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Early decay of Peccei–Quinn fermion and the IceCube neutrino events



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ABSTRACT

IceCube observed high-energy neutrino flux in the energy region from TeV to PeV. The decay of a massive long-lived particle in the early universe can be the origin of the IceCube neutrino events, which we call an “early decay scenario.” In this paper, we construct a particle physics model that contains such a massive long-lived particle based on the Peccei–Quinn model. We calculate the present neutrino flux, taking account of realistic initial energy distributions of particles produced by the decay of the massive long-lived particle. We show that the early decay scenario naturally fits into the Peccei–Quinn model, and that the neutrino flux observed by IceCube can be explained in such a framework. We also see that, based on that model, a consistent cosmological history that explains the abundance of the massive long-lived particle is realized.

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1. Introduction

After the observations of high energy cosmic ray neutrino events at the IceCube [1–5], the understanding of the origin of such high energy neutrinos has become an important task in the field of astrophysics and particle cosmology. This is because, for the energy range of $\mathcal{O}(10 \text{ TeV}) \lesssim E_\nu \lesssim \mathcal{O}(1 \text{ PeV})$, the observed flux is significantly larger than that expected from the background, *i.e.*, the atmospheric neutrinos. There have been many attempts to understand the origin from astrophysical [6–23]¹ and particle physics [26–41] points of view, although no scenario is considered to be conclusive yet.

As one of the possibilities, it was pointed out that a long-lived particle, with its mass being much higher than the PeV scale and its lifetime shorter than the present age of the universe, may be the origin of flux of high energy cosmic ray neutrinos [42].² (Hereafter, such a scenario is called as the “early decay scenario.”) In the early decay scenario, even though the neutrino produced by the decay of the long-lived particle has much larger energy than the PeV scale, which is the maximum energy scale of the deposited energy observed by IceCube, the energy is red-shifted so that the present energy can be $\mathcal{O}(10 \text{ TeV}) \lesssim E_\nu \lesssim \mathcal{O}(1 \text{ PeV})$. In Refs. [42,44], it was discussed that the decay of a long-lived particle with its mass of $\mathcal{O}(10^{10} \text{ GeV})$ may well explain the IceCube

events, without specifying the particle-physics model for the long-lived particle.

In considering the early decay scenario, one of the questions is the origin of the mass scale of the long-lived particle. In particular, the IceCube events may suggest the existence of a new physics at such a scale, into which the long-lived particle is embedded. The suggested mass scale is, as we mentioned above, $\sim 10^{10} \text{ GeV}$; one of the well-motivated new physics at such a scale is the Peccei–Quinn (PQ) symmetry [45,46] as a solution to the strong CP problem. The suggested scale of the PQ symmetry breaking is $10^9 \text{ GeV} \lesssim f_a \lesssim 10^{12} \text{ GeV}$ [see, *e.g.*, Refs. [47–49]], which well agrees with the mass scale required to realize the early decay scenario. Importantly, in hadronic axion models [50,51], the PQ mechanism requires the existence of new colored fermions. Then, if the model is embedded into grand unified theories (GUTs), there exist GUT partners of the new colored fermions. Some of these fermions (which we call PQ fermions) are stable if they do not have any mixing with the standard model (SM) fermions. With the mixing being suppressed enough, those fermions have long-enough lifetime required from the early decay scenario. As we will see below, some of the PQ fermion may play the role of the long-lived particle whose decay explains the IceCube events.

In this paper, we study the possibility of explaining the flux of cosmic ray high energy neutrinos observed by the IceCube using the decay of the PQ fermion. We assume that a color- and charge-neutral fermion (called X) embedded into the PQ fermions becomes long-lived, and that the decay of X becomes the dominant source of the cosmic ray neutrinos in the energy range of $\mathcal{O}(10 \text{ TeV}) \lesssim E_\nu \lesssim \mathcal{O}(1 \text{ PeV})$. We concentrate on the case where

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E-mail address: ema@hep-th.phys.s.u-tokyo.ac.jp (Y. Ema).¹ For more on this respect, see, *e.g.* Refs. [24,25] and references therein.² See Ref. [43] as a similar scenario.

Table 1

Particle contents and their representation under the SM gauge group as well as $U(1)_{PQ}$ and $U(1)_X$ in our model. They are embedded into $\mathbf{5}$ and $\bar{\mathbf{5}}$ representations of the $SU(5)$ GUT theory. Here, they are written as left-handed Weyl spinors.

	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{PQ}$	$U(1)_X$
X_3	$\bar{\mathbf{3}}$	$\mathbf{1}$	+1/3	+1	+1
X_2	$\mathbf{1}$	$\mathbf{2}$	-1/2	+1	+1
\bar{X}_3	$\mathbf{3}$	$\mathbf{1}$	-1/3	+1	-1
\bar{X}_2	$\mathbf{1}$	$\mathbf{2}$	+1/2	+1	-1

the lifetime of X is much shorter than the present age of the universe, and calculate the present flux of the neutrinos originating from the decay of X . For an accurate calculation of the present neutrino flux, we carefully study the spectrum of the secondary neutrinos produced from the decays of leptons and hadrons emitted from X . Then, we derive the mass, lifetime, and the relic abundance of the long-lived particle to explain the observed neutrino flux. We also discuss the production mechanism of the right amount of the long-lived particle in the early universe. We will see that the neutral PQ fermion can be a good candidate of the long-lived particle to realize the early decay scenario, if its mass is $\mathcal{O}(10^{10}$ GeV) and also if its lifetime is $\mathcal{O}(10^{11})$ sec.

