\( \nu \text{AMDM: A Model for Sterile Neutrino and Dark Matter Reconciles Cosmological and Neutrino Oscillation Data after BICEP2} \)

P. Ko and Yong Tang

School of Physics, Korea Institute for Advanced Study, Seoul 130-722, Korea

We propose a ultraviolet complete theory for cold dark matter(CDM) and sterile neutrinos that can accommodate both cosmological data and neutrino oscillation experiments within 1\( \sigma \) level. We assume a new \( U(1)_X \) dark gauge symmetry which is broken at \( \sim O(\text{MeV}) \) scale resulting light dark photon. Such a light mediator for DM’s self-scattering and scattering-off sterile neutrinos can resolve three controversies for cold DM on small cosmological scales: cusp vs. core, too-big-to-fail and missing satellites. We can also accommodate \( \sim O(1) \) eV scale sterile neutrinos as the hot dark matter(HDM) and can fit some neutrino anomalies from neutrino oscillation experiments within 1\( \sigma \). Finally the right amount of HDM can make a sizable contribution to dark radiation, and also helps to reconcile the tension between the data on the tensor-to-scalar ratio reported by Planck and BICEP2 Collaborations.

**INTRODUCTION**

The standard model of cosmology, the so-called \( \Lambda \)CDM with the minimal six parameters, can explain well a wide range of cosmological observations, such as primordial abundance of light elements, cosmic microwave background(CMB) and large scale structures(LSS). Meanwhile, there are still some hints that indicate that new physics beyond the minimal \( \Lambda \)CDM model maybe is needed in order to explain CDM sector better.

There are three controversies for CDM paradigm on small cosmological scales, cusp vs. core, too-big-to-fail and missing satellites (see Ref. [1] for a review), which have triggered both astrophysical [2–10] and DM-related investigations [11–27]. A solution that resolves simultaneously these controversies has been proposed in Ref. [28], where both DM and active neutrino interact with a new gauge boson with mass round \( O(\text{MeV}) \). Then the DM’s velocity-dependent self-scattering cross section can be around \( 1 \text{cm}^2/\text{g} \) at the Dwarf satellites scale, and evades the constraints from Milky Way galaxy and galaxy cluster. Thus one can resolve the first two controversies. Meanwhile, the DM’s scattering off cosmic neutrino background leads to its late kinetic decoupling at temperature \( T_{\text{kin}} < O(\text{keV}) \), which is translated into a cut-off of the smallest protohalo mass \( M_{\text{cut}} \sim O(10^9)M_\odot \), resolving the 3rd puzzle, namely missing satellites problem. However, since active neutrino couples to a MeV particle, such scenario is restrictively constrained [29–31].

The CMB data indicates that a small amount of relativistic species or hot dark matter(HDM) could exist at CMB time [32–35], in addition to the standard three generations of active neutrinos. This is often parametrized as the effective number of additional neutrino species \( \Delta N_{\text{eff}}^{\text{cmb}} \). It has been shown in Ref. [34] that the best fit to all available data is given by

\[
\Delta N_{\text{eff}}^{\text{cmb}} = 0.61 \pm 0.30, \quad m_{\text{hdm}}^{\text{eff}} = (0.47 \pm 0.13) \text{ eV}, \quad (1)
\]

where \( m_{\text{hdm}}^{\text{eff}} \) is the effective HDM mass. Also, it was shown very recently that a similar amount of HDM can help to relieve the tension of tensor-to-scalar ratio \( (\equiv r) \) between Planck data [37] and the recently announced measurement of B-mode polarization by BICEP2 [38], without a running spectral index [39–41].

It is well known that sterile neutrino can serve as a HDM component of the universe. Sterile neutrino is also well motivated in order to solve accelerator [42–44], reactor [44] and gallium anomalies [45–46] in neutrino oscillation experiments. Both reactor and gallium anomalies prefer a new mass-squared difference, \( \Delta m^2 \gtrsim 1 \text{ eV}^2 \) [47], while accelerator experiments [48–50] prefer \( \Delta m^2 \sim 0.5 \text{ eV}^2 \). In all three cases the favored mixing angles are around \( \sin^2 2\theta \sim 0.1 \). Such a large mixing angle would in general lead to fully thermalized sterile neutrinos by oscillation and thus an increase of \( \Delta N_{\text{eff}}^{\text{cmb}} = 1 \) for each sterile neutrino. This is in some tension with the above cosmological data Eq. (1), as shown in global fit [51] including BBN, CMB and LSS data.

