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Majoron as the QCD axion in a radiative seesaw model

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The smallness of neutrino mass, the strong CP problem, and the existence of dark matter are explained in an economical way. The neutrino mass is generated by the colored version of a radiative seesaw mechanism by using color adjoint mediators. The Majorana mass term of the adjoint fermion, which carries lepton number $U(1)_{\mathbb{L}}$, is induced by its spontaneous breaking, resulting in a Majoron which doubles as the quantum chromodynamics axion, thereby solving the strong CP problem. The breaking of $U(1)_{\mathbb{L}}$ sets simultaneously the seesaw scale for neutrino mass and the Peccei-Quinn breaking scale. This axion is a good candidate for dark matter as usually assumed.

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I. INTRODUCTION

The standard model of elementary particles (SM) has succeeded in describing high energy phenomena up to the TeV scale. However, there are several experimental and observational evidences of new physics beyond the SM, i.e. tiny neutrino masses, the existence of dark matter (DM) and dark energy, the density fluctuation from cosmic inflation, the baryon asymmetry of the Universe and so on. From the theoretical point of view, variations of hierarchy problems, such as the strong CP problem, the naturalness of the Higgs boson mass, the cosmological constant, and the mass hierarchy of the SM fermions, are still issues to be explored and understood.

As for the smallness of the neutrino mass, many seesaw mechanisms have been proposed. Since neutrinos may be Majorana particles [1], many people focus on this possibility to explain the big differences between neutrino masses and the ordinary SM fermion masses. The Majorana neutrino mass is allowed by the unique dimension-five operator of the SM [2], which may be implemented by a new naturally large mass scale of the operator. The simplest realization of this operator is the so-called type-I seesaw [3], where the right-handed singlet partners of the SM neutrinos are introduced as mediators. Two more ways (type II and type III) exist to fulfil the seesaw mechanism at tree level [4,5]. It was recognized many years ago [6] that these are the only three ways and the type-I, -II, and -III nomenclature was first introduced, together with the observation that there are generically also only three ways to realize this dimension-five operator in a one-particle-irreducible one-loop diagram. Recently, a review (see [7] and the references therein) of the many varieties of such radiative seesaw models appeared.

There are many observational evidences of DM. Its existence is no longer in doubt. On the other hand, no candidate particle is available within the SM. Whereas primordial black holes remain a possibility, this solution is likely to be ruled out by future observations and numerical studies [8]. A new particle, sometimes introduced for a solution to a different problem of the SM, has been known as a good candidate for DM. Especially, a weakly interacting massive particle (WIMP) is a popular candidate, where its relic abundance is naturally fixed by thermal freeze-out [9]. The strongly interacting massive particle scenario was also spotlighted recently as the new candidate for thermal DM [10]. Many alternative candidates (WIMPzilla [11], Q-ball [12], and axion [13]) are also known in the broad range of the DM mass, where the right amount of DM can be achieved by nonthermal production.

The only unobserved parameter in the SM is the QCD θ term. It is related to the chiral rotations of the quark fields through their mass terms; thus it is natural to expect a nonzero value. However, it has a tiny upper bound of 10^{-11} [14], which is indicative of a fine-tuning problem or a new mechanism to forbid it. One way to solve the problem is to consider the massless up quark [15], where the θ term is rotated away. Another solution is the Peccei-Quinn (PQ) mechanism [16], where θ is promoted to a dynamical field [17]. The vanishing θ term is realized by its dynamical relaxation in the potential containing the vacuum expectation value (VEV) of the PQ symmetry breaking.

In this Letter, we propose a new model which explains the smallness of neutrino mass, the strong CP problem and the existence of DM. The Majorana neutrino mass is generated by a one-loop radiative seesaw mechanism, where new color octet scalar and fermion fields circulate in the loop. The lepton number conservation symmetry is identified as the PQ symmetry, and its spontaneous breaking produces a Majoron [18] as an axion for the solution to the strong CP problem. This basic idea goes back many years [19,20] and has recently been applied [21] to a different axion model. In this model, it gives rise to a Majorana mass term of the octet

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fermion. Therefore, the seesaw scale for neutrino mass and the PQ symmetry breaking are related to each other. The Majoron (QCD axion) is also used for DM as usual to make this a minimal model [13], although it is possible [22] to have an additional WIMP candidate.