The organization of this paper is as follows. In Sec. 2, we specify the particle physics model we consider in this paper, and calculate the present flux of neutrino produced by the decay of X based on that model. In Sec. 3, we discuss various phenomenological aspects of our scenario. In particular, as we will see in Sec. 2.3, the required abundance of X to explain the IceCube events is quite tiny, and hence we discuss how such a tiny amount of X is produced in the early universe. Other cosmological concerns such as the dark matter and the fate of massive particles other than X are also discussed there. The last section is devoted to the summary.

2. Present flux of neutrino

2.1. Model

Here we specify our particle physics model. We introduce a heavy vector-like fermion X_2 that has the same Standard Model (SM) gauge charges as the SM lepton doublet. We take the Peccei–Quinn (PQ) model [45,46] which solves the strong CP problem [see, e.g., Ref. [52]] as an example containing such a fermion.³ In the hadronic axion model [50,51], an additional heavy quark that carries the $U(1)_{PQ}$ charge is introduced. If the SM gauge groups are embedded in the grand unified theory (GUT) such as $SU(5)$, there must be an additional $SU(2)$ doublet fermion in the theory. After the PQ symmetry breaking, it acquires a mass of the order of the PQ scale. We summarize the particle contents of our model in Table 1.

We identify the neutral component of X_2 as a massive long-lived particle (which we call X). The long lifetime can be naturally realized as follows. We can assign a conserved $U(1)_X$ charge to the PQ fermions if there is no mixing term with the SM sector, and hence the lightest fermion in the PQ sector is stable in that case. In general, the colored partner of X_2 is heavier than X_2 because of the renormalization group (RG) running effect from the GUT scale to the PQ scale. The charged component of X_2 also acquires mass from radiative corrections after the electroweak phase transition as we will see in Sec. 3.3. Thus, X is typically the lightest fermion in the PQ sector, and is stable if there is no mixing term between the

SM and PQ sectors. If the $U(1)_X$ is (slightly) broken, the following Yukawa term is allowed as an interaction between the PQ-sector and the SM-sector⁴:

$$\mathcal{L}_{\text{mix}} = \epsilon_y \bar{X}_{2L} H e_R + \text{h.c.}, \quad (2.1)$$

where H and e_R are the SM Higgs doublet and right-handed lepton, X_{2L} is the left-handed component of X_2 , and we suppress the flavor indices for simplicity. It is technically natural in 't Hooft's sense [56] to take the coupling ϵ_y small because $U(1)_X$ symmetry is restored in the limit of $\epsilon_y \rightarrow 0$. The decay rate of X is suppressed by ϵ_y , and X can be a massive long-lived particle.

Once the Yukawa term is fixed, X decays as $X \rightarrow l^\pm W^\mp$ after the electroweak symmetry breaking, where l^\pm is the SM charged lepton and W^\pm is the W gauge boson. The decay rate Γ_X is given by

$$\tau_X^{-1} \equiv \Gamma_X \sim \alpha \epsilon_y^2 m_X, \quad (2.2)$$

where τ_X is the lifetime of X and α is a numerical constant of order $\mathcal{O}(10^{-2})$. For simplicity, we assume that X equally couples to electron, muon and tau lepton. Also we assume that the CP symmetry is conserved in the X decay so that particles and anti-particles are equally produced. As we will see in Sec. 2.3, this decay mode is enough to explain the high-energy neutrino events observed by IceCube.

2.2. Boltzmann equation

In this subsection, we explain the time evolution of the high-energy neutrinos produced by the decay of the massive long-lived particle X in the early universe. We follow the calculation procedure given in Ref. [44].

Neutrinos propagate after produced by the decay of X , being sometimes scattered by the background particles (especially neutrinos). In order to obtain the present flux of neutrino, we numerically solve the following Boltzmann equation:

$$\begin{aligned} & \left(\frac{\partial}{\partial t} + 2H - HE \frac{\partial}{\partial E} \right) \Phi_{\nu,l}(t, E) \\ &= -\gamma_{\nu,l}(t; E) \Phi_{\nu,l}(t, E) \\ &+ \int dE' \Phi_{\nu,n}(t, E') \frac{d\gamma_{\nu,nm}(t; E', E)}{dE} P_{ml}(t, E) \\ &+ S_{\nu,m}(t, E) P_{ml}(t, E), \end{aligned} \quad (2.3)$$

where t is the cosmic time, H is the Hubble parameter, $\Phi_{\nu,l}(t, E)$ is the neutrino flux in l -th flavor, $S_{\nu,l}(t, E)$ is the source term, $\gamma_{\nu,l}(t; E)$ is the scattering rate and $d\gamma_{\nu,nm}(t; E', E)/dE$ is the differential neutrino production rate with E' (E) being the energy of the initial-state (final-state) neutrino. The Roman indices denote the flavors of the neutrinos. The flux is related to the number density of neutrino as $n_{\nu,l}(t) = 4\pi \int dE \Phi_{\nu,l}(t, E)$ where $n_{\nu,l}(t)$ is the number density of neutrino in l -th flavor. Moreover, we have introduced the ‘‘transition probability’’ $P_{ml}(t, E)$ to take into account the neutrino oscillation effects. For more details, see Ref. [44].

