Light spinless particle coupled to photons

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A pseudoscalar or scalar particle ϕ that couples to two photons but not to leptons, quarks, and nucleons would have effects in most of the experiments searching for axions, since these are based on the $a\gamma\gamma$ coupling. We examine the laboratory, astrophysical, and cosmological constraints on ϕ and study whether it may constitute a substantial part of the dark matter. We also generalize the ϕ interactions to possess SU(2)×U(1) gauge invariance, and analyze the phenomenological implications.

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I. INTRODUCTION AND MOTIVATION

Peccei-Quinn symmetry [1] is still the most attractive solution to the strong CP problem of QCD. As a consequence of the spontaneous breaking of that symmetry, the axion is born [2]. The axion properties and their phenomenological consequences have been studied in depth (for a review see [3]), and some experiments trying to discover the axion are under way (for a review see [4]). Axions might be constituents of the dark mass of the Universe, and this makes the search experiments even more fascinating.

Almost all experiments so far designed to search for light axions make use of the coupling of the axion to two photons:

$$\mathcal{L} = \frac{1}{8}\, g_{a\gamma\gamma}\, \varepsilon_{\mu\nu\alpha\beta} F^{\mu\nu} F^{\alpha\beta}\, a. \tag{1}$$

The coupling $g_{a\gamma\gamma}$ is proportional to the axion mass m_a:

$$g_{a\gamma\gamma} \approx \frac{\alpha}{2\pi}\, \frac{m_a}{1\text{ eV}}\, 10^{-7}\text{ GeV}^{-1}. \tag{2}$$

An interesting question is whether these dedicated experiments are (1) only sensitive to the axion or (2) could discover another class of particles. The answer is (2). Indeed, any light pseudoscalar particle ϕ coupled to two photons,

$$\mathcal{L} = \frac{1}{8}\, g\, \varepsilon_{\mu\nu\alpha\beta} F^{\mu\nu} F^{\alpha\beta}\, \phi, \tag{3}$$

with a strong enough coupling g would induce a positive signal in some of the axion searches. Of course, a scalar particle coupled to two photons would also be detected in such experiments. To simplify the presentation of the paper, first we will thoroughly discuss the pseudoscalar case. In Sec. VI we will compare the scalar to the pseudoscalar case.

With all this in mind, we have studied the phenomenology and consequences of a light particle ϕ that couples *only* to two photons with strength g.

We consider exclusively this type of interaction, Eq. (3), since the existence of this interaction is the only requirement for having a signal in the axion experiments. By making this assumption, however, we are not generalizing the axion. Our particle ϕ cannot be identified with the axion, since the axion couples to leptons, quarks, and nucleons, and ϕ does not. In this spirit, we will also assume that the coupling g and the mass m of the ϕ particle are not related, as they are for the axion, Eq. (2). In principle, we should consider arbitrary ϕ masses, but since we know that axion experiments are sensitive to very light axions, we will restrict the range of masses; we will only consider $m \le 1$ GeV.

In this paper, we will investigate the laboratory, astrophysical, and cosmological constraints on ϕ. Some of the axion constraints can be directly translated into constraints on g and m, but some cannot. We will also answer the question whether the relic ϕ particles can be, for some range of parameters, the dark matter of the Universe. Another issue we will study is the consequence of adding other couplings to $\gamma\gamma\phi$ in such a way that the full SU(2)×U(1) gauge invariance holds at high energies.

We finish this section with some general remarks. As we said, the motivation that has led us to assume a light particle coupled only to photons is the fact that experiments are sensitive to such a possibility. As far as we know, there is no current theoretical model where such peculiarities arise. In fact, one may even wonder whether it can ever occur. The point is that we know that the coupling of (quasi) Goldstone bosons to photons proceeds through anomalous triangle graphs, where the boson couples to charged particles. This is the situation for the neutral pion and in the axion model. One may argue that, in order to couple ϕ to photons, ϕ has to couple to charged particles, and one may conclude that our assumption of the absence of couplings to matter is inconsistent.

We would like to point out that one may think of scenarios where the only coupling that one may constrain at low energies is g in Eq. (3). We need to introduce particles that are very heavy and carry a new quantum number. We also have to impose that ϕ carries also this quantum number, and that the known leptons and quarks do not. The anomalous graphs with a triangle loop of new particles would then induce the effective coupling of Eq. (3). For heavy enough new particles, the important coupling of ϕ at low energies would be to photons, and all the constraints discussed in this paper do not need to be modified or reconsidered.

A related point is the fact that the effective Lagrangian (3) can only be used for energies $E \ll g^{-1}$. We keep in mind this restriction in all the calculations.

II. LABORATORY CONSTRAINTS: NONDEDICATED EXPERIMENTS

In the next section we will discuss the consequences that the experiments designed to look for axions have for our ϕ particle. Before that discussion, we show in this section that there are other laboratory experiments that also give limits on the parameters of ϕ: its mass m and coupling g.

The quarkonium data can be used to constrain g. The most restrictive limit comes from the absence of the decay

$$\Upsilon(1S) \longrightarrow \phi\ \gamma, \tag{4}$$

where ϕ does not decay in the detector. In Fig. 1(a) we show the diagram giving rise to that decay, and we would like to stress that it is *not* the same diagram used to constrain the axion properties in quarkonium decays $\Upsilon(1S) \to a\ \gamma$ [2]. It is easy to find the branching ratio corresponding to the exotic decay (4):

$$B(\Upsilon(1S) \to \phi\ \gamma) = \frac{1}{8\pi\alpha}\ g^2 m_b^2\ B(\Upsilon(1S) \to e^+e^-), \tag{5}$$

where we have used that $m_b \gg m$. In order to use the experimental data, we need to impose that g be small enough in such a way that ϕ does not decay inside the detector,

$$\frac{p_\phi}{m}\tau > L_{\rm det} \tag{6}$$

where the lifetime τ is given by

$$\tau = \frac{64\pi}{g^2 m^3}. \tag{7}$$

Using the Crystal Ball data [5] we find the limit

$$g \le 4.2 \times 10^{-3}\ \text{GeV}^{-1}. \tag{8}$$

The forbidden region is shown in Fig. 2. We set $L_{\rm det} = 2$ m as a rough estimate; our final result does not strongly depend on the precise value of $L_{\rm det}$.