This paper is organized as follows. In Sec. II, the particle content and the brief sketch of our new model are given with the relevant Lagrangian terms. In Sec. III, the neutrino mass generation mechanism, the solution to the strong *CP* problem, and the axion DM scenario are shown. The possible signatures and constraints of the model are also discussed. Conclusion and discussion are given in Sec. IV.

II. MODEL

To realize the PQ mechanism, colored fermions are needed which couple anomalously to $U(1)_{\mathrm{PQ}}$, and the existence of a singlet scalar is also assumed which breaks it spontaneously. In addition to the well-known KSVZ [23] and DFSZ [24] axion models, a third option exists in supersymmetry, using the gluino, i.e. a color octet fermion, assuming that its mass is dynamically generated [25]. In this gluino axion model, the gluino plays the role of the heavy quark in the KSVZ model. The idea of our new model is to use a *gluino* for the neutrino mass generation.

The particle content of the model is given in Table I. A singlet scalar with the lepton number $\mathbb{L} = -2$ is added to the radiative seesaw model proposed by Fileviez Perez and Wise [26], which is the color octet version of the simple scotogenic model [27]. The color adjoint fermions $\Psi_R^A (A = 1, 2, \ldots, 8)$ and scalars Φ^A for the radiative seesaw mechanism are analogs to the right-handed neutrinos and the inert Higgs doublet in the scotogenic model. Whereas an *ad hoc* dark parity was imposed originally to guarantee the stability of DM, it was shown more recently [28,29] that this dark parity is in fact derivable from lepton parity, a phenomenon applicable to many simple dark matter models proposed since 30 years ago. Unlike the scalar in the scotogenic model, the new colored scalar bosons may decay into the SM quarks through the Yukawa interactions,

$$\mathcal{L}_{Q\Phi q_R} = g_u^{ij}\overline{Q_i}\,\widetilde{\Phi^A}\,T^A u_{jR} + g_d^{ij}\bar{Q}_i\Phi^A T^A d_{jR} + \mathrm{H.c.} \quad (1)$$

where $\widetilde{\Phi^A} = i\sigma_2\Phi^{A\star}$, i, $j = 1, 2, 3$ are the flavor indices, $g_q^{ij} (q = u, d)$ are the arbitrary Yukawa coupling constants, and the $SU(2)_L$ and $SU(3)_C$ indices are summed implicitly. For definiteness, we assume that Φ^A is much heavier than the weak scale, so that the flavor changing neutral current (FCNC) problem does not happen.[1] The new colored fermions also have Yukawa interactions with the SM left-handed lepton doublets and the scalar color octet,

$$\mathcal{L}_{L\Phi\Psi_R} = h_\Psi^{ij}\tilde{\Phi}^{A\dagger}\overline{\Psi_{jR}^A}L_i + \mathrm{H.c.} \quad (2)$$

where h_Ψ^{ij} are the Yukawa coupling constants whose structure is related to the observed neutrino mass and mixing parameters [31]. The lepton number of the colored fermions is determined through this interaction, i.e. $\mathbb{L}(\Psi_R) = 1$. In order to fit the observed neutrino oscillation data, at least two flavors of new Majorana fermions are required. Hereafter, we assume three generations of gluino-like particles just for simplicity. The color octet $SU(2)_L$ doublet scalar field is parametrized as

$$\Phi^A = \begin{pmatrix} H^{+A} \\ (H^A + iA^A)/\sqrt{2} \end{pmatrix}. \quad (3)$$

The Majorana mass term for the colored fermions is forbidden by the lepton number conservation, while the Yukawa interactions with the SM singlet scalar are allowed,

$$\mathcal{L}_{S\Psi_R\Psi_R} = -\frac{1}{2}y_\Psi^i S\overline{(\Psi_{iR}^A)^c}\Psi_{iR} + \mathrm{H.c.} \quad (4)$$