From now, we explain more details about the treatment of the source term. Since we assume that the source of the high-energy neutrinos is the decay of X , the source term is given by

$$\begin{aligned} S_{\nu,l}(t, E) &= \frac{1}{4\pi} \frac{n_X(t)}{\tau_X} \frac{dN_{\nu,l}^{(X)}}{dE} \\ &= \frac{1}{4\pi} \frac{Y_X}{\tau_X} s(t) e^{-t/\tau_X} \frac{dN_{\nu,l}^{(X)}}{dE}, \end{aligned} \quad (2.4)$$

³ The gauge-mediated supersymmetry breaking model [53–55] is another candidate which also includes such a vector-like fermion, but we do not discuss it in this paper.

⁴ A mass mixing term is equivalent to the Yukawa term after diagonalizing the mass matrix.

where $dN_{\nu,l}^{(X)}/dE$ is the energy distribution of neutrinos produced by the decay of one X , τ_X is the lifetime of X , n_X is the number density of X and Y_X is the yield variable of X which is defined as

$$Y_X \equiv \left[\frac{n_X(t)}{s(t)} \right]_{t \ll \tau_X}, \quad (2.5)$$

with s being the entropy density. The energy distribution has two contributions:

$$\frac{dN_{\nu,l}^{(X)}}{dE} = \frac{dN_{\nu,l}^{(X,\text{dir})}}{dE} + \frac{dN_{\nu,l}^{(X,\gamma\gamma)}}{dE}. \quad (2.6)$$

Neutrinos are directly produced from the decay of X , and the contribution from such neutrinos is expressed as $dN_{\nu,l}^{(X,\text{dir})}/dE$. On the other hand, the decay of X also produces electromagnetic particles, especially photons, electrons and positrons. These electromagnetic particles produce high-energy neutrinos by the scattering with the cosmic microwave background (CMB) photons, and we denote such a contribution to the flux as $dN_{\nu,l}^{(X,\gamma\gamma)}/dE$. We estimate

$$\frac{dN_{\nu,l}^{(X,\gamma\gamma)}}{dE} = \int d\epsilon \left(\frac{dN_{\gamma}^{(X)}}{d\epsilon} + \frac{dN_{e^\pm}^{(X)}}{d\epsilon} \right) \frac{dN_{\nu,l}^{(\gamma\gamma)}(\epsilon, E)}{dE}, \quad (2.7)$$

where $dN_{\gamma}^{(X)}/d\epsilon$ and $dN_{e^\pm}^{(X)}/d\epsilon$ are energy distributions of photon and e^\pm produced by the decay of one X , and $dN_{\nu,l}^{(\gamma\gamma)}/dE$ is the spectrum of neutrinos produced by the double-photon pair creation processes with CMB photons. Here e^\pm is converted to photons with (almost) identical energy via the inverse Compton scattering process. We have used PYTHIA package [57,58] to calculate $dN_{\nu,l}^{(X)}/dE$, with the decay mode determined by Eq. (2.1). For more details on the treatment of high-energy electromagnetic particles, see Ref. [44].

Once the X 's mass m_X is fixed, we can calculate the energy distribution $dN_{\nu,l}^{(X)}/dE$. Thus, the source term, and also the present flux of neutrino are determined by the following three parameters:

$$m_X, z_* \equiv z(\tau_X), Y_X. \quad (2.8)$$

We use the redshift z_* , at which the cosmic time is equal to τ_X , to parametrize τ_X . It depends on the Yukawa coupling ϵ_Y .

2.3. Present flux and IceCube results

From now, we discuss implications of the early-decay scenario for the high-energy neutrino events observed by IceCube.

The IceCube collaboration observed high-energy neutrinos [1–5]. In the three-year observation, they detected 37 events in the energy region of $30 \text{ TeV} \lesssim E \lesssim 2 \text{ PeV}$, three of which are with the deposited energy of $1 \text{ PeV} \lesssim E \lesssim 2 \text{ PeV}$. Given that the expected backgrounds are 8.4 ± 4.2 events for the atmospheric muon and $6.6_{-1.6}^{+5.9}$ events for the atmospheric neutrinos, it indicates a new source of the high-energy neutrinos. In particular, in Ref. [59], they performed a detailed analysis of the high-energy neutrino flux.⁵ Assuming the power-law, the best fit all-flavor flux⁶ for the energy region between 25 TeV and 2.8 PeV is given by⁷

$$E^2 \sum_{l=e,\mu,\tau} \Phi_{\nu,l}(E) = \left(6.7_{-1.2}^{+1.1} \right) \times 10^{-8} \left(\frac{E}{100 \text{ TeV}} \right)^{-0.50 \pm 0.09} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \quad (2.9)$$

In addition, in Ref. [4], the IceCube collaboration shows their preliminary results of four-year observation of high-energy neutrinos. They detected additional 17 events in the energy region of $30 \text{ TeV} \lesssim E \lesssim 400 \text{ PeV}$ in the fourth year. The best fit value of the spectral index for all four-year events with $60 \text{ TeV} < E < 3 \text{ PeV}$ is given by $\gamma = -2.58 \pm 0.25$, and hence the expected flux might be even softer after taking into account the fourth-year results. We also note that, recently, the highest energy charged current muon neutrino event with the reconstructed muon energy of $4.5 \pm 1.2 \text{ PeV}$ was reported [5].

In order to show that the early-decay scenario can explain the IceCube neutrino events, we take the following two sample points in our analysis:

- point 1 [left panel]: $(m_X, 1 + z_*, Y_X) = (4 \times 10^{10} \text{ GeV}, 8 \times 10^3, 8 \times 10^{-27})$.
- point 2 [right panel]: $(m_X, 1 + z_*, Y_X) = (6 \times 10^{10} \text{ GeV}, 1 \times 10^4, 1 \times 10^{-26})$.

