Feebly Interacting $U(1)_{B-L}$ Gauge Boson Warm Dark Matter and XENON1T Anomaly

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The recent observation of an excess in the electronic recoil data by the XENON1T detector has drawn many attentions as a potential hint for an extension of the Standard Model (SM). Absorption of a vector boson with the mass of $m_{A^\prime} \in (2\text{ keV},3\text{ keV})$ is one of the feasible explanations to the excess. In the case where the vector boson explains the dark matter (DM) population today, it is highly probable that the vector boson belongs to a class of the warm dark matter (WDM) due to its suspected mass regime. In such a scenario, providing a good fit for the excess, the kinetic mixing $\kappa \sim 10^{-15}$ asks for a non-thermal origin of the vector DM. In this letter, we consider a scenario where the gauge boson is nothing but the $U(1)_{B-L}$ gauge boson and its non-thermal origin is attributed to the decay of the scalar talking to the SM sector via a portal with the SM Higgs boson. We discuss implications for the dark sector interactions that the vector DM offers when it serves as a resolution to both the small scale problems that ΛCDM model encounters and the XENON1T anomaly.

I. INTRODUCTION

Recently, the XENON1T collaboration reported an excess in the electronic recoil data for the energy regime ranging from 1 keV to 7 keV [1]. Especially, the prominence of the excess for $m_{A^\prime} \in (2\text{ keV},3\text{ keV})$ aroused many interesting interpretations based on various extensions of the SM. Absorption of a vector boson is one of the plausible possibilities for the excess in which case its suspected mass and kinetic mixing read $m_{A^\prime} \in (2\text{ keV},3\text{ keV})$ and $\kappa \sim 10^{-15}$ respectively [2–5]. Provided this vector boson serves as a dominant component of DM today, its suspected mass regime could be of interest in regard to the small scale problems (e.g. core/cusp problem [6], missing satellite problem [7, 8], too-big-to-fail problem [9]); keV scale WDM can alleviate some of the small scale problems if the free-streaming length travelled by the WDM amounts to $O(0.1)\text{ Mpc}$ [10].

Note that, however, the dark photon ($A^\prime_\mu$) mass $m_{A^\prime} \in (2\text{ keV},3\text{ keV})$ is actually outside of the allowed thermal WDM mass regimes inferred from Lyman-α forest observation $m_{\text{wdm,thermal}} > 5.3\text{ keV}$ [11] and redshifted 21cm signals in EDGES observations $m_{\text{wdm,thermal}} > 6.1\text{ keV}$ [12, 13]. Therefore in order for the dark photon to be a candidate for WDM to address the small scale problems, it should be a non-thermally originated one. Very interestingly, this line of reasoning to have a non-thermal dark photon DM (DPDM) gathers momentum when it is taken into account that a fit of good quality for the excess can be accomplished for the kinetic mixing $\kappa \sim 10^{-15}$ [1–4]. Being the most direct and dangerous coupling to thermalize $A^\prime_\mu$ with the SM thermal bath, the observed kinetic mixing $\kappa \sim 10^{-15}$ ensures that $A^\prime_\mu$ can avoid to join the SM thermal bath unless there is another significant indirect coupling to the SM sector. Now there arises an interesting question: what could be a non-thermal production mechanism that allows the mass regime $m_{A^\prime} \in (2\text{ keV},3\text{ keV})$ consistent with the aforementioned constraints on $m_{\text{wdm,thermal}}$? As we shall discuss, an answer to this question concerns momentum space distribution of $A^\prime_\mu$ which decides the map from $m_{\text{wdm,thermal}}$ to $m_{A^\prime}$. Different non-thermal production mechanisms will be characterized by different momentum spaces for $A^\prime_\mu$. Namely, given the same constraint on $m_{\text{wdm,thermal}}$, this implies that we would have different lower bounds for $m_{A^\prime}$, depending on what kind of non-thermal production mechanism produces $A^\prime_\mu$.