Without any loss of generality, the Yukawa coupling matrix y_Ψ is taken to be diagonal. Indeed, the Majorana mass for each colored fermion is obtained after developing the VEV of the singlet, i.e., $M_{\Psi i} = y_\Psi^i\langle S\rangle$. Since the global lepton number symmetry is broken spontaneously, a Nambu-Goldstone boson, so-called Majoron, appears. Thanks to the existence of the new colored fermions (gluinolike particles), the Majoron is identified as an axion. Note that the lepton number symmetry $U(1)_{\mathbb{L}}$ plays the role of the $U(1)_{\mathrm{PQ}}$ symmetry in this model.

The model is a minimal setup to solve the strong *CP* problem, the existence of DM, and the smallness of neutrino masses at the same time. In the normal approach, the strong *CP* problem and the neutrino mass generation are considered as different problems, so that the mass scales are introduced separately for each problem. In our model, the seesaw scale and the PQ symmetry breaking scale have the common origin.[2] The only energy scales introduced in

TABLE I. New fields introduced to the maxion model.

	S	Ψ_R^A	Φ^A
$SU(3)_C$	1	8	8
$SU(2)_L$	1	1	2
$U(1)_Y$	0	0	1/2
$U(1)_{\mathbb{L}}$	−2	1	0
spin	0	1/2	0

[1]If $g_{u,d}^{ij}$ are small enough while keeping the prompt decay of colored particles, Φ^A can become somewhat light. Further suppression of the FCNC is also possible by applying the minimal flavor violation hypothesis [30], i.e. the Yukawa coupling matrices $g_{u,d}^{ij}$ are proportional to the quark Yukawa matrices $Y_{u,d}^{ij} = \sqrt{2}M_{u,d}^{ij}/v$ in the SM.

[2]In Ref. [19] (see also Ref. [32]), the identification of the PQ symmetry and the lepton number symmetry is discussed in the KSVZ realization with the type-I seesaw mechanism. In their model, the Majorana mass for right-handed neutrinos and the Dirac mass for the singlet heavy quark are arranged separately, but are generated by the VEV of the same singlet.

our model are the negative mass squared of S for the PQ symmetry breaking and the dimensionful parameter for $\Phi^{A\dagger}\Phi^A$ term in addition to the one in the SM Higgs sector. From the viewpoint of the number of new fields, our model is comparable to the invisible axion models with the tree level seesaw mechanism. In addition to the common singlet field S and a new mediator for the neutrino mass generation, singlet chiral quarks with different PQ charges are introduced in the KSVZ model, while two Higgs doublets are required in the DFSZ model. In all conventional models with the seesaw extension as well as in our model, three kinds of new particles are required.

The scalar potential of this model is given by

$$\begin{aligned}\mathcal{V} = &-\mu^2 H^\dagger H - \mu_S^2 S^\star S + M_\Phi^2 \Phi^{A\dagger}\Phi^A + \lambda(H^\dagger H)^2 \\ &+ \lambda_S (S^\star S)^2 + \lambda_{SH}(S^\star S)(H^\dagger H) + \lambda_{S\Phi}(S^\star S)\Phi^{A\dagger}\Phi^A \\ &+ \lambda_3 (H^\dagger H)\Phi^{A\dagger}\Phi^A + \lambda_4 |H^\dagger \Phi|^2 \\ &+ \frac{1}{2}\{\lambda_5 (H^\dagger \Phi^A)^2 + \mathrm{H.c.}\} + \cdots \end{aligned} \tag{5}$$

where H is the Higgs doublet in the SM.[3] As long as $M_\Phi^2 \lesssim \lambda_{S\Phi}\langle S\rangle^2$, the mass of the new scalar doublet is controlled by the singlet VEV similarly to the DFSZ model. In this case, the model essentially has one new physics scale.