We show the present neutrino per-flavor flux for these two sample points in Fig. 1. The left panel corresponds to the point 1, and the right panel does to the point 2. The flux is close to the E^{-2} power-law with an exponential cut-off in the case of the point 1, while the spectral index of the flux is rather close to $\gamma \simeq -2.5$ in the case of the point 2. Both cases give the spectral index currently consistent with the IceCube results. We emphasize here that the spectral index is quite sensitive to m_X and z_* if $z_* \sim \mathcal{O}(10^3)$ and $m_X \sim \mathcal{O}(10^{10} \text{ GeV})$. This is because the neutrino scattering effects begin to be effective at around such an epoch. An $\mathcal{O}(10\%)$ tuning of the parameters is necessary to obtain the neutrino flux with the spectral index in the range, e.g., between -2 and -2.5 . In addition, the highest energy muon neutrino event may be also due to the neutrinos produced in the early decay scenario if the neutrino spectrum extends up to 3–4 PeV (which is the case with the above two sample points), assuming that a significant fraction of the energy of the primary muon neutrino is carried away by the muon; using the total exposure of the IceCube detector given in Ref. [5], the expected number of events for $2.5 < E < 4 \text{ PeV}$ is 0.5 for both the points 1 and 2.

We also note here that, for the energy of $E \simeq 6.3 \text{ PeV}$, the event rate of the anti-electron neutrino is enhanced because of the on-pole exchange of the W -boson (so-called the ‘‘Glashow resonance’’ [65]), which gives a stringent bound on the neutrino flux. With four years of observation [4], no event was reported with the energy deposit of $E \simeq 6.3 \text{ PeV}$. Using the effective area provided by the IceCube collaboration,⁸ the expected number of events at the Glashow resonance is calculated to be 0.05 and 0.1 for the points 1 and 2, respectively, for 1347 days of lifetime (which is the lifetime of the four years of data). Thus, in the early decay scenario, the neutrino flux at the Glashow resonance can be suppressed enough due to the exponential suppression of the neutrino flux above a few PeV. Future study of neutrino events at the Glashow resonance will be helpful to further constrain the spectral shape of the neutrino flux.

The neutrino scattering with the background neutrinos produces electromagnetic particles, and such electromagnetic particles may slightly modify the thermal history of the universe, leaving

⁵ See also Refs. [60–64].

⁶ Assuming the flavor universality, we obtain the per-flavor flux by dividing it by three.

⁷ A power-law flux with a spectral index $\gamma = 2$ and an exponential cut-off might be disfavored compared to a free spectral index one.

⁸ For the data, see <https://icecube.wisc.edu/science/data/HE-nu-2010-2012>.

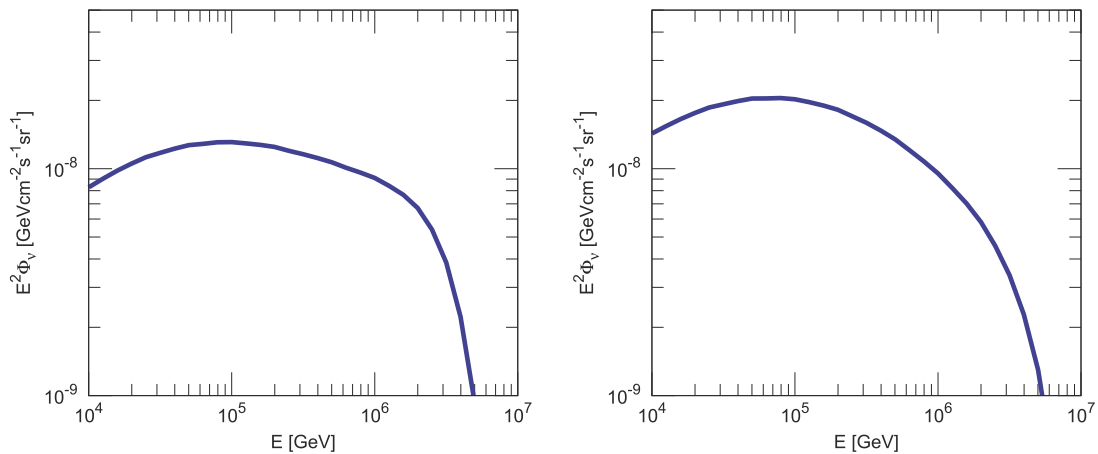


Fig. 1. The present per-flavor neutrino flux for two sample points. Left panel: $(m_X, 1 + z_*, Y_X) = (4 \times 10^{10} \text{ GeV}, 8 \times 10^3, 8 \times 10^{-27})$. Right panel: $(m_X, 1 + z_*, Y_X) = (6 \times 10^{10} \text{ GeV}, 1 \times 10^4, 1 \times 10^{-26})$. The branching ratio is $\text{Br}(X \rightarrow l^\pm W^\mp) = 1$. We take the produced leptons as flavor-universal, and also assume that particles and anti-particles are equally produced. We plot the flavor-averaged value of the neutrino flux. The present neutrino flux is almost flavor-universal after taking account of the neutrino scattering and the neutrino oscillation.

the trail in the present universe. In the case of our interest, constraints come from the CMB y - and μ -type distortions [66–72] and the big bang nucleosynthesis (BBN) [73,74]. We estimate the y and μ parameters as $(y, \mu) = (4 \times 10^{-9}, 2 \times 10^{-10})$ for the point 1 and $(y, \mu) = (6 \times 10^{-9}, 5 \times 10^{-10})$ for the point 2. Although they are consistent with the current the CMB observations [75,76], future experiments such as PIXIE [77] and PRISM [78] might have a chance to detect the CMB distortions caused by the decay of X . We have also checked that the points 1 and 2 are not excluded by the BBN constraints [73,74].