The above tension can be relieved by introducing new interaction for sterile neutrino. The new interaction can generate a matter potential \( V_{\text{eff}} \) that results in a tiny effective mixing angle \( \theta_m \) in matter [52] for \( V_{\text{eff}} \gg \Delta m^2/2\pi^2 \),

\[
\sin^2 \theta_m = \frac{\sin^2 2\theta_0}{(\cos 2\theta_0 + \frac{2\pi}{\Delta m} V_{\text{eff}})^2 + \sin^2 2\theta_0},
\]

where \( \theta_0 \) is the mixing angle in vacuum. As a result, the thermalization of sterile neutrino by oscillation can be suppressed and \( \Delta N_{\text{eff}} < 1 \) is easily obtained [53–54]. Recently it has been shown that the tension in the data can be reconciled at 2\( \sigma \) level within an effective theory [55] where a dim-5 operator is responsible for the active-sterile neutrino mixing.

In this paper, we propose a ultraviolet complete theory for DM and sterile neutrino that can accommodate the aforementioned cosmological data and neutrino oscillation experiments within 1\( \sigma \) level. The model includes both CDM and HDM, and we call it the \( \nu \text{AMDM}(\text{the first M stands for mixed}) \).
MODEL FOR CDM AND STERILE NEUTRINO

We consider the standard seesaw model with two right-handed (RH) neutrinos (gauge singlet) $N_i (i = 1, 2)$ and add a dark sector with $U(1)_X$ gauge symmetry and coupling $g_X$, and dark photon field $\hat{X}_\mu$, and dark Higgs field $\phi_X$ and two different Dirac fermion $\psi$ and $\chi$ in the dark sector. All the new fields are SM gauge singlets. We assign equal $U(1)_X$ charges to $\phi_X$ and $\psi$, which is normalized to 1. Then the most general gauge invariant renormalizable Lagrangian is given by

$$
\mathcal{L} = \mathcal{L}_{SM} + \hat{N}_i i \overline{\psi} N_i - \left( \frac{1}{2} m_{\tilde{H}}^2 \tilde{H} \tilde{H}^\dagger + y_{\alpha i} \tilde{H}_\alpha H N_i + h.c. \right) - \frac{1}{4} \hat{X}_{\mu \nu} \hat{X}^{\mu \nu} - \frac{1}{2} \sin \epsilon \hat{X}_{\mu \nu} \hat{B}^{\mu \nu} + \bar{\chi} (i \not{\partial} - m_\chi) \chi + \bar{\psi} (i \not{\partial} - m_\psi) \psi + D_\mu \phi^i \not{D}^\mu \phi_X - \left( f_i \phi^i \tilde{N}_i \psi + g_i \phi_X \bar{\psi} N_i + h.c. \right)
$$

$$
- \lambda_\phi \left[ \phi_X^2 - \frac{v_\phi^2}{2} \right] - \lambda_H \left[ \phi_X^2 - \frac{v_\phi^2}{2} \right] \left[ H^\dagger H - \frac{v_h^2}{2} \right],
$$

(2)

where $L_\alpha$ are the SM left-handed lepton doublets, $H$ is the SM Higgs doublet, and $\hat{B}$ is the field strength for SM $U(1)_Y$. The covariant derivative on a field $K$ is defined as

$$
D_\mu K = (\partial_\mu - i Q_K g_X \hat{X}_\mu) K \quad \text{(with } K = \chi, \psi, \phi_X).$$

We have chosen the $U(1)_X$ charge for $\chi$ in such a way that the $\phi_X \chi N_i$ term is forbidden by $U(1)_X$ gauge symmetry (otherwise $\chi$ may decay if kinematically allowed). Thus $\chi$ would be stable and the DM candidate.

The local gauge symmetry is broken by the following vacuum configurations:

$$
\langle H \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_h \end{pmatrix}, \langle \phi_X \rangle = \frac{v_\phi}{\sqrt{2}},
$$

(3)

where $v_h \simeq 246 \text{ GeV}$ and $v_\phi \sim O(\text{MeV})$ for our interest. There will be mixings among various fields after the spontaneous gauge symmetry breaking. The gauge kinetic mixing term results in tiny mixings among the physical gauge fields, $A_\mu, Z_\mu$ and $X_\mu$. Also there is a mixing between Higgs fields $h$ and $\phi$ with

$$
H \rightarrow \frac{v_h + \phi}{\sqrt{2}} \quad \text{and} \quad \phi_X \rightarrow \frac{v_\phi + \phi}{\sqrt{2}}.
$$

Two scalar excitations $h$ and $\phi$ can be expressed in terms of mass eigenstates, $H_1$ and $H_2$, as

$$
h = H_1 \cos \alpha - H_2 \sin \alpha, \quad \phi = H_1 \sin \alpha + H_2 \cos \alpha,
$$

(4), (5)

with a mixing angle $\alpha$. Because of the Higgs portal interaction ($\lambda_{\phi H}$ term) and the additional scalar $\phi$, the electroweak vacuum could be stable up to Planck scale without additional new physics beyond the particle contents presented in Eq. (2) (see Refs. [57] for example).