In this letter, we study the scenario where the keV-scale DPDM with the suppressed kinetic mixing arises with a non-thermal origin. Especially for $m_{A^\prime} \in (2\text{ keV},3\text{ keV})$, the DPDM can induce the XENON1T anomaly through absorption analogous to the photoelectric effect. We identify $A^\prime_\mu$ with the massive gauge boson of the broken $U(1)_{B-L}$ gauge theory which is the most well-motivated minimal extension of the SM. When incorporated with a scalar and three heavy right handed neutrinos, $U(1)_{B-L}$ gauge theory provides us with the most elegant explanation for the origin of the smallness of the active neutrino masses based on the seesaw mechanism [16–18] and the leptogenesis [19, 20]. In spite of this concreteness of the model, our mechanism and result can be easily generalized to a hidden broken $U(1)$ gauge theory, emphasizing its usefulness in the study of the DPDM. To enable the observed mass regime $m_{A^\prime} \in (2\text{ keV},3\text{ keV})$ to be consistent with the

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† Other examples of the use of $U(1)_{B-L}$ gauge symmetry for addressing the XENON1T anomaly can be found in Refs [14, 15].
given constraint on $m_{\text{wdm}}^{\text{thermal}}$, we consider the production mechanism where the non-relativistic scalar charged under $U(1)_{B-L}$ produces $A'_\mu$ via its decay after $U(1)_{B-L}$ gets spontaneously broken.$^2$ The parent scalar particle is assumed to be produced non-thermally from the SM thermal bath. We show the mapping between $m_{\text{wdm}}^{\text{thermal}}$ and $m_{A'}$ explicitly, corroborating that the production mechanism of our interest indeed provides a converted experimental constraint on $m_{A'}$ making mass regime $m_{A'} \in (2 \text{keV}, 3\text{keV})$ survive. By computing the free-streaming length based on the thermal history of $A'_\mu$ in the model, we further show presence of the parameter space in which $A'_\mu$ becomes the WDM candidate resolving some of the small scale problems.

II. MODEL

We consider the broken $U(1)_{B-L}$ gauge theory of which the massive gauge boson $A'_\mu$ is taken to be the DM candidate in the model. On top of the SM particle contents, we introduce a scalar $\Phi$ (-2) and three right-handed neutrino Weyl fields $\overline{N}_{i=1,2,3}$ (+1) where the numbers in the parenthesis denote $U(1)_{B-L}$ charges assigned to each field. These additional fields form the following Yukawa coupling

$$L_{\text{Yuk}} = \sum_{i=1}^{3} \frac{1}{2} y_{N,i,j} \Phi \overline{N}^{(i)} \overline{N}^{(j)} + \text{h.c.} ,$$

by which the right handed neutrinos acquire masses when the condensation of $\Phi$ induces the spontaneous breaking of $U(1)_{B-L}$. Stemming from the large vacuum expectation value (VEV) of $\Phi$, the heaviness of these right handed neutrino fields can explain the tiny active neutrino masses via the seesaw mechanism [16–18]. Moreover, the out-of-equilibrium decay of the heavy right-handed neutrinos creates the lepton asymmetry of the universe which is converted into the baryon asymmetry later with the help of the sphaleron transition [19]. For the consistency of the model shown to be later, we consider the reheating temperature $T_{\text{RH}} \gtrsim 10^{12}\text{GeV}$.

We assume that two of mass eigenvalues of $\Phi$ are comparable to $<\Phi> \equiv V_{B-L}/\sqrt{2}$ while the mass of the remaining one lies between $T_{\text{RH}}$ and $V_{B-L}$ ($T_{\text{RH}} < V_{B-L}$). Then at the reheating era, the lightest $\overline{N}$ serves as the source of the SM thermal bath via its decay into leptons and the SM Higgs (non-thermal leptogenesis) [20, 34].

Since $A'_\mu$ in the model is identified with the hypothetical vector boson triggering the XENON1T excess, we set $m_{A'} = 2g_{B-L}V_{B-L} \approx 2 - 3\text{keV}$ where $g_{B-L}$ is the gauge coupling of $U(1)_{B-L}$. Since we assume $V_{B-L} \gtrsim 10^{14}\text{GeV}$ for the leptogenesis, we further set $g_{A'} \lesssim 10^{-20}$. On the other hand, the scalar sector of the model is described by the following potential

$$V_{\text{scalar}} = -m_{\Phi}^2 |\Phi|^2 + \lambda |\Phi|^4 + \lambda_{\Phi H} H^* H |\Phi|^2 + V(H) ,$$

where $H$ denotes the SM SU(2) Higgs doublet and $V(H)$ is the Higgs potential in the SM. We see that there appear two new dimensionless couplings $\lambda$ and $\lambda_{\Phi H}$ as compared to the SM scalar sector. In the coming next sections, we study how these new couplings enable $A'_\mu$ to play a role of the non-thermally produced WDM which can resolve the small scale problems and explain the XENON1T anomaly.