There are three comments. First, we comment on the angular distribution of the neutrino flux. In the early-decay scenario, the neutrino flux is isotropic, and hence it is currently consistent with the IceCube results. It is in contrast to the case of the dark matter (DM) decay scenario [26–37]. In the DM decay scenario, the IceCube events are assumed to be explained by the decay of DM, and hence the galactic contribution dominates over the extragalactic one. Thus, a large fraction of neutrino comes from the galactic center region, resulting in a sizable dipole asymmetry [42]. Further study on the angular distribution may be helpful to distinguish these two scenarios.

Second, we comment on the flavor composition of the neutrino flux. So far the flavor ratio of the neutrino flux observed by IceCube is consistent with the universal one $\Phi_{\nu,e} : \Phi_{\nu,\mu} : \Phi_{\nu,\tau} = 1 : 1 : 1$ [59,79]. In the early-decay scenario, the present neutrino flux is almost flavor-universal after taking into account the neutrino oscillation as well as the scattering with background neutrinos, even if it is not at the onset of the production. Thus, it is currently consistent with the IceCube results. Some deviation from the flavor-universality may be helpful to exclude the early-decay scenario in the future.

Finally, we comment on the multi-messenger approach. It is a powerful way to explore the property of the source to combine the IceCube data with the gamma-ray observation such as the Fermi collaboration [80]. In fact, regardless of production mechanism pp or $p\gamma$, the authors in Refs. [10,23] argued that the astrophysical source should be opaque to 1–100 GeV gamma-rays if it is the origin of the IceCube neutrino events. The same method is also useful in the case of the DM decay scenario. In Ref. [34], it is discussed that although the DM decay scenario is currently consistent with the gamma-ray observation, it will be tested by near-future gamma-ray observations. In contrast to these scenarios, in the early-decay scenario, photons produced by X are absorbed into the CMB since they are produced in the early epoch. Thus, the present

gamma-ray observations are not affected by the early-decay scenario. Instead, as we discussed before, these photons might cause the CMB spectral distortions.

3. Phenomenology

3.1. Production mechanism

As we saw in Sec. 2.3, the desired yield value of X is

$$Y_X \sim 10^{-26}, \quad (3.1)$$

to explain the neutrino flux observed by the IceCube experiment. From now we discuss how such a tiny amount of X is produced in the early universe. Our purpose here is to show one specific example, and hence the discussion is not generic.

We first give an overview of our production scenario. In this paper, we concentrate on the case where the inflaton mass m_ϕ is larger than m_X . To be more specific, we take the inflaton mass to be $m_\phi \sim \mathcal{O}(10^{11} \text{ GeV})$, which is the case for, e.g., some class of new inflation models [81,82].⁹ Then, there are four ways to produce X after inflation:

- (1) Inflaton decay
The inflaton must decay to transfer its energy density to ordinary matters. If the branching ratio of the inflaton ϕ decaying into X is nonzero, X can be produced from the inflaton decay.
- (2) Thermal production by dilute plasma
There exists radiation called dilute plasma even before the reheating completes. The maximal temperature of the dilute plasma T_{max} is typically much higher than the reheating temperature T_R , and hence it may thermally produce the massive particle X .
- (3) Non-thermal production through inelastic scattering
Decay products of the inflaton with its energy $\sim m_\phi$ may produce X through inelastic scattering with the dilute plasma if $T_{\text{max}} \gg m_X^2/m_\phi$ [83].
- (4) Scattering among inflaton decay products before thermalization
High-energy particles with its energy $\sim m_\phi$ dominates the energy density of the radiation sector before the thermalization of the dilute plasma completes [84,85]. Scatterings of such high-energy particles among themselves can also produce X .

⁹ As we will see later, this type of models is also compatible with the axion dark matter scenario.

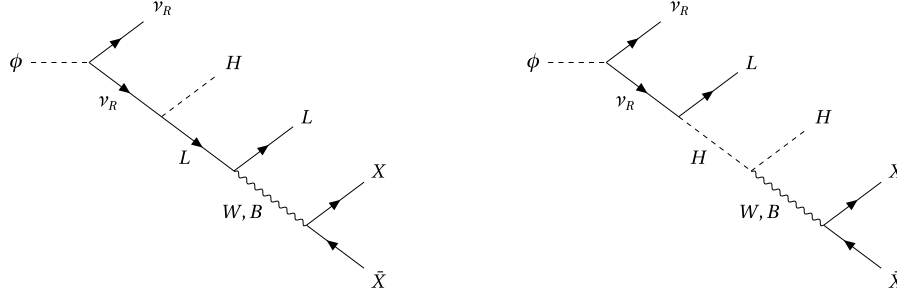


Fig. 2. The Feynman diagrams for the five-body decay of ϕ to X . The $SU(2)_L$ and $U(1)_Y$ gauge bosons are denoted as W and B , respectively. Note that the electroweak symmetry is restored at the reheating epoch, and hence we do not distinguish each component of the $SU(2)_L$ doublets. They are generated using `TikZ-Feynman` [91].