A novel feature of this model is that there can be mixing among three active neutrinos $\nu_\alpha$, sterile neutrinos $N_i$ and dark fermion $\psi$ due to $y_{\alpha i} \bar{L}_\alpha H N_i, f_i \phi^i \tilde{N}_i \psi$ and $g_i \phi_X \bar{\psi} N_i$ after spontaneous gauge symmetry breaking. In order to correctly explain the active neutrino oscillation data, at least two $N$’s are needed, in which case two of $\nu_\alpha$ are massive and the other one is massless. Then after diagonalization of $7 \times 7$ mass matrix for $\nu_\alpha, N_i$ and $\psi, m$ eigenstates are composed of 7 Majorana neutrinos, $\nu_\alpha (a = 1, 2, 3)$ and $\nu_{\alpha i} (i = 4, ..., 7)$, or collectively $\nu_i = \nu_{iL} + \nu_{iR}$:

$$
\begin{pmatrix}
\nu_\alpha \\
N_i^c \\
\psi_L \\
\psi_{\ell L}
\end{pmatrix} = U
\begin{pmatrix}
\nu_\alpha \\
\nu_{4L} \\
\vdots \\
\nu_{7L}
\end{pmatrix},
$$

$$
U^T M U =
\begin{pmatrix}
\begin{smallmatrix}
m_1 \\
\vdots \\
m_7
\end{smallmatrix}
\end{pmatrix},
$$

where $U$ is the unitary mixing matrix that diagonalizes the mass matrix $M$,

$$
M =
\begin{pmatrix}
\begin{smallmatrix}
v_{03 \times 3} & \frac{v}{\sqrt{2}} [y_{\alpha i}]_{3 \times 2} & 0_{3 \times 2} \\
\frac{v}{\sqrt{2}} [y_{\alpha i}]^T_{3 \times 3} & \frac{m_{\nu \ell}}{\sqrt{2}} & \frac{v_\phi}{\sqrt{2}} (f_i g_i)_{2 \times 2} \\
0_{2 \times 3} & \frac{v_\phi}{\sqrt{2}} (f_i g_i)_{2 \times 2} & \frac{v}{\sqrt{2}} [y_{\alpha i}]_{1 \times 2}
\end{smallmatrix}
\end{pmatrix}.
$$

In the following discussion, if not specified, we shall use $\nu_\alpha$ and $\nu_{\alpha i}$ to collectively denote three active neutrinos and four sterile neutrinos, respectively.

The mixing also distributes the new $U(1)_X$ gauge/Yukawa interaction to all neutrinos with actual couplings depending on the exact mixing angles. We assume that the mixing angles between $\nu_\alpha$ and $\psi$ are negligible, compared to the mixing between $N_i$ and

---

1 We could add more heavy $N$ in the Lagrangian for leptogenesis [56], which will not affect our discussions in the following.
ψ. This can be easily achieved by adjusting $y_{\nu i}$, $f_i$, and $g_i$. A more straightforward way is to work in the flavor basis, in which only $N_i$ and $\phi$ have dark Yukawa and gauge interactions, respectively. Because of the new dark interactions for $\nu$, all sterile neutrinos $\nu_s$ are not thermalized by oscillation from active neutrinos and thus can contribute to the number of effective neutrino by a proper amount, $\Delta N_{\text{eff}} < 1$ after BBN [53, 54].

The exact mass spectrum and mixing angles for $\nu_s$ are free, subject to conditions for fitting the data. We shall take at least one $\nu_s$ is around 1 eV and others as free, lighter or heavier, and the mixing angles among $\nu_s$ are large enough for suppressing their production by oscillation from active neutrino.

Based on a different setup, our model improves a similar attempt presented in a recent paper [55] in two aspects. First, our model is renormalizable and thus ultraviolet complete, while the model in Ref. [55] assumed a dim-5 operator for generating the active-sterile neutrino mixing and therefore depends on the UV completion. Second, we shall show below that the model presented in the present paper can reconcile the current cosmological data with neutrino oscillation experiments within 1σ rather than only within 2σ as discussed in [55].

**Thermal History and CDM Controversies**

Communication between dark sector and SM particles or thermal history before BBN time is determined mostly by two mixing parameters, $\sin \epsilon$ and $\lambda_{dH}$. $\sin \epsilon$ is constrained by DM direct searches around $\sin \epsilon < 10^{-9}$ for $\mathcal{O}(\text{TeV}) \chi$ and $\mathcal{O}(\text{MeV}) X_{\mu}$ [53]. And $\lambda_{dH}$ as small as $10^{-8}$ would be enough to thermalize the dark sector at $T \sim \text{TeV}$ [59]. After the cross sections of dark particles' scattering off SM particle drop below the expansion rate of the Universe, the dark sector decouples from the thermal bath of the visible sector and entropy density would be conserved separately in each sector. The decoupling temperature of the dark sector, $T_{\chi}^{\text{dec}}$, would determine how much $\Delta N_{\text{eff}}$ is left at a later time. The exact value for $\Delta N_{\text{eff}}$ will be given in the following.