III. SOLUTION TO THE PROBLEMS

A. Neutrino mass

After developing the VEV of the singlet, the model arrives at the Perez-Wise model, where the neutrino mass is generated at the one-loop level with colored mediators. The Feynman diagram for the neutrino mass generation is given in Fig. 1. By calculating this diagram, we obtain

$$\begin{aligned}(\mathcal{M}_\nu)_{ij} = &-\frac{1}{4\pi^2}\sum_k h_\Psi^{ik} h_\Psi^{jk} M_{\Psi k}\left(\frac{M_H^2}{M_{\Psi k}^2 - M_H^2}\ln\frac{M_H^2}{M_{\Psi k}^2}\right. \\ &\left. -\frac{M_A^2}{M_{\Psi k}^2 - M_A^2}\ln\frac{M_A^2}{M_{\Psi k}^2}\right),\end{aligned} \tag{6}$$

where the mass eigenvalues for the neutral component of the colored scalar are $M_{H,A}^2 = M_\Phi^2 + \frac{1}{2}\lambda_{S\Phi} f_a^2 + (\lambda_3 + \lambda_4 \pm \lambda_5)v^2$. The mass matrix takes the same form as the one in the scotogenic model [27] up to the additional color factor of 8. The structure of the mass matrix is easily maintained by the Yukawa coupling structure. The smallness of the neutrino mass is naturally explained not only by heavy colored particles but also by the radiative mechanism. Note that the magnitude of the Yukawa coupling constants h_Ψ^{ij} and the mass squared difference of the colored scalars ($\propto \lambda_5$) are additional sources of the suppression factor for the tiny neutrino mass.[4] Utilizing this freedom to maintain the small neutrino mass, it is also possible to keep masses of the new colored particles in the TeV scale.

[3] The complete scalar potential of H and Φ^A can be found in Ref. [30].

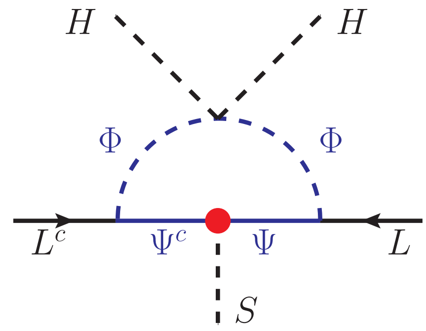

FIG. 1. One-loop diagram for neutrino mass generation.

Depending on how the neutrino mass is suppressed in the mass formula, varieties of the signature of the model are expected [33]. If the new colored particles are not superheavy, the new colored particle production can happen at the high energy frontier machine. Especially, the same-sign dilepton signature (without missing energy) will probe the lepton number violating nature of the Majorana neutrino mass. The displaced-vertex signature due to the long-lived color octet fermion is also interesting because it will probe the scale of the superheavy mediator. At the luminosity frontier, searches for the charged lepton flavor violations $\ell_i \to \ell_j \gamma$ and the electroweak precision test are also useful to explore these heavy particles. These different searches obtain information on different parameters in the neutrino mass formula, and are thus complementary.

[4] In a limit $2\lambda_5 v^2 \ll m_0^2 = (M_H^2 + M_A^2)/2$, the neutrino mass matrix is simplified as

$$(\mathcal{M}_\nu)_{ij} \simeq \frac{1}{4\pi^2}\lambda_5 v^2 \sum_k h_\Psi^{ik} h_\Psi^{jk} M_{\Psi k}\frac{M_{\Psi k}^2 \ln\frac{M_{\Psi k}^2}{m_0^2} - M_{\Psi k}^2 + m_0^2}{(M_{\Psi k}^2 - m_0^2)^2}. \tag{7}$$