Since the desired yield value (3.1) is tiny, X is typically overproduced if it is once in thermal equilibrium in the early universe. Even the non-thermal production by the radiation tends to overproduce X . Thus, we take the reheating temperature T_R to be low, of order $\mathcal{O}(10^3 \text{ GeV})$, so that the production mechanisms (2) and (3) can be neglected. As we will see below, the contribution from the production mechanism (4) is also negligible in such a case. Instead, we assume that X is mainly produced via a many-body decay process of ϕ . In such a case, the branching ratio of ϕ decaying to X is well-suppressed by couplings and the phase space factor, and hence we can naturally control the yield of X produced by the mechanism (1).

Now we explain details of the production mechanism. We consider the inflaton ϕ which couples only with the right-handed neutrino via the following interaction:

$$\mathcal{L}_{\phi\nu\nu} = -\frac{1}{2}\lambda\phi\overline{\nu}_R^c\nu_R + \text{h.c.}, \quad (3.2)$$

where ν_R is the right-handed neutrino and $\nu_R^c = C\overline{\nu}_R^T$ where the charge conjugation matrix C is given as $C = i\gamma_0\gamma_2$ in the Weyl representation. We denote the masses of ϕ and ν_R as m_ϕ and m_N , respectively. The right-handed neutrino couples to the SM Higgs and leptons as usual:

$$\mathcal{L}_Y = -y\overline{L}\tilde{H}\nu_R + \text{h.c.}, \quad (3.3)$$

where L is the SM left-handed lepton doublet, $\tilde{H} = (\epsilon H)^\dagger$ with ϵ being the antisymmetric tensor of $SU(2)$, and y is the matrix of the Yukawa coupling. Here we have suppressed the isospin and the generation indices for simplicity. Note that some components of y can be arbitrarily small even assuming the see-saw mechanism [86–90]. This is because only the mass squared differences of the left-handed neutrinos are currently observed.

The inflaton starts to oscillate around the bottom of its potential after inflation, and the reheating proceeds through the interaction (3.2). The decay rate of ϕ is estimated as

$$\Gamma_\phi \simeq \mathcal{O}(0.01)\lambda^2 m_\phi, \quad (3.4)$$

and the reheating temperature is given by

$$\begin{aligned} T_R &\sim \sqrt{M_P\Gamma_\phi} \\ &\sim 10^3 \text{ GeV} \left(\frac{\lambda}{10^{-11}} \right) \left(\frac{m_\phi}{10^{11} \text{ GeV}} \right)^{1/2}. \end{aligned} \quad (3.5)$$

We assume that the inequalities

$$m_\phi > m_X \sim \mathcal{O}(10^{10} \text{ GeV}) > m_N, \quad (3.6)$$

are satisfied.¹⁰ Then, ϕ directly decays into X through five-body decay processes $\phi \rightarrow X\overline{X}\nu_R LH$ [see Fig. 2] since X couples to the

Standard Model particles via the gauge interaction.¹¹ The five-body decay rate is estimated as

$$\Gamma_{\text{five}} \simeq \frac{c\lambda^2 y^2 g^4}{(4\pi)^7} m_\phi, \quad (3.7)$$

where c is a numerical factor and g is the $SU(2)$ or $U(1)$ gauge coupling constant. The factor $(4\pi)^{-7}$ comes from the phase-space suppression of the five-body decay. We have verified $c \lesssim \mathcal{O}(0.1)$ in the case of our interest by using `FeynRules` [92] and `MadGraph5` [93].¹² Note that only the off-shell processes contribute due to the inequality $m_N < m_X$. The yield of X produced via this process is estimated as

$$\begin{aligned} Y_X^{(1)} &= \left[\frac{n_X}{s} \right]_{T=T_R} \sim \frac{T_R}{m_\phi} \frac{\Gamma_{\text{five}}}{\Gamma_\phi} \\ &\simeq 10^{-26} \left(\frac{c}{0.1} \right) \left(\frac{y^2}{10^{-12}} \right) \left(\frac{m_\phi}{10^{11} \text{ GeV}} \right)^{-1} \left(\frac{T_R}{10^3 \text{ GeV}} \right). \end{aligned} \quad (3.8)$$

Thus, the tiny amount of X can be explained via the mechanism (1) for reasonable model parameters.

Next, we consider the production mechanism (2) and (3). It is well-known that there exists a dilute plasma even before the reheating completes. If $m_\phi \gg T_R$, the particles produced by ϕ are initially under-occupied compared with the thermal distribution. In such a case, these particles do not instantaneously thermalize. Instead, the so-called “bottom-up thermalization” process occurs [84, 85]. The maximal temperature of such a dilute plasma is estimated as

$$\begin{aligned} T_{\text{max}} &\sim \left(\frac{\alpha_s^4 M_P^2}{T_R m_\phi} \right)^{1/5} T_R \\ &\sim 10^7 \text{ GeV} \left(\frac{\alpha_s}{0.1} \right)^{4/5} \left(\frac{T_R}{10^3 \text{ GeV}} \right)^{4/5} \left(\frac{10^{11} \text{ GeV}}{m_\phi} \right)^{1/5}, \end{aligned} \quad (3.9)$$

where α_s is the fine structure constant of the $SU(3)$ gauge group. If T_{max} approaches to m_X , the production mechanisms (2) and (3) typically overproduce X . Thus, we limit ourselves to the case with low T_R , e.g. of order $\mathcal{O}(10^3 \text{ GeV})$. Then, the thermal production (2) is well suppressed by the Boltzmann factor. The non-thermal production (3) is also negligible since $T_{\text{max}} < m_X^2/2m_\phi$ in this case.