Chemical decoupling of DM from the heat bath sets its relic density today. After the temperature drops below $m_{\chi}$, $\chi$ starts to leave the chemical equilibrium and would finally freeze out at $T \sim m_{\chi}/25$. To account for the correct thermal relic density, the thermal cross section for $\chi \chi$ annihilation ($\sigma v$) should be around $3 \times 10^{-26}\text{cm}^3/\text{s}$. The dominant annihilation channel in this model is $\chi \chi \rightarrow X_{\mu} X_{\mu}$, and the relic density requires the gauge coupling $g_X$ to be [60]

$$g_X \sim \frac{0.50}{Q_X} \times \left( \frac{0.114}{T_{\text{cdm}}} \right)^{\frac{1}{2}} \left( \frac{m_{\chi}}{\text{TeV}} \right)^{\frac{3}{2}},$$

where $Q_X$ is the $U(1)_X$ charge of $\chi$ and shall be taken $\sim O(1)$ for definiteness in later discussion. We shall focus on the CDM $\chi$ with mass $\sim \text{TeV}$, which is preferred region as shown in Ref. [25].

Kinetic decoupling of $\chi$ from $\nu$, happens at much later time when the elastic scattering rate for $\chi \nu_s \leftrightarrow \nu_s$ drops below some value determined by Hubble parameter $H$. The Feynman diagram is shown in Fig. [1]. For a thermal distribution of sterile neutrino, the decoupling temperature is given by

$$T_X^{\text{kd}} \sim 1\text{keV} \left( \frac{0.1}{g_X} \right) \left( \frac{T_\gamma}{T_{\nu_s}} \right)^{\frac{1}{2}} \left( \frac{m_{\chi}}{\text{TeV}} \right)^{\frac{3}{2}} \left( \frac{m_{\nu_s}}{\text{MeV}} \right),$$

where $T_\gamma$ and $T_{\nu_s}$ are the temperatures of CMB and sterile neutrinos, respectively. Except that DM is dominantly scattering off sterile neutrinos in our model rather than active ones, the above formula is similar to the one in Ref. [25] and gives the approximate order-of-magnitude estimation, although the precise formula may depend on the neutrino mixing angles from the couplings $\bar{\nu}_i \gamma^\mu \nu_j X_{\mu}$.

The kinetic decoupling of DM from the relativistic particles imprints on the matter power spectrum, for which there are two relevant scales [61, 62]: the co-moving horizon $\tau_{kd} \propto 1/T_{\chi}^{\text{kd}}$ and free-streaming length $(T_{\chi}^{\text{kd}}/m_{\chi})^{1/2} \tau_{kd}$. For our interested regime, $\tau_{kd}$ is much larger and relevant. Thus $T_{\chi}^{\text{kd}}$ can be translated into a cutoff in the power spectrum of matter density perturbation with

$$M_{\text{cut}} = \frac{4\pi}{3} \rho_M (c\tau_{kd})^3 \sim 2 \times 10^8 \left( \frac{T_{\chi}^{\text{kd}}}{\text{keV}} \right)^{-3} M_\odot,$$

where $\rho_M$ is the sum of matter densities today, $\rho_{\text{CDM}} + \rho_{\text{baryon}}$. Then $M_{\text{cut}} \sim O(10^9) M_\odot$ can be easily obtained for explanation of missing satellites problem for $\mathcal{O}(\text{TeV}) \chi$ and $\mathcal{O}(\text{MeV}) X_{\mu}$.

Because of the light mediator $X_{\mu}$, the DM self-scattering $\chi \chi \rightarrow \chi \chi$ can have a large cross section, $\sigma \sim 1\text{cm}^2$ at small scales, while relative small values at Milky Way and larger scales. This can flatten the dark halo, decrease the total mass of halo centre and resolve both cusp vs. core and too-big-to-fail controversies. The
quantity that is usually used to describe the efficiency for the DM-DM self-scattering is the transfer cross section

$$\sigma_T \equiv \int d\Omega (1 - \cos \theta) \frac{d\sigma}{d\Omega}.$$ 

$\sigma_T$ can be easily calculated from Fig. 1(a) in the perturbative regime $\alpha_X m_X < m_\chi$ as,

$$\sigma_T = \frac{8\pi}{m_\chi^2} \left( \ln \left( 1 + R^2 \right) - \frac{R^2}{1 + R^2} \right),$$

$$\alpha_X = \frac{g_\chi^2}{4\pi}, \beta = \frac{2\alpha_X m_\chi}{m_\chi v_{rel}^2}, R = \frac{m_\chi v_{rel}}{m_\chi},$$

where $v_{rel}$ is the relative velocity of $\chi$ and $\bar{\chi}$. $v_{rel}$ is around 20, 200, 1000 km/s for Dwarf galaxies, Milky Way and the galaxy clusters, respectively. In the non-perturbative regime $\alpha_X m_X > m_\chi$, we have [22]

$$\sigma_T = \begin{cases} \frac{4\pi}{m_\chi^2} \beta^2 \ln \left( 1 + \beta^2 \right), & \beta \lesssim 0.2 \\ \frac{8\pi}{m_\chi^2} \beta^2 / \left( 1 + 1.5\beta^2 \right), & 0.2 \lesssim \beta \lesssim 1300 \\ \frac{8\pi}{m_\chi^2} \left( \ln \beta + 1 + \frac{1}{2} \ln \beta^2 \right), & \beta > 1300 \end{cases}$$

As an illustration, in Fig. 2 we show the case with $m_\chi = 1\,\text{TeV}, m_X = 4\,\text{MeV}$ and $g_\chi = 0.5$, in which $\sigma_T/m_\chi$ can be achieved properly for Dwarf galaxies with $v_{vel} \simeq 20\,\text{km/s}$.

