C_{a\gamma\gamma}=\frac{m_a m_e}{8\pi f_\pi m_\pi}\frac{1+z}{\sqrt{z}}\left[\frac{8}{3}-\frac{2}{3}\frac{4+z}{1+z}\right]\ , \tag{2}$$

where $f_\pi=94$ MeV is the pion decay constant, $z=m_u/m_d=0.568$ (Ref. 24), N is the number of generations, x is the ratio of the VEV's of the two Higgs fields whose neutral components couple to d-type and u-type quarks, respectively, and m_a is the axion mass given by

$$m_a=m_\pi\frac{f_\pi}{f_a}\frac{N\sqrt{z}}{1+z}\left[x+\frac{1}{x}\right]\ . \tag{3}$$

It should be stressed that, as pointed out in Ref. 19, f_a/X_f [X_f is the PQ charge of the fermion and it has been chosen to be x, $1/x$, $1/x$ for u, d, and e, respectively, in Eqs. (2) and (3)] is a physical quantity but the axion decay constant f_a itself is not since the latter depends on the absolute magnitude of PQ charges which is not fixed.

The axion-neutron coupling is of the form

$$g_{ann}=g^0_{aNN}-g^3_{aNN} \tag{4}$$

with

$$g^0_{aNN}=-\frac{m_N}{f_a}\left[\frac{N-1}{2}\right]\left[x+\frac{1}{x}\right]g^0_A\ ,$$
$$g^3_{aNN}=\frac{1}{2}\frac{m_N}{f_a}\left[x-\frac{1}{x}-\left[x+\frac{1}{x}\right]\frac{N(1-z)}{1+z}\right]g^3_A\ , \tag{5}$$

where g^0_A and g^3_A are isoscalar and isovector nucleon form factors, respectively. Neglecting the strange-quark contribution to $\langle N\,|\bar{s}\gamma_\mu\gamma_5 s\,|N\rangle$ (this is equivalent to assuming $m_s\to\infty$), the form factors at $q^2=0$ can be expressed in terms of two parameters F and D:

$$g^0_A=3F-D,\quad g^3_A=F+D\ . \tag{6}$$

A recent fit to neutron β decay and hyperon decay rates gives[25]

$$F+D=1.254,\quad \frac{D}{F+D}=0.61\ . \tag{7}$$

When effects of strange quarks are included, the form factor g^0_A is modified to $g^0_A=(3F-D+2S)/3$ and g^0_{aNN} becomes[26]

$$g^0_{aNN}=-\frac{1}{6}\frac{m_N}{f_a}\left[\left[x+\frac{1}{x}\right](N-1)(3F-D+2S)+(3F-D-S)\frac{2}{x}\right]\ , \tag{8}$$

where S is a new parameter characterizing the flavor-singlet coupling. It is estimated to be $0.1\le S\le 2.2$ from elastic neutrino scattering off nucleons.[27] In the limit $m_s\to\infty$, $3F-D=S$ (Ref. 28).

The KSVZ axion

At low energies, the color anomaly-free axial-vector current for KSVZ axions reads[29]

$$J^a_\mu = f_a \partial_\mu a + \tfrac{1}{2}\bar{Q}\gamma_\mu\gamma_5 Q - \frac{1}{2}\frac{1}{1+z}(\bar{u}\gamma_\mu\gamma_5 u + z\bar{d}\gamma_\mu\gamma_5 d) \ . \tag{9}$$

From this current it follows that

$$m_a = m_\pi \frac{f_\pi}{f_a}\frac{\sqrt{z}}{1+z} \tag{10}$$

and

$$C_{a\gamma\gamma} = \frac{m_a m_e}{8\pi f_\pi m_\pi}\frac{1+z}{\sqrt{z}}\left[6Q_{\rm em}{}^2 - \frac{2}{3}\frac{4+z}{1+z}\right] , \tag{11}$$

where $Q_{\rm em}$ is the charge of the color exotic quark Q. The effective KSVZ axion-electron interaction is generated at the one-loop level induced by $aF\tilde{F}$ couplings[28]

$$g_{aee} = \frac{3\alpha^2}{2\pi}\frac{m_e}{f_a}\left[6Q_{\rm em}{}^2 \ln\frac{f_a}{m_e} - \frac{2}{3}\frac{4+z}{1+z}\ln\frac{\Lambda}{m_e}\right] , \tag{12}$$

where $\Lambda \sim 1$ GeV is the QCD chiral-symmetry-breaking scale[30] and the second term in (12) arises from a-π^0 mixing. Finally, the axion-nucleon couplings are

$$g^0_{aNN} = \frac{1}{2}\frac{m_N}{f_a}g^0_A, \quad g^3_{aNN} = \frac{1}{2}\frac{m_N}{f_a}\frac{1-z}{1+z}g^3_A \ . \tag{13}$$