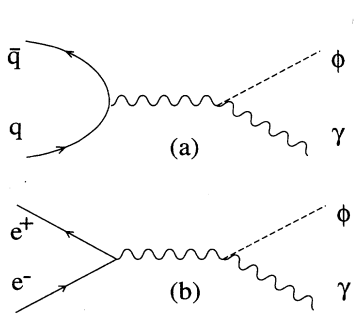

FIG. 1. (a) Diagram of quarkonium decay into ϕ and a photon. (b) Diagram of e^+e^- annihilation into ϕ and a photon.

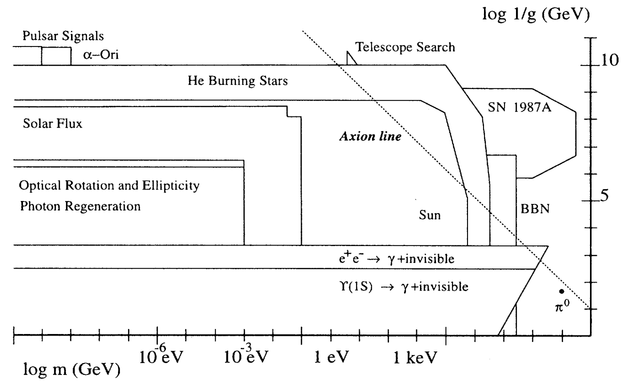

FIG. 2. Excluded regions for the mass m and coupling g of ϕ, coming from laboratory, astrophysical, and cosmological considerations. The line that relates coupling and mass for the axion is shown. We also show the coupling and mass of π^0. The logarithm is to base 10.

Positronium decay into a single photon plus invisible particles lead to limits that are less restrictive than quarkonium decays. The decay $K^+ \to \pi^+ +$ missing and similar were used to constrain axion decays [3] but do not restrict the coupling g of the ϕ particle.

We find, however, another high energy process from which more stringent limits can be extracted. The process is

$$e^+e^- \longrightarrow \phi\ \gamma, \tag{9}$$

and the Feynman diagram is shown in Fig. 1(b). The cross section in the limit $m, m_e \ll \sqrt{s}$ is

$$\sigma(e^+e^- \to \phi\ \gamma) = \frac{\alpha}{24}\ g^2\ f(\theta_{\rm min}), \tag{10}$$

where $f(\theta_{\rm min})$ is a factor less than 1 determined by the angular resolution of the experimental device.

The signature of the process (9) is a single photon observed in the final state, with unbalanced momentum. The standard model process $e^+e^- \to \nu\bar{\nu}\gamma$, with $\nu = \nu_e,\ \nu_\mu$, or ν_τ, has identical signature. Some experiments [6, 7] find a signal consistent with the standard model (SM) contribution, and therefore they are able to place bounds on the anomalous single-photon production. We use this fact to determine a forbidden region in the g versus m plot. In fact, it is the ASP results [6] that lead to the most restrictive limits on the ϕ parameters,

$$g \le 5.5 \times 10^{-4}\ \text{GeV}^{-1}, \tag{11}$$

provided $m \ll 29$ GeV. This is what is shown in Fig. 2.

III. LABORATORY CONSTRAINTS: DEDICATED EXPERIMENTS

In the last few years, some dedicated experiments have been designed to search for axions. Many of these experiments make use of the axion coupling to two photons:

laser experiments [8–10], solar flux detection [11–13], telescope search [14, 15], and microwave cavity experiments [16]. We will find that the first three types of experiments lead to constraints on the properties of ϕ, but not the last one. We will discuss the four types of experiments in turn. In the discussion we will see that some experiments need astrophysical or cosmological assumptions in order to extract consequences from their observations.

The laser experiments consist in the study of laser beam propagation through a transverse magnetic field [8, 10]. The production of real ϕ particles, as in Fig. 3(a), would produce a rotation of the beam polarization, while the emission and absorption of a virtual ϕ, as in Fig. 3(b), would contribute to the ellipticity of the laser beam. Such effects are not observed and their absence implies the constraint [10]

$$g \leq 3.6 \times 10^{-7}\ \mathrm{GeV}^{-1}. \tag{12}$$

This limit is valid provided $m \leq 1$ meV, which is the condition for coherent ϕ production.

A slightly less restrictive limit is obtained when considering the photon regeneration effect that would also occur in the presence of a coupling of ϕ to two photons [9, 10] [see Fig. 3(c)].

The laser experiment limits are shown in Fig. 2. The limits on g based on optical techniques are expected to improve by a factor of about 40 when the PVLAS experiment [17] will take and analyze data. We would like to stress that the constraints from these laser experiments do not depend on astrophysical or cosmological assumptions.