**EFFECTIVE NUMBER OF EXTRA NEUTRINOS**

After the decoupling of dark sector from the visible thermal bath, relativistic particles can still contribute to the radiation density. For 4 light sterile neutrinos, their contributions to $\Delta N_{\text{eff}}$ can be parametrized as

$$\Delta N_{\text{eff}} (T) = 4 \times \frac{T^4_{\nu_s}}{T^4_{\nu_a}} = 4 \times \frac{g_{s*} (T)}{g_{a*} (T)} \times \frac{T^3_{\nu_s}}{T^3_{\nu_a}}$$

$$= 4 \times \left[ \frac{g_{s*} (T)}{g_{a*} (T)} \times \frac{T^3_{\nu_s (T_{\text{dec}})}}{T^3_{\nu_a (T_{\text{dec}})}} \right]^{\frac{4}{3}},$$

(8)

where $T$ is the photon temperature, and $g_{s*}$ counts the total number of relativistic degrees of freedom for entropy ($g_{s*}$ for dark sector). Conservation of entropy density has been used in the last step of the above equations.

When only sterile neutrinos are relativistic at the time just before BBN epoch, we have

$$g_{s*} (T_{\text{dec}}) = 3 + 1 + \frac{7}{8} \times (4 \times 2) = 11,$$

$$g_{s*} (T_{\text{dec}}) = \frac{7}{8} \times (4 \times 2) = 7.$$

The parameter $g_{s*} (T_{\text{dec}})$ is well-known in SM [63] and depends on the decoupling temperature. For example, $g_{s*} (T_{\text{dec}}) \simeq 72$ for $m_c < T_{\text{dec}} < m_\tau$. Together with

$$g_{s*} (T_{\text{dec}}) = 2 + \frac{7}{8} \times (3 \times 2 + 2 \times 2) = \frac{43}{4},$$

we can get

$$\Delta N_{\text{eff}} = 4 \times \left[ \frac{43}{4} \times \frac{11}{7} \times \frac{2}{72} \right]^{\frac{3}{4}} \simeq 0.579.$$
Increasing (decreasing) $T_x^{\text{dec}}$ gives smaller (larger) $\Delta N_{\text{eff}}$ due to the changes in $g_{ss}(T_x^{\text{dec}})$. For instance, if $T_x^{\text{dec}} > m_{\nu_s}$, we would have $g_{ss}(T_x^{\text{dec}}) \simeq 107$ and $\Delta N_{\text{eff}} = 0.341$. If $T_x < T_x^{\text{dec}} < m_{\nu_s}$, we would have $g_{ss}(T_x^{\text{dec}}) \simeq 41$ and $\Delta N_{\text{eff}} = 1.23$. Here $T_x$ is the temperature for confinement-deconfinement transition between quarks and hadrons in QCD.

Decoupling temperature lower than $T_x$ would give too large $\Delta N_{\text{eff}} \geq 3.96$ and therefore is excluded at high confidence level. The available range for $\Delta N_{\text{eff}}$ is the region between two red vertical lines in Fig. 3.

If $X_\mu$ and $H_2$ are also relativistic around BBN time, we have $g^x_{ss}(T) = g^x_{ss}(T_x^{\text{dec}})$ in Eq. (8) and additional contributions from the bosonic part

$$\Delta N_{\text{eff}}^b = 2 \times \frac{8}{7} \times \frac{T_x^{\nu_s}}{T_x^{\nu_s}},$$

where the factor 2 accounts for bosonic degrees of freedom normalized to fermionic one, $g_b/g_\nu$. The ratio of $\Delta N_{\text{eff}}$ for two cases is about

$$\text{ratio} = \frac{4 \times (\frac{11}{7})^4}{4 + 2 \times \frac{8}{7}} \simeq 1.16. \quad (10)$$

So the difference is small and we shall not distinguish two cases in the later discussion.

These extra sterile neutrinos can also be relativistic even at CMB time with $T_\gamma \simeq O(1)$ eV and play the role of HDM. Their effects on cosmology can be parametrized by the effective mass defined as

$$m_{\text{hdm}}^{\text{eff}} = \left( \frac{T_{\nu_1}}{T_{\nu_s}} \right)^3 \sum_{\nu_x} m_{\nu_x} = \left( \frac{\Delta N_{\text{eff}}}{4} \right)^2 \sum_{\nu_x} m_{\nu_x}, \quad (11)$$

where only relativistic sterile neutrinos are summed over.