Bounds on m_a and f_a

In Table II astrophysical constraints are translated into lower bounds on the PQ breaking scale and into upper bounds on the axion mass by the aid of Eqs. (2) and (3) and Eqs. (10)–(12). For DFSZ axions, constraints on g_{aee} give more restrictive limits on m_a and f_a than $C_{a\gamma\gamma}$. For KSVZ axions, no significant bounds can be set from the limits on g_{aee} since the KSVZ axion-electron coupling is very weak. For purpose of illustration in Table II we have chosen $x=1$ for DFSZ axions and $Q_{\rm em} = -\frac{1}{3}$ in the KSVZ model.[31]

In principle, restrictive limits on m_a and f_a can be provided by astrophysical constraints on axion-nucleon couplings. However, as pointed out by Kim,[20] since the parameter S lies in the range $0.1 \le S \le 2.2$, one could have $g_{ann}=0$ for a particular choice of S. Indeed, it is easily seen from Eqs. (4)–(8) and (13) that when $S=0.33$ and 0.17, respectively, in the DFSZ and KSVZ models, $g_{ann}=0$. In such a case, no useful bounds can be induced. Of course, even if $g_{ann}=0$ there is still a substantial axion flux due to proton-proton and proton-neutron bremsstrahlung.[13]

From Table II it is evident that the lower bound of the PQ breaking scale for KSVZ axions axions is generally one order of magnitude weaker than that in the DFSZ model.[32] As remarked in the last section, some of the bounds should not be taken seriously as the screening effects are not corrected. The best limits on m_a and f_a given in Table II are

$$\begin{aligned} &m_a^{\rm DFSZ} < 0.01 \text{ eV}, \quad f_a^{\rm DFSZ} > 3.7\times 10^9 \text{ GeV (helium ignition in red giants)} , \\ &m_a^{\rm DFSZ} < 0.03 \text{ eV}, \quad f_a^{\rm DFSZ} > 1.3\times 10^9 \text{ GeV (white-dwarf cooling times)} , \\ &m_a^{\rm KSVZ} < 0.42 \text{ eV}, \quad f_a^{\rm KSVZ} > 1.4\times 10^7 \text{ GeV (helium-burning lifetimes)} . \end{aligned} \tag{14}$$

IV. CONSTRAINTS ON MAJORONS

A spontaneously broken global symmetry of lepton number will lead to massive Majorana neutrinos and a Nambu-Goldstone boson—the Majoron. This can be accomplished by extending the standard model with an additional gauge-single Higgs field,[33] or SU(2)-triplet Higgs multiplet,[34] or SU(2)-doublet Higgs superpartner.[35] The respective Goldstone bosons are the Chikashige-Mohapatra-Peccei (CMP) Majoron, the Gelmini-Roncadelli (GR) Majoron, and the Aulakh-Mohapatra (AM) Majoron, respectively. In the following we discuss CMP and GR Majorons and their astrophysical bounds.[36]

The CMP Majoron

This model requires the addition of a gauge-singlet Higgs field σ and a right-handed heavy neutrino. The CMP Majoron ϕ_M is the phase field of σ and it does not have a mixing with the Higgs doublet. The effective Majoron-fermion interaction induced at one-loop level is given by[33]

$$\mathcal{L} = \epsilon_f \frac{h_2 G_F}{16\pi^2} m_f m_\nu \bar{f} i\gamma_5 f \phi_M \ , \tag{15}$$

where h_2 is a Yukawa coupling in the Lagrangian, $\epsilon_f = 1$ for e^- and u and -1 for d. From the current upper bound $m_{\nu_e} < 18$ eV (Ref. 37), it is easily seen that the coupling of the CMP Majoron to matter is extremely small. Even if h_2 is of order of unity, it still turns out that

$$g_{\phi_M ee} < 1.7\times 10^{-18} \ . \tag{16}$$

Hence, astrophysical constraints on $g_{\phi ee}$ are trivially satisfied and no useful information is gained as the CMP Majoron is a type of "hadronic" Majoron.[38]