III. PRODUCTION OF THE DARK PHOTON

At the reheating era with $T_{\text{RH}} \gtrsim 10^{12}\text{GeV}$, the scattering process $H + H^* \rightarrow \phi + \phi^*$ based on the Higgs portal interaction in Eq. (2) can produce the radial mode of $\Phi$ non-thermally for a sufficiently small $\lambda_{\Phi H}$. Here $\phi$ is defined via $\Phi \equiv (\phi/\sqrt{2} + V_{B-L}/\sqrt{2}) e^{i\theta}/v_{\text{av}}$ with $\theta$ Nambu-Goldstone mode that is absorbed by $A'_\mu$ on the spontaneous $U(1)_{B-L}$ breaking. We shall take $\phi$ as the parent particle of which decay produces a pair of $A'_\mu$ when it becomes non-relativistic.

For our purpose of having $A'_\mu$ as a non-thermal WDM candidate, we figured out that the model is consistent with a large enough $\lambda$ which enables $\phi$ to form the dark thermal bath. The tiny gauge coupling and the small Higgs portal coupling ($\lambda_{\Phi H}$) keeps the purity of the dark thermal bath until $\phi$ becomes non-relativistic so that the dark thermal bath is made up of $\phi$ only. $\phi$ avoids to be thermalized by the SM plasma since it is produced from the scattering process $H + H^* \rightarrow \phi + \phi^*$ if the associated interaction rate continues to be smaller than the Hubble expansion rate, i.e. $\Gamma < H$ until the SM Higgs particles are integrated out. As we shall see in the discussion of the DM relic density, we found that for the reheating temperature we consider in the model, $\lambda_{\Phi H}$ is sufficiently small not to thermalize $\phi$ until $T_{\text{SM}} \simeq 100\text{GeV}$ is reached. Below $T_{\text{SM}} \simeq 100\text{GeV}$, both the SM Higgs and $\phi$ do not exist and thus the scattering process $H + H^* \rightarrow \phi + \phi^*$ cannot take place physically.

As soon as $\phi$s are produced via the Higgs portal at the reheating era, due to the self-quartic interaction ($\lambda$) shown in Eq. (2), the dark thermal bath forms via the scattering process $\phi + \phi \rightarrow \phi + \phi$ if the associated interaction rate wins against the Hubble expansion rate. $^3$

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$^2$ A variety of ways other to produce the vector DM non-thermally has been suggested: resonant decays from an axion-like scalar via a tachyonic instability [21–24], production from a dark Higgs via a parametric resonance [25], decays from a network of the cosmic strings [26], vector coherent oscillation [27–30], and the gravitational production [31–33].

$^3$ As for non-thermal productions of dark sector particles via the SM particle scattering processes, see Refs. [35–37].
where \( \xi \equiv T_{\text{DS}} / T_{\text{SM}} \) is used. Eq. (3) gives us the maximum allowed reheating temperature \( T_{\text{SM, max}}(a_{\text{RH}}) \sim \lambda^2 \xi M_P \) below which \( \Gamma_{\phi\rightarrow 2\phi} > H \) always holds. As will be shown later, the correct DM relic abundance matching gives us \( \xi \simeq 0.5 \). Apart from Eq. (3), when the dark thermal bath formation requires \( \phi \) to be relativistic at the reheating era. Thus, we take the following hierarchy as the condition for formation of the dark thermal bath

\[
(\lambda^2 \xi M_P) \xi > T_{\text{DS}}(a_{\text{RH}}) > m_{\phi},
\]

where \( T_{\text{DS}} \) is the dark thermal bath temperature. The parameter space meeting the condition in Eq. (4) is shown as the area below the green dotted line in Fig. 1.