Sterile neutrino masses can be chosen to fit the neutrino oscillation data. We take the face values from the global fit [64–66]: for instance, with $3 + 2$ scenario [66] gives $\Delta m^2_{12} = 0.46$ eV$^2$ and $\Delta m^2_{31} = 0.87$ eV$^2$. Since $\nu_1$ is massless in our model, we have $m_{\nu_4} \simeq 0.68$ eV and $m_{\nu_5} \simeq 0.93$ eV. Then using Eq. (11), we depict the central value of $\sum m_{\nu_x}$ as the horizontal dotted line in Fig. 3. We can see that cosmological data can be reconciled with neutrino oscillation experiments within 1$\sigma$ level in our model, which is quite remarkable.

The crucial difference between our model and Ref. [55] is due to Eqs. (8) and (11), because of “4” sterile neutrinos in our model. Usually, only one sterile neutrino is responsible for $\Delta N_{\text{eff}}$ and the relation among $m_{\nu_4}$, for which one would have

$$m_{\text{hdm}}^{\text{eff}} = (\Delta N_{\text{eff}})^{3/4} m_{\nu_4}.$$ 

Then this is consistent with neutrino oscillation data only at 2$\sigma$ level as shown in Ref. [55].

In the above discussion we have assumed that $\Delta N_{\text{eff}} (\text{BBN}) = \Delta N_{\text{eff}} (\text{CMB})$ for illustration. This assumption may not be necessarily true when either oscillation brings all neutrinos into equilibrium or some sterile neutrinos are heavy enough such that they become non-relativistic at the time before CMB and heat other neutrinos. In both cases we have $\Delta N_{\text{eff}} (\text{CMB}) < \Delta N_{\text{eff}} (\text{BBN})$, and our model predictions are still consistent with neutrino oscillation data within 1$\sigma$ level.

**FURTHER TESTS OF THE MODEL**

There are a few different ways to test our model. Direct detection of CDM $\chi$ will be possible for no vanishing sin $\epsilon$. Also, $\chi \bar{\chi}$ will annihilate to two $X_\mu$s, which in turn decay into sterile neutrinos immediately. These high energy sterile neutrinos can oscillate to active neutrinos which can be detected by neutrino telescopes, such as IceCube, whose current limit on $\langle \nu v \rangle$ is around $10^{-22}\text{cm}^3/\text{s}$ [67]. Taking into account boost factors due to light mediators in our model, future detection of these neutrino flux will be possible. Since we have more sterile neutrino species than other models, oscillation experiments could also be used to test the model even though this depends on the exact mixing angles and mass spectrum.

**SUMMARY**

In this paper, we have proposed a *ultraviolet* complete renormalizable model for self-interacting CDM and sterile neutrinos that can accommodate the cosmological data and neutrino oscillation experiments simultaneously within 1$\sigma$ level. The model is based on a dark sector with local $U(1)_X$ dark gauge symmetry that is spontaneously broken at $\mathcal{O}(\text{MeV})$ scale. The resulting $\mathcal{O}(\text{MeV})$ gauge boson (dark photon) can mediate a DM self-scattering cross section around $\sigma \sim 1\text{cm}^2/\text{g}$ which is of right order to resolve two issues for CDM at small cosmological scales, *cusp vs. core* and *too-big-to-fail*.

In our model, two light RH gauge singlet neutrinos ($N_{i=1,2}$) can mix with a dark fermion $\psi$ and therefore can interact with DM through the new dark gauge boson. The relics of these sterile neutrinos serve as the hot dark matter with a right amount of $\Delta N_{\text{eff}}$ (see Fig. 1), which relieves the tension between Planck and BICEP2. The masses of these sterile neutrinos are consistent with neutrino oscillation experiments within 1$\sigma$. Meanwhile, the interaction between DM and sterile neutrino delays the DM’s kinetic decouple to sub-keV temperature and induces a lower cut-off in the primordial matter power spectrum, resolving the *missing satellites problem*. The model could be tested further by (in)direct detection of CDM $\chi$, and also through neutrino oscillation experiments if favorable parameters are realized in Nature.
We are grateful to Wan-Il Park for useful discussions. This work is supported in part by National Research Foundation of Korea (NRF) Research Grant 2012R1A2A1A01006053, and by the NRF grant funded by the Korea government (MSIP) (No. 2009-0083526) through Korea Neutrino Research Center at Seoul National University (PK).