The search for solar axions [12, 13] is based on Sikivie's idea [11] that axions produced in the Sun can be converted into x rays in a static magnetic field. One can use the absence of such signal to place a limit on the coupling g of ϕ to two photons. The Sun produces ϕ particles through the Primakoff effect [see Fig. 4(a)], and these particles reconvert to x rays by means of the inverse process shown in Fig. 4(b). The limit is [13]

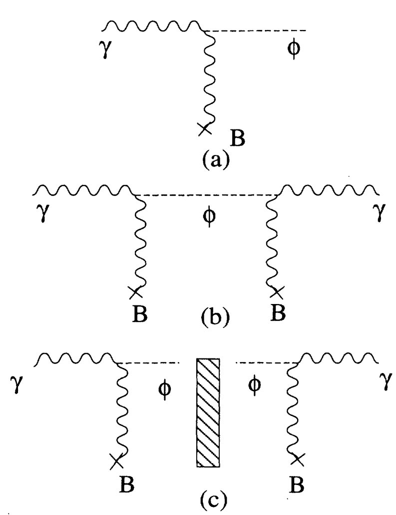

FIG. 3. Diagrams contributing to (a) rotation, (b) ellipticity, and (c) photon regeneration effects of the laser experiments.

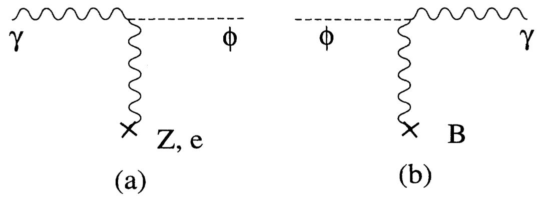

FIG. 4. (a) ϕ are produced in the Sun and (b) reconverted in x rays.

$$g \leq 3.6 \times 10^{-9}\ \mathrm{GeV}^{-1}. \tag{13}$$

Again, this limit is only true when the conversion process is coherent in the magnetic telescope. Coherence is preserved for $m \leq 0.03$ eV. When $0.03\ \mathrm{eV} \leq m \leq 0.11$ eV a slightly less restrictive limit is obtained [13]. The constraints on the ϕ properties obtained from the solar flux experiment are plotted in Fig. 2. To reach the numerical limit on g, Eq. (13), one relies on the calculated axion luminosity from the Sun [18].

The axion search program in BINP (Novosibirsk) will use a similar procedure to probe the coupling g [19].

We turn now our attention to the telescope search experiment [14, 15], which was motivated by the axion window in the eV range [20]. The unsuccessful search can be used to constrain the properties of ϕ when m is also in the eV range. The assumptions one needs are reasonable [14]: ϕ particles have been in thermal equilibrium in the early Universe due to their photon interaction, leaving a relic density after decoupling. These cosmological ϕ particles will be found in galactic clusters, since they fall with baryons into the potential wells that arose in the evolution of the early Universe. The decay $\phi \to \gamma\gamma$, provided ϕ has a mass in the range 3–8 eV, should produce detectable lines when observing rich clusters of galaxies. Bershady, Ressell, and Turner [15] have reported a null signal for such line of radiation, which excludes the region indicated in our Fig. 2.

We finally consider the microwave cavity experiments [16], which search for the signal of the conversion of halo axions, with $E \sim 10^{-5}$ eV, in a magnetic field following the ideas of Ref. [11] [see Fig. 4(b)]. The search, up to now, has been unsuccessful. To get strong limits on the coupling of axions to photons, it is assumed that axions form the dark galactic halo. While this assumption is tenable in the case of axions, it is far from being realistic when considering ϕ, as we now discuss.

For the range of g and of m of interest in microwave experiments, relic ϕ's are relativistic. As we will see in Sec. V, today's temperature of ϕ is bounded by $T_\phi \leq 10^{-4}$ eV. It follows that the contribution of ϕ to the total energy of the Universe, normalized to the critical density, is

$$\Omega_\phi = \frac{\rho_\phi}{\rho_{\mathrm{crit}}} \leq 10^{-5}; \tag{14}$$

i.e., the abundance of ϕ particles fails by several orders of magnitude to account for the halo dark matter. We conclude that the microwave experiments give very poor constraints on the properties of ϕ.

IV. ASTROPHYSICAL CONSTRAINTS

Stellar interiors can produce ϕ particles, provided ϕ is light enough, and these particles can escape almost freely, provided ϕ interacts weakly enough with radiation. This effect is a potentially important energy loss mechanism. The confrontation of astrophysical observations with the modifications caused by nonstandard energy losses on star evolution leads to astrophysical constraints on the ϕ properties.

In the case of the Sun and red giant evolution it turns out that the detailed studies that have been done in the literature for the axion can be used without changes to get limits on the coupling g and the mass m of ϕ. To be more specific, the so-called hadronic or Kim-Shifman-Vainshtein-Zhakharov (KSVZ) axion is dominantly produced in the Sun and in a red giant interior through the Primakoff process. It is also through this process that ϕ is produced, so that we can use the axion limits. We will borrow the numerical results found by Raffelt and Dearborn [21]. The reader interested in more details can consult Raffelt's review [22], where a full list of references can be found.

To get a rigorous limit on g from the Sun we use the numerical studies of [21], which follow the evolution of a star with $M = M_\odot$ that loses energy with the emission of light particles produced by the Primakoff effect. For reasonable values of the presolar helium concentration the present luminosity L of the star corresponds to the Sun luminosity $L = L_\odot$ as long as

$$g \leq 2.5 \times 10^{-9}\ \text{GeV}^{-1}. \tag{15}$$

The limit is valid when the ϕ mass is less than the central temperature of the Sun, $T \approx 1$ keV.

When m is larger than the solar core temperature one has to take into account that the number density of photons with enough energy to produce particles ϕ quickly decreases when m increases. Therefore, for 1 keV $\leq m \leq$ 60 keV, the corresponding bound on g is less stringent than for smaller masses, and for masses larger than 60 keV there is no bound. This effect can be seen in our results plotted in Fig. 2.