Thereafter, \( \phi \) in the dark thermal bath continues to experience redshifting to become the non-relativistic particle when \( m_{\phi} \simeq T_{\text{DS}} \) is reached.\(^4\) Thereafter, when the time for the \( \phi \)'s decay rate \( \Gamma_{\phi \rightarrow A^{' +} A^{' -}} \) to be comparable to the Hubble expansion rate is reached, the non-relativistic \( \phi \) starts to decay to a pair of \( A^{' +} \)'s. Since then, \( A^{' +} \) begins free-streaming as the WDM. By equating the decay rate \( \Gamma_{\phi \rightarrow A^{' +} A^{' -}} \) to the Hubble expansion rate during radiation dominated era, we obtain the temperature at which the free-streaming of \( A^{' +} \) begins

\[
\Gamma_{\phi \rightarrow A^{' +} A^{' -}} \simeq \frac{Q_{\phi}^4 g_{B-L} V_{B-L}^2 m_{\phi}^3}{32 \pi m_{A^{' +}}} \simeq \frac{T_{\text{SM}}(a_{\text{FS}})}{M_P}
\]

\( \Leftrightarrow T_{\text{SM}}(a_{\text{FS}}) \simeq 2.2 \times 10^8 \times \sqrt{\lambda} \times \left( \frac{m_{\phi}}{1 \text{GeV}} \right)^{1/2} \text{GeV} \) (5)

where we used \( m_{\phi} = \sqrt{2} \lambda V_{B-L} \), \( Q_{\phi} \) is the \( U(1)_{B-L} \) charge of \( \phi \), and \( a_{\text{FS}} \) is the scale factor for the onset of the free-streaming of \( A^{' +} \). For the second line in Eq. (5), we used \( Q_{\phi} = -2 \). To make it sure that the time corresponding to \( a_{\text{FS}} \) is preceded by the time when \( \phi \) becomes non-relativistic, we impose the condition \( m_{\phi} > \xi T_{\text{SM}}(a_{\text{FS}}) \), which results in

\[
m_{\phi} > \xi^2 \times (4.8 \times 10^{16} \times \lambda) \text{ GeV}.
\]

where we referred to Eq. (5) for \( T_{\text{SM}}(a_{\text{FS}}) \). This condition is satisfied by the region above the red dotted line in Fig. 1.

\( \text{ IV. DARK PHOTON AS THE WDM} \)

Here we study physical quantities that characterize the dark photon \( A^{' +} \) as the WDM. These quantities include the relic density, \( \Delta N_{\text{eff}} \) contributed by \( A^{' +} \), the free-streaming length \( L_{A^{' +}} \) of \( A^{' +} \); and the current constraint on the mass of \( A^{' +} \) based on the Lyman-\( \alpha \) forest observation.

\( \text{A. Relic density} \)

We assume that \( A^{' +} \) explains the whole of DM population today. The fraction of the energy density of the universe today attributed to the DM reads

\[
\Omega_{\text{DM},0} = \frac{\rho_{\text{DM},0}}{\rho_{c,0}} = \frac{s_0}{\rho_{c,0}} \times m_{\text{DM}} \times Y_{\text{DM}} \simeq 0.27,
\]

where the comoving number density of DM, \( Y_{\text{DM}} \equiv n_{\text{DM}} / s \), is the conserved quantity since the production of the DM. From the values of \( s_0 = 2.21 \times 10^{-10} \text{eV}^3 \) and \( \rho_{c,0} = 8.02764 \times 10^{-11} \times h^2 \text{ eV}^4 \), one obtains

\[
Y_{\text{DM}} = 9.8 \times 10^{-4} \times \left( \frac{m_{\text{DM}}}{1 \text{ keV}} \right)^{-1} \times h^{-2},
\]

where \( h \) is defined to be \( h_0 = 100 \text{ km}/\text{sec}/\text{Mpc} \). On the other hand, having the decay process \( \phi \rightarrow A^{' +} + A^{' -} \) as the origin of DM in the model, regarding number densities of particles, we have the relation \( n_{A^{' +}} = 2 n_{\phi} \). Given the non-thermally produced \( \phi \) from the SM Higgs scattering, the comoving number density of DM discussed in Eq. (8) can be written as \([36]\)

\[
Y_{\text{DM}} \equiv \frac{n_{A^{' +}}}{\Delta N_{\text{DM}}} = \frac{2 n_{\phi}}{s_{\text{SM}}} \sim 2 n_{\text{H}} \Gamma / H \bigg|_{T = T_{\text{RH}}} 
\]

\[
\simeq 2.7 \times 10^9 \times \lambda_{H}^2 \times \left( \frac{T_{\text{RH}}}{10^9 \text{GeV}} \right)^{1/2} \times \left( \frac{m_{A^{' +}}}{1 \text{ keV}} \right)^{-1/2}. \tag{9}
\]