REFERENCE

[1] D. H. Weinberg, J. S. Bullock, F. Governato, R. K. de Naray and A. H. G. Peter, arXiv:1306.0913 [astro-ph.CO].
[2] G. Efstathiou, J. Bond, and S. D. White, Mon.Not.Roy.Astron.Soc. 258 (1992) 1–6.
[3] C. Pfisterer, P. Chang, and A. E. Broderick, Astrophys.J. 752 (2012) 24, arXiv:1106.5505 [astro-ph.CO].
[4] J. Silk and A. Nusser, Astrophys.J. 725 (2010) 556–560, arXiv:1004.0857 [astro-ph.CO].
[5] P. R. Shapiro, I. T. Illiev, and A. C. Raga, Mon. Not. Roy. Astron. Soc. 348 (2004) 753, arXiv:astro-ph/0307266.
[6] M.-M. Mac Low and A. Ferrara, Astrophys. J. 613 (1999) 142, arXiv:astro-ph/9804297.
[7] P. F. Hopkins, E. Quataert, and N. Murray, Mon.Not.Roy.Astron. Soc. 421 (2012) 3522–3537, arXiv:1110.4638 [astro-ph.CO].
[8] A. Pontzen and F. Governato, Mon. Not. Roy. Astron. Soc. 421 (2012) 3464, arXiv:1106.0499 [astro-ph.CO].
[9] C. A. Vera-Ciro, A. Helmi, E. Starkenburg, and M. A. Breddels, arXiv:1202.6061 [astro-ph.CO].
[10] J. Wang, C. S. Frenk, J. F. Navarro, and L. Gao, arXiv:1203.4097 [astro-ph.GA].
[11] D. N. Spergel and P. J. Steinhardt, Phys. Rev. Lett. 84 (2000) 3760–3763, arXiv:astro-ph/9909386.
[12] P. Bode, J. P. Ostriker, and N. Turok, Astrophys. J. 556 (2001) 93–107, arXiv:astro-ph/0010389.
[13] J. J. Dalcanton and C. J. Hogan, Astrophys. J. 561 (2001) 35–45, arXiv:astro-ph/0004381.
[14] A. R. Zentner and J. S. Bullock, Astrophys.J. 598 (2003) 49, arXiv:astro-ph/0304292.
[15] K. Sigurdson and M. Kamionkowski, Phys. Rev. Lett. 92 (2004) 171302, arXiv:astro-ph/0311486.
[16] M. Kaplinghat, Phys. Rev. D72 (2005) 063510, arXiv:astro-ph/0507300.
[17] F. Borzumati, T. Bringmann, and P. Ullio, Phys. Rev. D77 (2008) 063514, arXiv:hep-ph/0701007.
[18] C. Boehm, P. Fayet, and R. Schaeffer, Phys. Lett. B518 (2001) 8–14, arXiv:astro-ph/0102504.
[19] M. Kaplinghat, L. Knox, and M. S. Turner, Phys. Rev. Lett. 85 (2000) 3335, arXiv:astro-ph/0005210.
[20] W. Hu, R. Barkana, and A. Gruzinov, Phys. Rev. Lett. 85 (2000) 1158–1161, arXiv:astro-ph/0003365.
[21] M. R. Lovell, V. Eke, C. S. Frenk, L. Gao, A. Jenkins, et al., Mon.Not.Roy.Astron.Soc. 420 (2012) 2318–2324, arXiv:1104.2929 [astro-ph.CO].
[22] J. L. Feng, M. Kaplinghat, and H.-B. Yu, Phys. Rev. Lett. 104 (2010) 151301, arXiv:0911.0422 [hep-ph].
[23] M. R. Buckley and P. J. Fox, Phys.Rev. D81 (2010) 083522, arXiv:0911.3898 [hep-ph].
[24] A. Loeb and N. Weiner, Phys.Rev.Lett. 106 (2011) 171302, arXiv:1011.6374 [astro-ph.CO].
[25] M. Vogelsberger, J. Zavala, and A. Loeb, Mon.Not.Roy.Astron.Soc. 423 (2012) 3740, arXiv:1201.5892 [astro-ph.CO].
[26] S. Tulin, H.-B. Yu, and K. M. Zurek, arXiv:1302.3898 [hep-ph].
[27] P. Ko and Y. Tang, arXiv:1402.6494 [hep-ph].
[28] L. G. van den Aarssen, T. Bringmann and C. Pfisterer, Phys. Rev. Lett. 109, 231301 (2012) [arXiv:1205.5809 [astro-ph.CO]].
[29] B. r. Ahlgren, T. Ohlsson and S. Zhou, Phys. Rev. Lett. 111, 199001 (2013) [arXiv:1309.0991 [hep-ph]].
[30] I. M. Shoemaker, Phys. Dark Univ. 2, no. 3, 157 (2013) [arXiv:1305.1936 [hep-ph]].
[31] R. J. Wilkinson, C. Boehm and J. Lesgourgues, JCAP 1405, 011 (2014) [arXiv:1401.7597 [astro-ph.CO]].
[32] R. A. Burenin, Astron. Lett. 39, 357 (2013) [arXiv:1301.4791 [astro-ph.CO]].