Also, when g is large enough, ϕ particles are trapped in the Sun mainly due to their decay into two photons if m is larger than a few keV, or to rescattering through the inverse Primakoff effect if m is smaller. To bound g, we have calculated the effective mean opacity of ϕ and compared it with the photon mean opacity, along the lines described in Ref. [23]. It turns out that ϕ would be allowed for quite a strong coupling g, and in fact much of the excluded region in this trapping regime was already ruled out by the laboratory experiments presented in Sec. II; when there is such overlap we do not show the astrophysical bound in Fig. 2.

Raffelt and Dearborn [21] have also followed the evolution of stars with properties similar to the stars of the open cluster M67, paying attention to plasmon mass, screening, and degeneracy effects. The helium burning lifetime is shortened when ϕ is emitted from the stellar interior. The star count in M67 and other open clusters would be in sharp conflict with the effect of energy loss in form of ϕ particles unless

$$g \leq 1 \times 10^{-10}\ \text{GeV}^{-1} \tag{16}$$

for masses less than $T \approx 10$ keV. For 10 keV$\leq m \leq$ 300 keV, the suppression in the number density of photons with enough energy to produce ϕ particles makes the bound on g less and less stringent until it disappears.

Also, ϕ would be trapped in the red giant interiors for a relatively large value of the coupling g. Again, we use the procedure presented in Ref. [23] to limit g, as we did in the solar case. In Fig. 2 we display these limits when there is no overlapping with the limits coming from laboratory experiments.

Finally, we will consider the constraints coming from the observation of the neutrino pulse from the Supernova 1987A. The limits on g are less stringent than the ones obtained from the Sun and red giants, and for this reason an estimation of the limit will be enough for our purposes. Let us start assuming that m is smaller than the temperature of the supernova core. In this core, ϕ production is mainly due to the processes (see Fig. 5)

$$e\gamma \longrightarrow e\phi, \tag{17}$$

$$p\gamma \longrightarrow p\phi, \tag{18}$$

$$pn \longrightarrow pn\gamma\phi. \tag{19}$$

We stress that these are *not* the main processes for axion production in supernovae.

The cross sections for these processes can be estimated as

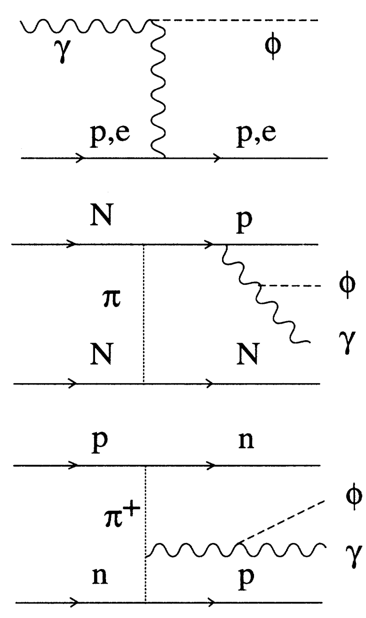

FIG. 5. Processes contributing to ϕ emission in the supernova core. In the second diagram, N stands for n or p. There are also similar diagrams not displayed where the photon is attached to an initial p.

$$\sigma(e\gamma \to e\phi) \approx \sigma(p\gamma \to p\phi) \tag{20}$$
$$\approx \alpha\, g^2, \tag{21}$$
$$\sigma(pn \to pn\gamma\phi) \approx \frac{\alpha}{\pi^3}\, \sigma(pn \to pn)\, g^2\, T^2, \tag{22}$$

where $T \approx 50$ MeV is the temperature of the supernova core and $\sigma(pn \to pn) \simeq 100$ mb is the cross section for the process $pn \to pn$. These cross sections can be used in turn to estimate the rate of energy drain:

$$\dot{E}(e\gamma \to e\phi) \approx \dot{E}(p\gamma \to p\phi) \tag{23}$$
$$\approx V\, n_p\, n_\gamma\, \sigma(p\gamma \to p\phi)\, T, \tag{24}$$
$$\dot{E}(pn \to pn\gamma\phi) \approx V\, n_p\, n_n\, \sigma(pn \to pn\gamma\phi)\, \sqrt{\frac{T}{m_N}}\, T, \tag{25}$$

Here $V \simeq 4 \times 10^{18}$ cm^3 stands for the volume of the supernova core and m_N is the nucleon mass. The number densities we will adopt are $n_n \simeq 7 \times 10^{38}$ cm^{-3} and $n_p \simeq n_e \simeq 3 \times 10^{38}$ cm^{-3}. The constraint from the neutrino signal emitted by SN 1987A is [20, 22, 24] $\dot{E} \leq 10^{52}$ erg/sec, or, in terms of g,

$$g \leq 10^{-9}\ \mathrm{GeV}^{-1}, \tag{26}$$

valid for $m \leq 50$ MeV. As we already said, the SN 1987A limit is less stringent than other astrophysical limits. However, the range of m for which the limit holds is larger. For 50 MeV $\leq m \leq$ 500 MeV this bound becomes smaller due to the low number of photons with enough energy to produce ϕ particles.

To find the limits in the trapping regime, we follow Ref. [23] to calculate the effective ϕ opacity κ_ϕ, dominated by decay, and impose [25]

$$\kappa_\phi > \kappa_\nu \simeq 8 \times 10^{-17}\mathrm{cm}^2/\mathrm{g}. \tag{27}$$

All these constraints are shown in Fig. 2.

One can use data on pulsar signals to probe the g coupling, provided the mass is very small, $m \leq 10^{-10}$ eV [26]. Magnetic fields of pulsars create a ϕ background, and pulsar signals propagating through this background show a time lag between different modes of polarization. The limit from present data is $g \lesssim 2 \times 10^{-11}$ GeV^{-1}.