For $2\lambda_5 v^2 \ll m_0^2 \ll M_{\Psi k}^2$ and $\lambda_5 \simeq 1$, $h_\Psi^{ik} \simeq 0.1$, $y_\Psi^i \simeq 1$, the axion decay constant f_a becomes $\mathcal{O}(10^{12})$ GeV. The neutrino mass vanishes if we take one of three parameters, λ_5, $M_{\Psi k}$, h_Ψ^{ik}. Each of them corresponds to the symmetry violating parameter in the lepton number broken phase depending on the choices of the global $U(1)_{\mathbb{L}}$ charges of Ψ_R^A and Φ^A,

$$\mathbb{L}(\Psi_R^A) = 1, \qquad \mathbb{L}(\Phi^A) = 0 \quad \Rightarrow \mathbb{L}(M_{\Psi k}) \neq 0, \tag{8}$$

$$\mathbb{L}'(\Psi_R^A) = 0, \qquad \mathbb{L}'(\Phi^A) = 1 \quad \Rightarrow \mathbb{L}(\lambda_5, g_{u,d}^{ij}) \neq 0, \tag{9}$$

$$\mathbb{L}''(\Psi_R^A) = 0, \qquad \mathbb{L}''(\Phi^A) = 0 \quad \Rightarrow \mathbb{L}(h_\Psi^{ik}) \neq 0. \tag{10}$$

Note that the operator $(H^\dagger \Phi^A)^2$ can be generated by the quark loop effect even if λ_5 is 0 at tree level.

B. Strong *CP* problem

The effective axion-gluon-gluon coupling is generated by the triangle anomaly diagrams via the interaction between the Majoron and the color adjoint fermions,

$$\mathcal{L}_a = -\frac{g^2}{32\pi^2}\left(\theta - \frac{3n_\Psi a(x)}{f_a}\right)\tilde{G}^{A\mu\nu}G^A_{\mu\nu}, \qquad (11)$$

where we have also included the QCD θ term in the Lagrangian, and $n_\Psi(=3)$ is the number of the color adjoint fermions. The gluon field strength tensor is $G^{A\mu\nu}$, f_a is the axion decay constant, the axion field $a(x)$ is the phase of the electroweak singlet for the PQ symmetry breaking, i.e. $S(x) = \frac{1}{\sqrt{2}}(f_a + \sigma(x))e^{ia(x)/f_a}$, and $\sigma(x)$ is a real scalar field with a mass of order f_a. A factor of 3 in front of n_Ψ is the consequence of the adjoint representation.[5] After the QCD phase transition, the axion potential becomes [34]

$$\mathcal{V}_a = \left(\frac{f_a}{3n_\Psi}\right)^2 m_a^2\left\{1 - \cos\left(\theta - \frac{3n_\Psi a(x)}{f_a}\right)\right\}, \qquad (12)$$

by the nonperturbative effect of QCD. The axion mass is related to the decay constant similarly to the standard QCD axion as [35]

$$m_a \simeq 6\ \mu\mathrm{eV} \times \left(\frac{10^{12}\ \mathrm{GeV}}{f_a/(3n_\Psi)}\right). \qquad (13)$$

By minimizing the axion potential, the *CP* invariance of the strong interaction is achieved dynamically.

C. Dark matter

The axion is known as a candidate for cold DM. In the cosmic evolution, we assume that PQ symmetry breaking occurs before or during inflation. Under this assumption the axion field becomes homogeneous, so domain walls and axion strings are absent in our Universe. Thus the only process relevant to axion DM production is coherent oscillation due to the vacuum misalignment. The current axion energy density is given by [36,37]

$$\Omega_a h^2 \approx 0.18\theta_i^2\left(\frac{f_a/(3n_\Psi)}{10^{12}\ \mathrm{GeV}}\right)^{1.19}, \qquad (14)$$

where h is the present Hubble parameter in units of $100\ \mathrm{km/s/Mpc}$, and θ_i is the initial axion misalignment angle, which takes the range $(-\pi, \pi)$. Since we assume that the PQ symmetry is broken before inflation ends, θ_i takes the same constant value in the whole Universe and is considered as a free parameter. Hence the observed value $\Omega_{\mathrm{DM}}h^2 \sim 0.12$ [38] of the energy density for DM is easily explained. A robust lower bound on the decay constant $f_a/(3n_\Psi) \gtrsim 4\times 10^8$ GeV is known from the measured duration time of the neutrinos from the supernova SN 1987A [39].