The GR Majoron

An SU(2)-triplet Higgs multiplet Δ is introduced in this model. The wave function of the GR Majoron is primarily the phase field of the neutral component of Δ, but it has a small admixture of the Higgs doublet with the mixing angle $2v_T/v_D$ (v_T and v_D being the VEV's of

TABLE II. Astrophysical constraints on the mass and the decay constant (or the Peccei-Quinn breaking scale) of DFSZ and KSVZ axions. Use of $x=1$ and $Q_{\rm em}=-\frac{1}{3}$ has been made for DFSZ and KSVZ axions, respectively ($N=3$ being assumed). The bounds with an asterisk should not be taken seriously as screening effects have not been included.

Stellar object	DFSZ axion	KSVZ axion
Sun	$m_a<3.7$ eV	
	$f_a>1.0\times10^7$ GeV	
	$m_a<2.9$ eV*	$m_a<1.8$ eV*
	$f_a>1.3\times10^7$ GeV*	$f_a>3.3\times10^6$ GeV*
	$m_a<3.3$ eV	$m_a<9.8$ eV
	$f_a>1.1\times10^7$ GeV	$f_a>6.2\times10^5$ GeV
	$m_a<1.1$ eV	
	$f_a>3.2\times10^7$ GeV	
Red giant	$m_a<0.07$ eV	
	$f_a>5.2\times10^8$ GeV	
	$m_a<0.06$ eV	$m_a<0.45$ eV
	$f_a>6.4\times10^8$ GeV	$f_a>1.4\times10^7$ GeV
	$m_a<0.01$ eV	$m_a<0.06$ eV*
	$f_a>3.7\times10^9$ GeV	$f_a>1.1\times10^8$ GeV*
		$m_a<0.42$ eV
		$f_a>1.4\times10^7$ GeV
Super giant	$m_a<1.0$ eV*	$m_a<2.1$ eV*
	$f_a>3.6\times10^7$ GeV*	$f_a>3.0\times10^6$ GeV*
White dwarf	$m_a<1.4$ eV*	$m_a<5.5$ eV*
	$f_a>2.7\times10^7$ GeV*	$f_a>1.1\times10^6$ GeV*
	$m_a<0.03$ eV	
	$f_a>1.3\times10^9$ GeV	
Neutron-star crust	$m_a<0.04$–0.06 eV	
	$f_a>(5.7$–$8.5)\times10^8$ GeV	
	$m_a<0.03$–0.04 eV	
	$f_a>(0.9$–$1.0)\times10^9$ GeV	
Neutron-star core		

the Higgs triplet and doublet, respectively). As a consequence, the GM Majoron has a tree-level coupling to fermions

$$\mathcal{L}=2\sqrt{2}\,G_F v_T \tilde{\epsilon}_f m_f \bar{f} i\gamma_5 f\phi_M \ , \tag{17}$$

where $\tilde{\epsilon}_f=1$ for d, e^- and -1 for u. It is easy to check that the anomalous coupling of the GR Majoron to photons vanishes.[39]

It is quite straightforward to write down the GR Majoron-nucleon couplings: All we have to do is to replace x/f_a and $1/f_a$ and $1/(xf_a)$ in Eqs. (5) and (8) by

$$\frac{x}{f_a}\to-2\sqrt{2}\,G_F v_T \quad\text{and}\quad \frac{1}{xf_a}\to 2\sqrt{2}\,G_F v_T \ , \tag{18}$$

respectively. Hence,

$$\begin{aligned} g^0_{\phi_M NN}&=-\frac{2\sqrt{2}}{3}G_F m_N v_T(3F-D-S)\ ,\\ g^3_{\phi_M NN}&=-2\sqrt{2}\,G_F m_N v_T(F+D)\ . \end{aligned} \tag{19}$$

From Eq. (4) it follows

$$g_{\phi_M nn}=\frac{2\sqrt{2}}{3}G_F m_N v_T(4D+S)\ . \tag{20}$$

From the relation $g_{\phi_M ee}=2\sqrt{2}\,G_F m_e v_T$, we obtain upper limits on v_T as listed in Table III. Unlike the case of invisible axions, we can however gain useful constraints from the consideration of Majoron emission rates from the neutron-star core. From Eqs. (7) and (20) and the lower bound $S\geq0.1$, we find a very stringent constraint on v_T. The best bound $v_T<2$ keV improves the previous best limit[18] $v_T<9$ keV derived from $g_{\phi ee}$. When S increases, the constraint on v_T becomes even better.