Now let us move on to the production mechanism (4). Since we take the reheating temperature to be low, it takes a long time for the dilute plasma to achieve the thermalization. Before the thermalization is achieved, the energy density of the decay products is dominated by high-energy particles whose typical energy is of order $\sim m_\phi$ [85]. These particles can produce X by scattering with

¹⁰ Since λ is tiny [see Eq. (3.5)], the right-handed neutrino mass induced by the interaction (3.2) is small compared to m_X even if $\phi \sim M_P$. Note that $\phi \ll M_P$ typically holds for new inflation models.

¹¹ The $SU(2)$ gauge symmetry is usually restored at the time of the inflaton decay, and hence we do not distinguish each component of the $SU(2)$ doublets here.

¹² If m_ϕ is close to m_X , the numerical factor c can be much smaller than unity.

each other. The number density of high-energy particles is estimated as

$$n_{\text{high}} \sim \frac{\rho_\phi \Gamma_\phi}{m_\phi H} \sim \frac{M_P H T_R^2}{m_\phi}, \quad (3.10)$$

and hence the number density of X produced during one Hubble time is given by

$$n_X \sim \frac{\langle \sigma v \rangle n_{\text{high}}^2}{H} \sim \frac{\alpha_s^2 M_P^2 H T_R^4}{m_\phi^4}, \quad (3.11)$$

where we have estimated the averaged cross section as $\langle \sigma v \rangle \sim \alpha_s^2 / m_\phi^2$. The number density of X produced at $H = H_p$ contributes to the yield as

$$Y_X^{(4)} = \left[\frac{n_X}{s} \right]_{T_R} \sim \frac{T_R n_X}{m_\phi n_\phi} \Big|_{T_R} \sim \frac{T_R n_X}{m_\phi n_\phi} \Big|_{H=H_p} \sim \alpha_s^2 \left(\frac{T_R}{m_\phi} \right)^4 \left(\frac{T_R}{H_p} \right). \quad (3.12)$$

Thus, the late time production gives the dominant contribution. Just before the thermalization, the Hubble parameter is estimated as

$$H(T = T_{\text{max}}) \sim \frac{T_{\text{max}}^4}{M_P T_R^2} \sim \left(\frac{\alpha_s^4 M_P^2}{T_R m_\phi} \right)^{4/5} \frac{T_R^2}{M_P}, \quad (3.13)$$

and hence the yield of X produced by the mechanism (4) is given by

$$Y_X^{(4)} \sim \alpha_s^{-6/5} \left(\frac{T_R}{M_P} \right)^{3/5} \left(\frac{T_R}{m_\phi} \right)^{16/5} \sim 10^{-34} \left(\frac{0.1}{\alpha_s} \right)^{6/5} \left(\frac{T_R}{10^3 \text{ GeV}} \right)^{19/5} \left(\frac{10^{11} \text{ GeV}}{m_\phi} \right)^{16/5}. \quad (3.14)$$

Therefore, it is negligible for $T_R \sim \mathcal{O}(10^3 \text{ GeV})$ and $m_\phi \sim \mathcal{O}(10^{11} \text{ GeV})$.

Finally, we comment on the baryogenesis. In our scenario, the reheating temperature should be low enough to avoid the overproduction of X . In such a case, the thermal [94] and non-thermal [81, 95, 96] leptogenesis scenarios do not work in its original form. Still, if the heavy neutrino masses are sufficiently degenerate, the leptogenesis may be possible [97]. The Affleck–Dine [98] or electroweak baryogenesis [99–101] may be other possibilities to explain the baryon to photon ratio of the universe.

3.2. Dark matter

In this subsection, we show that our scenario naturally accommodates the axion dark matter (DM). In order to see this point, let us consider the dynamics of the PQ symmetry breaking field in the early universe.

We assume that the inflation scale is low, such as those of the models in Refs. [81, 82]. Note that the mass scale of the inflaton is $m_\phi \sim \mathcal{O}(10^{11} \text{ GeV})$ in these models, and hence it is consistent with our production scenario. Then, the PQ symmetry is already broken during inflation. In such a case, the coherent oscillation of the axion provides a natural candidate of DM.¹³ The present energy density of the axion is given by [102]

$$\Omega_a h^2 = 0.18 \theta_a^2 \left(\frac{f_a}{10^{12} \text{ GeV}} \right)^{1.19} \left(\frac{\Lambda_{\text{QCD}}}{400 \text{ MeV}} \right), \quad (3.15)$$

where θ_a is the misalignment angle at the onset of the coherent oscillation and Λ_{QCD} is the energy scale of the QCD phase transition. For $f_a \sim \mathcal{O}(10^{11} - 10^{12} \text{ GeV})$, the axion coherent oscillation can provide the present DM abundance if $\theta_a \sim \mathcal{O}(1)$. Thus, our scenario is consistent with the axion DM scenario.

Here we comment on the isocurvature perturbation. If the PQ symmetry is broken during inflation, the axion already exists and acquires fluctuations during that epoch. They contribute to the DM isocurvature perturbation, which is constrained by the observation [103]. In our case, however, the inflation scale is taken to be low, and thus the isocurvature perturbation is suppressed by the factor H_{inf}/f_a where H_{inf} is the Hubble parameter during inflation. It is consistent with the observation for $H_{\text{inf}}/f_a \lesssim 10^{-6}$.