Thus, using \( h = 0.68 \), equating Eq. (8) and Eq. (9) gives \( \lambda_{H} \simeq 8.86 \times 10^{-7} \times \left( \frac{T_{\text{RH}}}{10^9 \text{GeV}} \right)^{1/2} \times \left( \frac{m_{A^{' +}}}{1 \text{ keV}} \right)^{-1/2} \).\(^{10}\)

\(^{4}\) Note that \( A^{' +} \) cannot join the dark thermal bath by scattering among \( \phi \)s due to the small \( g_{B-L} \). This allows \( \phi \) to decay after it becomes non-relativistic, which is the key aspect that makes \( A^{' +} \)'s mass regime inferred from the XENONsurvive against the Lyman-\( \alpha \) forest constraint on the WDM mass. The similar non-thermal WDM production mechanism was also used in Refs. 3, 37, 48.
Note that for a given $T_{\text{RH}} \gtrsim 10^{12}\text{GeV}$, we see that $\lambda_{T\Phi}$ in Eq. (10) is still small enough to ensure that $\Gamma_{T\Phi,\gamma,\rightarrow\phi} < H$ holds until $H$ becomes not energetic enough to produce $\phi$s via the scattering and $\phi$ disappears in the dark sector via its decay.

Using the forms of $n_\phi$ and $s_{\text{SM}}$ as functions of $T_{\text{DS}}$ and $T_{\text{SM}}$, we can write $Y_{\text{DM}}$ given in Eq. (9) as the function of $\xi$. By equating the result to $Y_{\text{DM}}$ inferred from Eq. (8), we obtain $\xi \simeq 0.5$ for the mass regime of $m_{\lambda} \in (2\text{keV}, 3\text{keV})$.

B. $\Delta N_{\text{eff}}$ contributed by the dark photon

As the keV-scale DM candidate, $\lambda_{\mu}$ is the relativistic particle during the BBN era, behaving as the radiation. Therefore, its energy density contributes to the background expansion during the radiation dominated era, which is parametrized by its contribution to the extra effective number of neutrinos, $\Delta N_{\text{eff}}$. The amount of energy density of $\lambda_{\mu}$ at BBN time is given by

$$\rho_{\text{DM}}(a_{\text{BBN}}) = \sqrt{\frac{m_{\mu}^2}{2a_{\text{BBN}}} + \left(\frac{m_{\Phi}}{2a_{\text{BBN}}}\right)^2} Y_{\text{DM}} \times \frac{2\pi^2}{45} g_{s,\text{SM}}(a_{\text{BBN}}) T_{\text{SM}}(a_{\text{BBN}})^3,$$  \hspace{1cm} (11)

where $Y_{\text{DM}}$ is given in Eq. (8) and $g_{s,\text{SM}}$ is the effective number of relativistic degrees of freedom for the entropy density.\footnote{By referring to Eq. (5), we obtain $a_{\text{FS}}$}

$$\Delta N_{\text{eff}} \simeq \rho_{\text{DM}}(a_{\text{BBN}}) / \rho_{s}(a_{\text{BBN}}) \times 8 \left(\frac{11}{4}\right)^{4/3}. \hspace{1cm} (13)$$

From the constraint on $\Delta N_{\text{BBN}} \leq 0.364$ [39] (95% C.L.), Eq. (11) and Eq. (13), we obtain

$$m_{\Phi} \lesssim (1.32 - 2.96) \times 10^{19} \lambda \text{GeV}, \hspace{1cm} (14)$$

for $m_{\lambda} \in (2\text{keV}, 3\text{keV})$.

C. Free-streaming length of the dark photon

As the WDM candidate, $\lambda_{\mu}$ can alleviate the small scale problems provided its free-streaming length amounts to $\lambda_{FS} \simeq 0(0.1)\text{Mpc}$ [10].\footnote{In Ref. [10], WDM is characterized by $0.01\text{Mpc} \lesssim \lambda_{FS} \lesssim 0.1\text{Mpc}$.} Thus, we focus on the parameter space which can produce $\lambda_{FS} \simeq O(0.1)\text{Mpc}$ to make $\lambda_{\mu}$ DPDM WDM candidate addressing the small scale problems.