V. CONSTRAINTS ON FAMILONS

The familon is the Goldstone boson associated with the spontaneous breaking of a global family symmetry G_F (i.e., horizontal symmetry). One motivation of reinforcing a global flavor symmetry is attributed to the possibility of embedding the PQ invariance in G_F. It has

TABLE III. Astrophysical constraints on the triplet vacuum expectation value in the GR Majoron model, on the family symmetry-breaking scale F, and on the strength of the $1/r$ potential mediated by GR Majorons. Use of $\bar{\theta}=10^{-10}$ has been made.

Stellar object	GR Majoron	Familon	$1/r$ potential
Sun	$v_T<3.1$ MeV	$F>2.0\times10^7$ GeV	$g_s{}^2<6\times10^{-40}$
	<5.1 MeV	$>1.2\times10^7$ GeV	$<2\times10^{-39}$
	<2.8 MeV	$>2.2\times10^7$ GeV	$<5\times10^{-40}$
	<0.9 MeV	$>6.4\times10^7$ GeV	$<5\times10^{-41}$
Red giant	<55 keV	$>1.1\times10^9$ GeV	$<2\times10^{-43}$
	<49 keV	$>1.3\times10^9$ GeV	$<1\times10^{-43}$
	<9 keV	$>7.3\times10^9$ GeV	$<5\times10^{-45}$
Super giant	<0.9 MeV	$>6.8\times10^7$ GeV	$<5\times10^{-41}$
White dwarf	<1.2 MeV	$>5.4\times10^7$ GeV	$<8\times10^{-41}$
	<24 keV	$>2.6\times10^9$ GeV	$<3\times10^{-44}$
Neutron-star crust	$<37-55$ keV	$>(1.1-1.7)\times10^9$ GeV	$<(1-2)\times10^{-43}$
	$<30-37$ keV	$>(1.7-2.0)\times10^9$ GeV	$<(5-8)\times10^{-43}$
Neutron-star core	$<13-19$ keV		$<(1-2)\times10^{-44}$
	$<1.9-2.8$ keV		$<(2-5)\times10^{-46}$

been shown that if G_F is a maximal family group[3] or a chiral group,[40] the theory will tend to lead to a $U_{PQ}(1)$ symmetry. Moreover, the troublesome domain-wall problem which exists in most axion models can be avoided.[3]

For phenomenological purposes, let us consider the following effective interaction of familons at low energies:

$$\mathcal{L}=\frac{1}{F}[\bar{s}\gamma_\mu(1-\gamma_5)d+\bar{\mu}\gamma_\mu(1-\gamma_5)e+\bar{e}\gamma_\mu(1-\gamma_5)e+\cdots]\partial^\mu\phi_F\ , \tag{21}$$

where ϕ_F refers to the familon species $\phi_F^{sd},\phi_F^{\mu e},\ldots$ and F denotes generically the scale at which flavor symmetry is broken. Applying Dirac equations of motion, we obtain

$$\mathcal{L}\simeq\frac{1}{F}[m_s\bar{s}(1-\gamma_5)d+m_\mu\bar{\mu}(1-\gamma_5)e+2m_e\bar{e}\gamma_5e+\cdots]\phi_F\ . \tag{22}$$

For flavor-diagonal coupling, astrophysical limits on $g_{\phi_F ee}(=2m_e/F)$ put stringent constraints on the breaking scale of F (see Table III). The best bound on F is $F>7\times10^9$ GeV, corresponding to $g_{\phi_F ee}<1.4\times10^{-13}$ (Ref. 18). As axions, no useful information can be induced from the couplings $g_{\phi nn}$.