We note that the gluino axion model suffers from the cosmological domain wall problem [40], because the domain wall number is $N_{\mathrm{DW}} = 3n_\Psi$ and cannot be one, as in the KSVZ model. If the inflation finishes before the PQ symmetry breaking, the axion field does not become homogeneous. As a result, domain walls are formed by the axion potential, Eq. (12). For this reason, it is necessary to assume that the PQ symmetry is broken before or during the inflation. Conversely, the color adjoint axion model can be verified if the inflation scale is determined by future observation.

A constraint can be derived from the isocurvature fluctuation. From Planck result [38],

$$\sqrt{\mathcal{P}_S/\mathcal{P}_\zeta} \lesssim 0.18, \qquad \mathcal{P}_\zeta \simeq 2.2\times 10^{-9}, \qquad (15)$$

where $\mathcal{P}_S$ and $\mathcal{P}_\zeta$ are the dimensionless power spectrum of the DM isocurvature and curvature perturbations, respectively. In our model, scalar S has nonzero VEV during inflation, so that $\mathcal{P}_S$ becomes

$$\mathcal{P}_S \simeq \left(\frac{H_{\mathrm{inf}}}{\pi(f_a/(3n_\Psi))\theta_i}\right)^2\left(\frac{\Omega_a h^2}{\Omega_{\mathrm{CDM}}h^2}\right)^2, \qquad (16)$$

where H_{inf} is the Hubble parameter during inflation. Therefore, H_{inf} is bounded to be

$$H_{\mathrm{inf}} \lesssim 2\times 10^7\ \mathrm{GeV}\theta_i^{-1}\left(\frac{10^{12}\ \mathrm{GeV}}{f_a/(3n_\Psi)}\right)^{0.19}. \qquad (17)$$

IV. CONCLUSION AND DISCUSSION

We have constructed a model which explains the smallness of neutrino mass, the existence of cosmic DM, and the absence of strong *CP* violation at the same time. Color octet fermions (which carry lepton number) and scalars (which do not) are introduced to obtain Majorana neutrino masses by the radiative seesaw mechanism. In addition, a SM singlet scalar (which carries two units of lepton number) is chosen to break the lepton number symmetry dynamically. The color octet fermions obtain masses as a result, and the associated Goldstone boson plays the dual role of the Majoron as well as the QCD axion, because PQ symmetry is now identified with lepton number symmetry. The neutrino seesaw scale is thus also the PQ breaking scale. This axion is assumed to provide the necessary relic abundance to account for the DM of the Universe by a nonthermal production mechanism.

This model also has the potential to explain other issues beyond the SM. The real component of the singlet scalar

[5]For one flavor of the fundamental representation, the factor is 1 as in the KSVZ model.

may be identified as the inflaton, whereas the decay of color octet fermions may be used to facilitate leptogenesis [41]. These topics are beyond the scope of this Letter, and are discussed elsewhere.

As an aside, we point out a possible realization of the PQ symmetry in the radiatively induced Dirac neutrino mass model [42]. Leptoquark fields Φ_{LQ} and φ are introduced to the KSVZ model so as to close the one-loop diagram for the neutrino mass generation. To be specific, the terms $\bar{L}(\Psi_Q)_R i\sigma_2\Phi^\star_{\mathrm{LQ}}$, $\overline{(\Psi_Q)_L}N_R\varphi$, and $\Phi^\dagger_{\mathrm{LQ}}H\varphi$ are added, where Ψ_Q is a color triplet vectorlike fermion. By requiring Yukawa interactions (or the vanishing PQ charge) for $(\Psi_Q)_R$ with SM particles, the PQ charges of N_R, Ψ_L, S are determined to be the same and nonzero, which forbid the tree level neutrino mass automatically. An axion in this extension is no longer Majoron because of the lepton number conservation. Strong *CP* problem and the DM relic abundance and other topics beyond the SM can be explained in an analogous fashion.

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