A similar limit, $g \leq 2.5 \times 10^{-11}$ GeV^{-1}, also valid for small masses, $m \leq 10^{-9}$ eV, has been obtained by Carlson [27]. He has studied x-ray conversion of ϕ produced in stellar cores and found the limit on g using HEAO1 satellite data on α-Ori x-ray emission.

V. COSMOLOGY OF ϕ

In this section, we will study the cosmological evolution of the ϕ species until the present time. We will discuss to which extent ϕ can be the dark matter of the Universe, and also the constraints from primordial nucleosynthesis on the ϕ parameters.

In the early Universe, ϕ is in equilibrium due to the processes

$$A\phi \longleftrightarrow A\gamma \quad \text{and changing } A \text{ by } \bar{A},$$
$$\gamma\phi \longleftrightarrow A\bar{A}. \tag{28}$$

Here A stands for any particle with electric charge $Q_A \neq 0$ present in the primordial plasma. The interaction rate can be written as

$$\Gamma \simeq \frac{\alpha}{24}\, g^2\, n, \tag{29}$$

with n given by

$$n \simeq \frac{\zeta(3)}{\pi^2}\, f(T)\, T^3, \tag{30}$$

where

$$f(T) \approx \sum_A g_A Q_A^2 \tag{31}$$

and g_A is related to the internal degrees of freedom of A. In this formula the sum runs over all the charged particles with $m_A < T$. We are making the usual approximation of considering only the relativistic degrees of freedom, whose contribution is dominant (the contribution of the nonrelativistic degrees of freedom is exponentially suppressed).

The expansion of the Universe is characterized by the rate

$$H = 1.66\, \sqrt{g_*(T)}\, \frac{T^2}{m_{\mathrm{Pl}}}, \tag{32}$$

where $g_*(T)$ is the relativistic degrees of freedom contributing to the energy density. We follow the notation of [20]. The ϕ species decouples when both rates meet

$$H(T_f) = \Gamma(T_f). \tag{33}$$

The freeze-out temperature is a (decreasing) function of g, and also depends on m. It can be calculated numerically, but before we do it, we present two general peculiarities that will help to understand the cosmological evolution of ϕ.

First, we can show that T_f is bounded by $T_f \geq 0.2$ MeV. Using that for $T_f \leq 0.2$ MeV the only particle contributing to the processes (28) is $A = e$, that $n_e \sim 10^{-10}T^3$, and that $g_* \simeq 3.4$ we can deduce, from (33),

$$T_f \sim 0.2\, \left(\frac{0.04\ \mathrm{GeV}^{-1}}{g}\right)^2 \mathrm{MeV}. \tag{34}$$

Since laboratory experiments exclude $g \geq 0.04$ GeV^{-1} for $m \leq 10$ MeV, we conclude that $T_f \geq 0.2$ MeV. The ϕ freeze out occurs before e^+e^- annihilation.

Second, we can show that ϕ can only be a hot relic. Writing Γ as a function of the ϕ lifetime τ and $H = (2t)^{-1}$ [radiation-dominated (RD) universe and we set $\Omega = 1$, but other values of Ω lead to the same conclusion], the freeze-out condition (33) can be written as

$$\frac{3T_f}{m} \approx \left(\frac{\tau}{f(T_f)\, t_f}\right)^{1/3}. \tag{35}$$

We now force that $\tau \geq t_0 \sim 10^{17}$ sec, since we want ϕ to survive until the present time. Also, since $T_f \geq 0.2$

MeV, we have that $t_f \leq 1$ sec. It follows that

$$\frac{3T_f}{m} \geq \left(\frac{10^{17}}{f(T_f)}\right)^{1/3} \gg 1; \tag{36}$$

i.e., if ϕ is a relic, it is a hot relic. [These conclusions are based on the processes (28). One can show that the process $\phi \to \gamma\gamma$ does not modify them.]

After showing these two general features, we turn our attention to the ϕ abundance. Specifically, we would like to elaborate on the question whether ϕ can be a substantial component of the dark matter of the Universe.

If ϕ is today a relativistic particle, it has a very small abundance. The density would be less than the relic photon density

$$\Omega_\phi h^2 < \Omega_\gamma h^2 \simeq 2.6 \times 10^{-5}. \tag{37}$$

Thus we will concentrate on the case that ϕ is a nonrelativistic particle, $m \gtrsim 10^{-4}$ eV. The abundance is

$$\Omega_\phi h^2 = 7.8 \times 10^{-2} \frac{m(\text{eV})}{g_{*S}(T_f)}, \tag{38}$$

where $g_{*S}(T)$ is the effective degrees of freedom that contribute to the entropy of the Universe at a temperature T [20]. The calculation of Ω_ϕ when 0.2 MeV $\leq T_f \leq$ 300 GeV is straightforward since we know the functions $g_*(T)$, $g_{*S}(T)$ [28] and $f(T)$. We first calculate T_f solving Eq. (33) and then we introduce $g_{*S}(T)$ in Eq. (38) to get the ϕ contribution to the energy density of the Universe. Freeze out at $T_f \sim 300$ GeV corresponds to a coupling $g \sim 10^{-9}$ GeV^{-1}. We would like to know the abundances for smaller values of g, say, until $g \sim 10^{-12}$, that would correspond to a freeze-out temperature $T_f \leq 10^9$ GeV. We have to extrapolate $g_*(T)$, $g_{*S}(T)$, and $f(T)$ for $10^9 \geq T \geq 300$ GeV. We will consider three plausible scenarios.

(A) The SM desert: there are no particles with masses above ~ 300 GeV so that $g_*(T)$, $g_{*S}(T)$, and $f(T)$ are constant at high energies.

(B) The MSSM desert: one finds the supersymmetric (SUSY) partners at ~ 300 GeV so that the functions $g_*(T)$, $g_{*S}(T)$, and $f(T)$ increase in one step, and stay flat at high temperatures.