Restrictive limits on F also arise from the consideration of flavor-changing interactions such as $K^+\to\pi^+\phi_F$ and $\mu\to e\phi_F$. From Eq. (22) we obtain the branching ratio

$$B(K^+\to\pi^+\phi_F)\simeq\frac{(m_K{}^2-m_\pi{}^2)^2}{F^2}\frac{(1-4m_\pi{}^2/m_K{}^2)^{1/2}}{16\pi m_K}\times\frac{1}{\Gamma(K^+\to\text{all})}=\frac{3.3\times10^{13}\ \text{GeV}^2}{F^2}\ . \tag{23}$$

F is constrained to be $\geq 3\times10^{10}$ GeV from the KEK limit,[41] 3.8×10^{-8}. Similarly,[42]

$$B(\mu\to e\phi_F)=\frac{1.6\times10^{14}\ \text{GeV}^2}{F^2} \tag{24}$$

and $F>8\times10^9$ GeV is required from the recent experimental limit[43] on the branching ratio of $\mu\to e\phi_F$, 2.6×10^{-6}.

If the PQ invariance is embedded in flavor symmetry, then the cosmological constraint on the PQ breaking scale[44] applies to familons as well, namely, $F\leq10^{12}$ GeV. This implies that $B(K^+\to\pi^+\phi_F)$ has an upper limit, 3×10^{-11}, which should be testable soon. A measurement of this decay mode to the precision 10^{-10} is now underway at Brookhaven.

VI. CONSTRAINT ON THE $1/r$ POTENTIAL

Although Goldstone bosons have only derivative couplings to fermionic matter, nevertheless, the Goldstone theorem can be violated by anomalies which exist in realistic theories, e.g., by the QCD anomaly. By the aid of the color anomaly, Goldstone bosons can actually mediate the $1/r$ potential.[5] A heuristic argument to see the induced scalar coupling, say $\bar{N}N\phi$, is as follows. The Goldstone boson ϕ first mixes with π^0 and η, which in turn have strong *CP*-violating scalar couplings to nucleons in the presence of the QCD anomaly.

Let us consider a realistic theory that contains a strong *CP*-nonconserving term $\bar{\theta}G\tilde{G}$. This anomaly can be transformed away by a proper chiral rotation of quark fields. This chiral rotation will induce not only the usual strong *CP*-odd Lagrangian, but also a scalar coupling from the interaction $\bar{q}i\gamma_5 q\phi$. In axion models, the scalar interaction reads [Eq. (3.85) of Ref. 19]

$$\tilde{\mathcal{L}}=\frac{\bar{\theta}}{f_a}\left[\frac{m_u m_d}{m_u+m_d}(X_u\bar{u}u+X_d\bar{d}d)-(X_u+X_d)\frac{Nz}{(1+z)^2}(m_d\bar{u}u+m_u\bar{d}d)\right]a\ , \quad (25)$$

where X_u and X_d are the PQ charges of u and d quarks, respectively. How large is $\bar{\theta}$ in the axion model? In the presence of complex higher-dimensional operators due to weak interactions, the minimum of the scalar potential is expected to shift to a nonvanishing $\bar{\theta}$ (see Ref. 19 for a detailed discussion). Unfortunately, a reliable calculation of $\bar{\theta}$ is impossible at present, although some very crude estimate of $\bar{\theta}$ (about $\sim 10^{-14}$) in the Kobayashi-Maskawa model has been made in the past.[45] It was suggested recently[46] that $\bar{\theta}$ lies in the range $10^{-17}-10^{-19}$. If $\bar{\theta}$ is less than 10^{-17}, the axion scalar interaction is undetectable by any practical means.

Next, let us consider the GR Majoron since it has tree-level coupling to quarks. Substituting Eq. (18) into (25) gives the scalar interaction of the GR Majoron with quarks

$$\tilde{\mathcal{L}}=-2\sqrt{2}\,\bar{\theta}\frac{m_u m_d}{m_u+m_d}G_F v_T(\bar{u}u-\bar{d}d)\phi_M\ . \quad (26)$$

Using this Lagrangian we are ready to compute the scalar coupling $g_s\bar{N}\tau_3 N\phi_M$. Noting that the matrix elements $\langle N\,|\,\bar{u}u\,|\,N\rangle,\ldots$ can be determined from the linear first-order baryon mass formulas [Eq. (9.2) of Ref. 24], we find[47]

$$|\,g_s\,|=2\sqrt{2}\,\bar{\theta}G_F v_T\frac{m_u m_d}{m_u+m_d}\,\frac{m_{\Xi^0}+m_{\Xi^-}-m_{\Sigma^+}-m_{\Sigma^-}}{2m_s-m_u-m_d}\ . \quad (27)$$