[33] M. Wyman, D. H. Rudd, R. A. Vanderveld and W. Hu, Phys. Rev. Lett. 112, 051302 (2014) [arXiv:1307.7715 [astro-ph.CO]].
[34] J. Hamann and J. Hasenkamp, JCAP 1310, 044 (2013) [arXiv:1308.3255 [astro-ph.CO]].
[35] R. A. Battye and A. Moss, Phys. Rev. Lett. 112, 051303 (2014) [arXiv:1308.5870 [astro-ph.CO]].
[36] S. Garazza, C. Giunti and M. Laveder, JHEP 1311, 211 (2013) [arXiv:1309.3192 [hep-ph]].
[37] P. A. R. Ade et al. [Planck Collaboration], arXiv:1303.5076 [astro-ph.CO].
[38] P. A. R. Ade et al. [BICEP2 Collaboration], arXiv:1403.3985 [astro-ph.CO].
[39] E. Giussman, E. Di Valentino, M. Lattanzi, A. Melchiorri and O. Mena, arXiv:1403.4852 [astro-ph.CO].
[40] J. -F. Zhang, Y. -H. Li and X. Zhang, arXiv:1403.7028 [astro-ph.CO].
[41] Cora Dvorkin, Mark Wyman, Douglas H. Rudd, Wayne Hu, arXiv:1403.8049 [astro-ph.CO].
[42] C. Athanassopoulos et al. [LSND Collaboration], Phys. Rev. Lett. 77, 3082 (1996) [nucl-ex/9605003].
[43] A. Aguilar-Arevalo et al. [LSND Collaboration], Phys. Rev. D 84, 112007 (2011) [arXiv:1101.2755 [hep-ex]].
[44] M. Vogelsberger, J. Zavala, and A. Loeb, Phys. Rev. D 86, 045032 (2012) [arXiv:1204.5379 [hep-ph]].
[45] J. N. Abdurashitov, V. N. Gavrin, V. V. Gorbashchev, P. P. Gurkina, T. V. Ibragimova, A. V. Kalikhov and N. G. Khairnasov et al., Phys. Rev. C 73, 048505 (2006) [nucl-ex/0512014].
[46] C. Giunti, M. Laveder, Y. F. Li, Q. Y. Liu and H. W. Long, Phys. Rev. D 83, 073006 (2011) [arXiv:1101.2755 [hep-ex]].
[47] A. A. Aguilar-Arevalo et al. [BICEP2 Collaboration], Phys. Rev. D 83, 073006 (2011) [arXiv:1101.2755 [hep-ex]].
[48] M. Antonello, B. Baibussinov, P. Benetti, E. Calligaris,
[51] J. Hamann, S. Hannestad, G. G. Raffelt and Y. Y. Y. Wong, JCAP 1109, 034 (2011) [arXiv:1108.4136 [astro-ph.CO]].
[52] E. K. Akhmedov, hep-ph/0001264.
[53] S. Hannestad, R. S. Hansen and T. Tram, Phys. Rev. Lett. 112 (2014) 031802 [arXiv:1310.5926 [astro-ph.CO]].
[54] B. Dasgupta and J. Kopp, Phys. Rev. Lett. 112 (2014) 031803 [arXiv:1310.6337 [hep-ph]].
[55] T. Bringmann, J. Hasenkamp and J. Kersten, arXiv:1312.4947 [hep-ph].
[56] M. Fukugita and T. Yanagida, Phys. Lett. B 174, 45 (1986).
[57] S. Baek, P. Ko and W. -I. Park, JHEP 1307, 013 (2013) [arXiv:1303.4280 [hep-ph]].
[58] B. Batell, M. Pospelov and A. Ritz, Phys. Rev. D 79 (2009) 115019 [arXiv:0903.3396 [hep-ph]].
[59] S. Baek, P. Ko, W. -I. Park and Y. Tang, arXiv:1402.2115 [hep-ph].
[60] J. L. Feng, M. Kaplinghat and H. -B. Yu, Phys. Rev. D 82, 083525 (2010) [arXiv:1003.4678 [hep-ph]].
[61] A. Loeb and M. Zaldarriaga, Phys. Rev. D 71 (2005) 103520 [astro-ph/0504112].
[62] E. Bertschinger, Phys. Rev. D 74 (2006) 063509 [astro-ph/0607319].
[63] J. Beringer et al. [Particle Data Group Collaboration], Phys. Rev. D 86, 010001 (2012).
[64] J. Kopp, M. Maltoni and T. Schwetz, Phys. Rev. Lett. 107, 091801 (2011) [arXiv:1103.4570 [hep-ph]].
[65] M. Archidiacono, N. Fornengo, C. Giunti and A. Melchiorri, Phys. Rev. D 86, 065028 (2012) [arXiv:1207.6515 [astro-ph.CO]].
[66] J. Kopp, P. A. N. Machado, M. Maltoni and T. Schwetz, JHEP 1305, 050 (2013) [arXiv:1303.3011 [hep-ph]].
[67] M. G. Aartsen et al. [IceCube Collaboration], arXiv:1309.7007 [astro-ph.HE].