(C) Power-law extrapolation: the function $g_*(T)$ for 200 MeV $\leq T \leq$ 300 GeV can be approximated by a power-law function $g_*(T) \sim T^{0.1}$. We can extrapolate $g_*(T)$ at high temperatures by letting it increase according to this power-law function. The functions $g_{*S}(T)$ and $f(T)$ are extrapolated in a similar way.

It should be clear that we ignore how to extrapolate $g_*(T)$, since that depends on the unknown new physics and new particles above the weak scale. We will in turn consider the three scenarios described above, calculate Ω_ϕ, and see the consequences. In a sense scenarios (A) and (C) represent two extreme possibilities: $g_*(T)$ does not increase at all in (A), while it increases at high temperatures at the same pace as it does at low temperatures in (C). Case (B) is a model-motivated scenario. In Fig. 6 we present our results in the form of lines of constant $\Omega_\phi h^2$. Two possibilities are obviously interesting: ϕ being the cosmological dark matter of a critical Universe or ϕ being only the dark matter in galactic halos. For the first possibility to make sense we should have $\Omega_\phi = 1$; for the second a necessary condition is $\Omega_\phi \geq 0.02h^{-1}$. Since $0.5 \leq h \leq 1.0$, the interesting range is $0.01 \leq \Omega_\phi h^2 \leq 1$. The two extremes of the range are plotted in Fig. 6, where also the requirement that ϕ has survived until the present time, $\tau > H_0^{-1}$, is plotted.[1] Both conditions $\Omega_\phi h^2 = 1$ and $\Omega_\phi h^2 = 0.01$ are worked out and plotted for each of the three scenarios (A), (B), and (C) presented above.

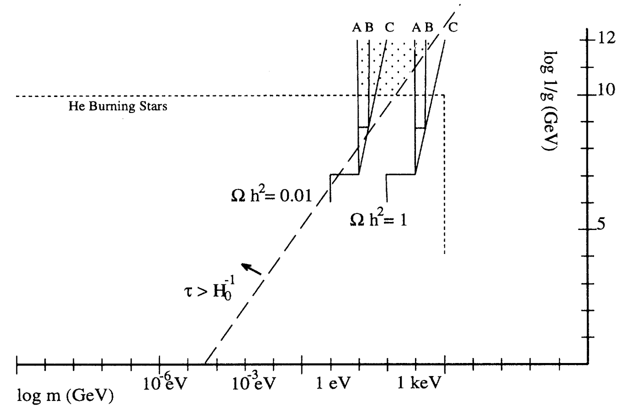

FIG. 6. The lines $\Omega_\phi h^2 = 0.01$ and $\Omega_\phi h^2 = 1$ are represented as a function of g and m. Each condition is calculated in three different cases, according to how we extrapolate the degrees of freedom in the early Universe for $10^9 \geq T \geq 300$ GeV. In (A) there are no more excited degrees of freedom, while in (C) they increase at the same rate as a function of T as they do for $T \leq 300$ GeV. In (B) they increase at $T \sim 300$ GeV as in a SUSY theory and stay constant at higher temperatures. The condition that ϕ has survived until the present time, $\tau > H_0^{-1}$, is also shown. The logarithm is to base 10.

The dotted region in Fig. 6 is the range of couplings and masses such that ϕ is the dark matter. We see that ϕ must have at least a mass $m \geq 10$ eV. The maximum mass that is interesting for dark matter depends on how small we allow g to be. Let us impose that $g \geq 10^{-12}$ GeV^{-1}. Then we have $m \leq 10$ keV. For 10 eV $\leq m \leq$ 100 eV, ϕ could be the galactic dark matter; for 1 keV $\leq m \leq$ 10 keV, ϕ is the cosmological dark matter.

The interesting ϕ masses for dark matter are unfortunately away from the mass range where undergoing experiments could be able to detect ϕ. As we mention in Sec. III, laser experiments work for $m \leq 10^{-3}$ eV, solar axion detection is possible for $m \leq 10^{-1}$ eV, and microwave experiments are restricted to $E \sim 10^{-5}$ eV. It is clear that these experiments are not sensitive to a

[1] The requirement should be $\tau > t_{\text{Univ}}$. The lifetime of the Universe, t_{Univ}, is proportional to H_0^{-1} with the proportionality factor depending on Ω. Also, H_0 is known up to a factor of 2. Our conclusions do not depend on these details.

particle ϕ coupling only to two photons and constituting the dark matter. The telescope search experiment has 3 eV $\leq m \leq$ 8 eV, which is near the mass range relevant for dark matter. Since the calculation of the ϕ abundances has some inherent uncertainties, we conclude that the telescope search may be sensitive to a ϕ that forms the galactic dark halo.

A last point relevant for ϕ being dark matter is the fact that ϕ would be hot dark matter, as was shown at the beginning of this section. Formation of the large structures of the Universe seems, however, to require a large proportion of cold dark matter, so that ϕ would not have a role in structure formation.

We finally discuss the constraint on the ϕ parameters that arises when considering the big bang nucleosynthesis (BBN) of light elements. The comparison of the observed primordial abundances with the predictions of the standard model of the early Universe places a bound on the expansion rate at $T \approx 1$ MeV. In terms of relativistic degrees of freedom this bound reads [29]

$$\Delta g_*(T = 1 \text{ MeV}) \leq 0.5. \tag{39}$$

The contribution of ϕ to the effective degrees of freedom at the nucleosynthesis epoch, provided $m \ll 1$ MeV, is given by

$$\Delta g_*^\phi(T) = \left(\frac{T_\phi}{T}\right)^4, \tag{40}$$

where $T \simeq 1$ MeV. We allow for the case that after the ϕ decoupling and before the nucleosynthesis epoch there has been pair annihilations to photons, in such a way that at $T \simeq 1$ MeV one has $T_\phi \neq T$. How different the two temperatures are is a function of T_d, the decoupling temperature. Entropy conservation gives the relation [20, 28]

$$\frac{T_\phi}{T} = \left(\frac{g_{*S}(T)}{g_{*S}(T_d)}\right)^{1/3}. \tag{41}$$

Since $g_{*S}(1 \text{ MeV}) = 10.75$, the bound (39) implies $g_{*S}(T_d) > 17.4$ which in turn forces $T_d > 200$ MeV.