3.3. Lifetimes of other particles

We consider the neutral component of a heavy SU(2) doublet as X , which is a mother particle of the high energy neutrinos observed by IceCube. There are also the charged component of the SU(2) doublet and a heavy SU(3) triplet in the theory. In this subsection, we discuss fates of these particles in the early universe.

First, let us consider the heavy SU(3) triplet (denoted as X_3). We assume that they are more massive than the SU(2) doublet that includes X . Note that it is usually the case due to the RG effects. Since X_3 and X_2 are embedded into the same GUT multiplet, each component of X_3 can decay into X_2 by exchanging massive gauge bosons. The decay rate is roughly estimated as

$$\Gamma_{X_3} \sim \frac{g^4 m_{X_3}^5}{(4\pi)^3 m_V^4}, \quad (3.16)$$

where g is the gauge coupling of the GUT, m_{X_3} is the mass of X_3 and m_V is the mass of the massive gauge bosons which is of the order of the GUT scale. If we take $m_V \sim \mathcal{O}(10^{16} \text{ GeV})$ and $m_{X_3} \sim \mathcal{O}(10^{11} \text{ GeV})$, for example, the decay rate is estimated as $\Gamma_{X_3} \sim 10^{-7} \text{ GeV}$. Thus, X_3 does not affect any observables because it decays well before the BBN. Moreover, the yield of X_3 produced in the early universe is at most comparable to that of X_2 , and hence the estimation of Y_X in the previous subsection still holds.

Next, let us consider the charged component of X_2 . We denote it as X^\pm here. Before the electroweak symmetry breaking, it is stable because the mass of X^\pm is exactly the same as that of X . However, a mass difference is generated radiatively after the electroweak symmetry breaking. In the limit $m_X \gg m_Z$ where m_Z is the Z -boson mass, the mass difference is given by [104, 105]

$$\delta m \equiv m_{X^\pm} - m_X \simeq \frac{1}{2} \alpha_2 \sin^2 \theta_W m_Z \simeq 350 \text{ MeV}, \quad (3.17)$$

where α_2 is the fine structure constant of the SU(2) gauge group, θ_W is the weak mixing angle and m_{X^\pm} is the mass of X^\pm . Due to this mass difference, X^\pm decays into X by mainly emitting a soft charged pion [104, 105]. The lifetime is of order $\mathcal{O}(10^{-10} \text{ sec})$, and hence it starts to decay just after the electroweak phase transition, well before the BBN. Thus, it has no effect on any observables. Again, the estimation of Y_X is not affected by this process since the yield of X^\pm is at most comparable to Y_X .

4. Summary

IceCube detected high-energy neutrino events whose origin is still unknown. There have been a lot of discussion on the origin from both astrophysical and particle physics point of view. As one

¹³ For a review on the axion cosmology, see, e.g. Ref. [49].

possibility, in Refs. [42,44], it was proposed that decay of a massive long-lived particle in the early universe can be the origin of the IceCube neutrino events. We have called such a scenario as “early decay” scenario, and the massive long-lived particle as X in this paper. Interestingly, it was shown that the mass scale of X can be as large as of order $\mathcal{O}(10^{10} \text{ GeV})$ if it decays well before the present universe, although the IceCube neutrino events are in the energy range of $\mathcal{O}(10 \text{ TeV}) \lesssim E_\nu \lesssim \mathcal{O}(1 \text{ PeV})$.

In this paper, we have extended the discussion in Refs. [42, 44]. We have constructed a specific particle physics model that includes X based on the Peccei–Quinn model. The hadronic axion models include an additional heavy quark, and there is also an SU(2) counterpart once we embed the model into the grand unified theory. We have identified the charge neutral component of the SU(2) doublet as X . We have calculated the energy spectrum of neutrinos produced from the decay of X , and followed the time evolution of the neutrino flux including the scattering processes with the background neutrinos. The decay of X also produces electromagnetic particles such as photons, electrons and positrons. They induce electromagnetic cascades with the background photons, and secondary neutrinos produced in such cascade processes are also included in our calculation. We have shown that our model can explain the IceCube neutrino events in the energy region of $\mathcal{O}(10 \text{ TeV}) \lesssim E_\nu \lesssim \mathcal{O}(1 \text{ PeV})$. The favored mass of X is of order $m_X \sim \mathcal{O}(10^{10} \text{ GeV})$, and the favored lifetime is of order $\tau_X \sim \mathcal{O}(10^{11} \text{ sec})$. The decay of X can also slightly modify the thermal history of the universe. We have checked that our scenario is currently consistent with the CMB and BBN constraints. Notice that our scenario predicts the y and μ parameters to be $y \sim \mathcal{O}(10^{-9})$ and $\mu \sim \mathcal{O}(10^{10})$, respectively. Thus, future CMB observations such as PIXIE [77] and PRISM [78] have a possibility to detect the CMB distortions caused by the present scenario.

We have also discussed the cosmological history based on our model. In particular, the yield of X should be of order $\mathcal{O}(10^{-26})$ to reproduce the normalization of the IceCube neutrino flux. This is a rather small value, and hence we have discussed in detail how such a tiny amount of X is produced in the early universe. The thermal as well as non-thermal production by background radiation typically overproduce X , and hence we limit ourselves to the case of the low reheating temperature. Instead, we have used a many-body decay of the inflaton as a production mechanism. In such a case, the phase-space factor as well as couplings suppress the abundance of X . The dark matter abundance and fate of additional particles other than X are also briefly discussed.

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