Bounds on $g_s{}^2$ from various stellar objects are given in Table III. The strongest limit on $g_s{}^2$ is $2\times10^{-26}\bar{\theta}^2$. From neutron electric dipole moment, $\bar{\theta}$ is constrained to be $\le 10^{-9}-10^{-10}$ (Ref. 48). Taking $\bar{\theta}=10^{-10}$ yields $g_s{}^2\le 2\times10^{-46}$, which is just on the verge of the observable experimental range in Eötvös-type experiment,[49] namely, $g_s{}^2$(Eötvös)$\le 10^{-46}$ (Ref. 5). Hence, this new type of force is in principle testable provided that $\bar{\theta}$ is in the vicinity of the present limit.

As pointed out in Ref. 5, a small mass of ϕ can be induced through a quark-loop diagram with scalar coupling vertices. It is estimated to be

$$m_\phi \approx g_s m_q\ ,$$

where m_q is the mass of the loop quark. For $m_q\sim1$ GeV, $m_\phi\sim10^{-14}$ eV and the range of the potential is as large as 10^6 km. The potential of the new force is of the form $V_s=g_s{}^2e^{-m_\phi r}/r$.

VII. SUMMARY AND CONCLUSIONS

In this paper the astrophysical constraints on light and weakly interacting pseudoscalar particles ϕ are summarized. More precisely, astrophysical limits on the couplings of ϕ to electrons, photons, and neutrons are tabulated. These limits are obtained from considerations of the Sun, red giants, super giants, white dwarfs, and neutron stars.

We then apply these astrophysical bounds to invisible axions, Majorons, and familons. For Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) axions, the upper bound on the axion decay constant f_a and the lower bound on the axion mass m_a are determined from g_{aee}, while constraints on $C_{a\gamma\gamma}$ are trivially satisfied. In contrast, stringent limits on Kim-Shifman-Vainshtein-Zakharov (KSVZ) axions are obtained from $C_{a\gamma\gamma}$. No significant information can be extracted from the axion-neutron couplings due to the inaccurately known coupling parameter S. Indeed, $g_{ann}=0$ when $S=0.33$ and 0.17, respectively, in the DFSZ and KSVZ models. The best limits on m_a and f_a are $m_a^{\rm DFSZ}<0.01$ eV, $f_a^{\rm DFSZ}>3.7\times10^9$ GeV, and $m_a^{\rm KSVZ}<0.42$ eV, $f_a^{\rm KSVZ}>1.4\times10^7$ GeV. From Table II it is evident that the lower bound of the PQ breaking scale for KSVZ axions is generally one order of magnitude weaker than that in the DFSZ model.

For Chikashige-Mohapatra-Peccei Majorons, no useful information is gained from astrophysical considerations since they do not have tree-level couplings to fermionic matter. For Gelmini-Roncadelli (GR) Majorons, a severe bound on the triplet vacuum expectation value v_T is obtained from $g_{\phi nn}$. Even with the smallest value of S, 0.1, we find a very restrictive bound $v_T<2$ keV which improves the previous best limit $v_T<9$ keV obtained by Dearborn *et al.*[18]

The family-symmetry-breaking scale F in the familon model is given in Table III. Similar stringent limits on F also arise from the consideration of the flavor-changing reactions such as $K^+\to\pi^+\phi_F$ and $\mu\to e\phi_F$. In either case, the lower bound of F is of order 10^{10} GeV.

Owing to the presence of the QCD anomaly, Goldstone bosons can in fact mediate the $1/r$, but $\bar{\theta}$-dependent long-range, potential ($\bar{\theta}$ being a strong *CP*-violating parameter). In axion models, the axion-scalar interaction is practically unobservable since $\bar{\theta}$ is expected to be very small (about $<10^{-17}$). Therefore, we focus on the scalar coupling of GR Majorons, which is one of the most interesting examples. The strongest bound on the GR-Majoron–scalar coupling is given by $g_s{}^2<2\times10^{-26}\bar{\theta}^2$. For $\bar{\theta}\sim10^{-10}$, the range of the potential is as large as 10^6 km and $g_s{}^2<2\times10^{-46}$, which is just on the verge of the observable range in Eötvös-type experiments. Therefore, this new force is testable provided that the parameter $\bar{\theta}$ is in the vicinity of the present upper limit $10^{-9}-10^{-10}$.

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