This result applies when $m \ll 1$ MeV. When $m \sim 1$ MeV, Eqs. (40) and (41) have to be modified but still one can restrict g when

$$m \leq 2.6 \text{ MeV}. \tag{42}$$

In this mass range we obtain the condition

$$\Gamma(T) < H(T) \text{ for } 1 \text{ MeV} \leq T \leq 100 \text{ MeV} . \tag{43}$$

It is now easy to extract a bound on g using the condition (43) together with the expressions for Γ and H, Eqs. (29) and (32). For that range of energies, $Q = e$ dominates the ϕ interaction. We get

$$g < 2 \times 10^{-7} \text{ GeV}^{-1} \tag{44}$$

for $m < 2.6$ MeV. The region limited by BBN is shown in Fig. 2.

VI. EXTENSIONS

In this section we will consider first the case that ϕ is a scalar particle, and second we will generalize the couplings in such a way that they have the full SU(2)$\times$U(1) gauge invariance.

The motivation to consider a particle ϕ coupled to two photons was that such particle could be detected in some of the undergoing axion experiments. In the previous sections, we have considered a pseudoscalar particle with a coupling as in Eq. (3). We could as well consider a scalar particle ϕ_s, with mass m_s, that would couple as

$$\mathcal{L} = \frac{1}{4}\, g_s\, F^{\mu\nu} F_{\mu\nu}\, \phi_s. \tag{45}$$

One can show that all the limits we have for the pseudoscalar coupling g hold exactly for the scalar coupling g_s. Our Fig. 2 showing the regions allowed for g can be used without changes for g_s (and m_s instead of m). The considerations regarding ϕ as dark matter can also be translated into identical statements for ϕ_s, and similarly in Fig. 6 we can interchange g by g_s. At this point it is instructive to know that the optical experiments [17], mentioned in Sec. III, will be able to distinguish between the effects of a scalar ϕ_s and a pseudoscalar ϕ boson since they lead to different signatures.

Our second point is related to gauge invariance. The Lagrangian (3), containing a dimension-five operator, represents an effective interaction. New physics at an energy scale $\Lambda \sim g^{-1}$ would lead to the low energy interaction expressed in (3). This scale is much larger than the weak scale, $\Lambda \gg G_F^{-1/2}$. For energies in between Λ and $G_F^{-1/2}$, the effective interaction should have the full SU(2)$\times$U(1) gauge invariance. Our purpose is to explore the consequences of the imposition of gauge invariance. To proceed, we need to specify the behavior of ϕ under the gauge symmetry. For simplicity, we take ϕ as a SU(2)$\times$U(1) singlet.

There are two types of operators that are interesting for our purposes and preserve SU(2)$\times$U(1) gauge invariance. In the form of pieces of the effective Lagrangian they are

$$\mathcal{L}_{BB} = \frac{1}{8}\, g_{BB}\, \varepsilon_{\mu\nu\alpha\beta} B^{\mu\nu} B^{\alpha\beta}\, \phi, \tag{46}$$

$$\mathcal{L}_{WW} = \frac{1}{8}\, g_{WW}\, \varepsilon_{\mu\nu\alpha\beta} \vec{W}^{\mu\nu} \vec{W}^{\alpha\beta}\, \phi, \tag{47}$$

where

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \tag{48}$$

$$\vec{W}_{\mu\nu} = \partial_\mu \vec{W}_\nu - \partial_\nu \vec{W}_\mu + \frac{e}{\sin\theta_W} \vec{W}_\mu \times \vec{W}_\nu, \tag{49}$$

and

$$B_\mu = -\sin\theta_W Z_\mu + \cos\theta_W A_\mu, \tag{50}$$

$$W_\mu^3 = \cos\theta_W Z_\mu + \sin\theta_W A_\mu, \tag{51}$$

$$W_\mu^1 = \frac{1}{\sqrt{2}} \left(W_\mu + W_\mu^\dagger\right), \tag{52}$$

$$W_\mu^2 = \frac{i}{\sqrt{2}} \left(W_\mu - W_\mu^\dagger\right), \tag{53}$$

Z_μ, A_μ, and W_μ being the fields corresponding to the particles Z, γ, and $W^\pm$; θ_W is the weak mixing angle. Both operators (47) lead to ϕ coupling to two photons exactly as in Eq. (3). However, they lead to more than this, as we will now explore for each piece of the Lagrangian (47) in turn.

$\mathcal{L}_{BB}$ leads to the couplings $\gamma\gamma\phi$: $ZZ\phi$, and $\gamma Z\phi$:

$$\mathcal{L}_{BB} = \frac{1}{8}\, g_{BB} \cos^2\theta_W\ \varepsilon_{\mu\nu\alpha\beta} F^{\mu\nu} F^{\alpha\beta}\ \phi + \frac{1}{8}\, g_{BB} \sin^2\theta_W\ \varepsilon_{\mu\nu\alpha\beta} Z^{\mu\nu} Z^{\alpha\beta}\ \phi - \frac{1}{4}\, g_{BB} \sin\theta_W \cos\theta_W\ \varepsilon_{\mu\nu\alpha\beta} F^{\mu\nu} Z^{\alpha\beta}\ \phi, \quad (54)$$

where $Z_{\mu\nu} = \partial_\mu Z_\nu - \partial_\nu Z_\mu$.

$\mathcal{L}_{WW}$ is richer in structure:

$$\mathcal{L}_{WW} = \frac{1}{8}\, g_{WW} \sin^2\theta_W\ \varepsilon_{\mu\nu\alpha\beta} F^{\mu\nu} F^{\alpha\beta}\ \phi + \frac{1}{8}\, g_{WW} \cos^2\theta_W\ \varepsilon_{\mu\nu\alpha\beta} Z^{\mu\nu} Z^{\alpha\beta}\ \phi + \frac{1}{4}\, g_{WW} \sin\theta_W \cos\theta_W\ \varepsilon_{\mu\nu\alpha\beta} F^{\mu\nu} Z^{\alpha\beta}\ \phi + \cdots ; \quad (55)$$

namely, it generates similar couplings to those generated by $\mathcal{L}_{BB}$ and in addition it generates $WW\phi$, $WWZ\phi$, and $WW\gamma\phi$ couplings not explicitly written in the above expression.

From the phenomenological point of view the most interesting interaction, apart from $\gamma\gamma\phi$, is $\gamma Z\phi$. It would contribute to $e^+e^- \to Z \to \gamma\phi$, giving rise to single photons with unbalanced momentum at the Z peak. Such a signature has been found at the CERN e^+e^- collider LEP, consistent with the standard model expectations, $e^+e^- \to Z \to \nu\bar{\nu}\gamma$. The observed signal is consistent with the standard model prediction ($N_\nu = 3$). This implies limitations on the strength of the different and *a priori* independent pieces of the effective Lagrangian (47). We will not allow for unnatural cancellations of the effects produced by different pieces, and thus we will consider one piece at a time.

Let us consider $\mathcal{L}_{BB}$. The observations made by the OPAL Collaboration at LEP [30] place a limit on the combination:

$$g_{\gamma Z\phi} = -g_{BB}\ \sin\theta_W \cos\theta_W. \quad (56)$$

Since we identify the $\gamma\gamma\phi$ coupling in Eq. (54) as

$$g = g_{BB}\ \cos^2\theta_W, \quad (57)$$

we have that

$$g = -g_{\gamma Z\phi}\ \cot\theta_W \quad (58)$$

and thus the experimental limit on $g_{\gamma Z\phi}$ translates into a limit on g. We get

$$g \leq 1.2 \times 10^{-4}\ \text{GeV}^{-1}. \quad (59)$$

This limit holds provided ϕ does not decay inside the detector [see Eqs. (6) and (7)] and provided $m < 64$ GeV [30].

A similar reasoning leads to a limit on g when considering $\mathcal{L}_{WW}$:

$$g \leq 3.6 \times 10^{-5}\ \text{GeV}^{-1}. \quad (60)$$

We take as our result the *less* restrictive bound (59). We stress that our hypotheses are that ϕ is an SU(2)$\times$U(1) singlet and that there are no unnatural cancellations of the effects caused by $\mathcal{L}_{BB}$ and $\mathcal{L}_{WW}$. As was discussed in Sec. II, the coupling g in the Lagrangian (3) is most restricted by the ASP results: $g \leq 5.5 \times 10^{-4}$ GeV^{-1}. In this section, we have extended the Lagrangian (3) such that it possesses gauge invariance. The appearance of a coupling $\gamma Z\phi$ makes possible to have a strongest bound $g \leq 1.2 \times 10^{-4}$ GeV^{-1}, using now the OPAL results.

VII. SUMMARY AND CONCLUSIONS

Most experiments searching for axions are based on its coupling to two photons. These experiments are also sensitive to a pseudoscalar (or scalar particle) ϕ that couples only to two photons, and not to leptons, quarks, and nucleons. Motivated by this fact, we have examined the constraints on such a particle, and investigated to which extent ϕ can be the dark matter of the Universe. Some of the constraints can be deduced quite easily from studies on the axion, and other constraints have been deduced in this paper.

The laboratory, astrophysical, and cosmological limits are shown in Fig. 2. High energy searches of $e^+e^- \to \gamma +$ invisible give the best constraints of what we have classified as nondedicated experiments. Among the dedicated experiments, the solar flux detection gives strong limits once one assumes ϕ production in the solar core. Laser experiments give poorer limits, but are free of any astrophysical assumption. The telescope search gives very strong constraints, but in a very limited range of ϕ masses. Consideration of He-burning stars allows one to place very stringent limits for $m \lesssim 10$ keV. For higher ϕ masses, one has to rely on the limits from SN 1987A observations and from considerations of big bang nucleosynthesis.

We have studied the cosmological evolution of the ϕ species, and calculated the relic ϕ density. The interesting range of masses and couplings that leads to a ϕ density such that ϕ can be at least the galactic dark matter is shown in Fig. 6. Unfortunately, the mass range interesting for dark matter is much higher than the masses to which most of the existing experiments are sensitive. Only the telescope search experiment is sensitive to masses that are close to the dark matter range.

Another conclusion that we have reached has been that, if ϕ is a relic species, it must be a hot relic.

The case that ϕ is a scalar particle is very similar to the pseudoscalar case, regarding the constraints on the coupling and our conclusions on dark matter.

A final aspect we have studied is the SU(2)$\times$U(1)-gauge-invariant generalization of the ϕ interactions. In addition to the vertex $\gamma\gamma\phi$, one then has a vertex of the type $\gamma Z\phi$, as well as other exotic couplings. Experimental data from $e^+e^- \to \gamma +$ missing at the Z peak lead to limits on the coupling g that are stronger than the limits obtained from this process without the gauge-invariant generalization.

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