The free-streaming length is computed by

$$\lambda_{FS} = \int_{t_{FS}}^{t_0} \frac{<v_{\text{DM}}(t)> dt}{a} = \int_{a_{FS}}^{1} \frac{da}{H(a) \sqrt{(p_{\text{DM}}(a_{FS})/a_{FS})^2 + m_{\text{DM}}^2 a^2}}, \hspace{1cm} (15)$$

where $F(a) \equiv \sqrt{\Omega_{\text{rad},0} + a^{4} \Omega_{\mu,0} + a^4 \Omega_{\Lambda,0}}$, $a_{FS}$ is given in Eq. (12) and $\langle p_{\text{DM}}(a_{FS})/a_{FS}\rangle \simeq m_{\phi}/2$ in the instantaneous decay limit. Since $a_{FS}$ is a function of $(\lambda, m_{\phi})$, so is $\lambda_{FS}$ in our model.

D. $\lambda_{\mu}$ mass constraint from the Lyman–α forest

From the observation that WDMs of different origins have similar effects on the linear matter power spectrum if their velocity variances ($\beta$) are close to each other [40],\footnote{In references we cite here, the velocity variance is denoted by $\sigma$. But here we use $\beta$ on behalf of $\sigma$.} one can map the mass constraint on the thermal WDM from Lyman–α forest observation to a mass constraint on a non-thermally originated WDM [41]. The velocity variance is defined to be

$$\beta \equiv \frac{\beta_{\text{DM}}}{m_{\text{DM}}} \text{ with } \beta \equiv \frac{\int dqq^4 f(q)}{\int dqq^2 f(q)}, \hspace{1cm} (16)$$

where $f(q)$ is a momentum space distribution for the DM of interest. In our model, $\lambda_{\mu}$ is produced from the non-relativistic particle’s decay. As such, $\lambda_{\mu}$ is characterized by the momentum space distribution of the form $f(q,t) = (B(q)/q(\mu)^2)$ with $q \equiv p/T$. [40, 42–45]. Here $B$ is a normalization factor. This yields $\beta_{\text{DM}} \simeq 1$ while $\beta_{\text{thermal}} \simeq 3.6$ for the fermionic thermal WDM. On the other hand, the thermal WDM’s temperature today is given by [46]

$$T_{\gamma,0}^{\text{therm},0} \simeq \left[0.036 \left(\frac{94\text{eV}}{m_{\text{WM}}^{\text{therm}}}\right)^{1/3} T_{\gamma,0}^{\text{therm}}, \hspace{1cm} (17)$$

where $T_{\gamma,0} \simeq 2.35 \times 10^{-4}\text{eV}$ is the photon temperature today. In addition, using Eq. (12), we obtain DM temperature today

$$T_{\lambda,0} \simeq \frac{m_{\phi}}{2} \left[\frac{m_{\phi}}{1\text{GeV}}\right]^{1/2} \text{GeV}. \hspace{1cm} (18)$$

Eventually equipped with the aforementioned information above, by equating $\beta_{\lambda}$ with $\beta_{\text{thermal}}$, we obtain the lower
bound of $m_{A'}$\(^8\)

$$m_{A',lb} = \left(1\right) \left(\frac{T_{A',0}}{T_{\text{thermal}}}\right) m_{\text{thermal,lb}}. \quad (19)$$

when we enter the lower bound (lb) of the thermal WDM mass into $m_{\text{thermal,lb}}$. Given $m_{\text{thermal}} > 5.3$ keV from Lyman-\(\alpha\) forest observation [11], we show the parameter space resulting in the lower bound of $m_{A'}$ smaller than 3 keV as the region below the blue dotted line in Fig. 1. (Applying the conservative bound $m_{\text{thermal}} > 2$ keV [46, 47], we also obtain the area below the cyan dotted line in Fig. 1.)\(^9\) This confirms that the observed mass regime $m_{A'} \in (2$ keV, 3 keV) is consistent with Lyman-\(\alpha\) forest observation if $A'_\mu$ is the WDM produced based on the mechanism we discussed in our model.

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\(^8\) Note that comparison of the velocity variance $\beta$ (warmness) of the thermal WDM and $A'_\mu$ should be done at the matter-radiation equality time. Since both WDM decouple at a time earlier than the matter-radiation equality time, $T_{A'/\text{wdm}} = T_{\text{thermal}} - a_{\text{eq}}^{-1}$. Since $m_{\text{thermal}} > 2$ keV, the mass region below the cyan dotted line in Fig. 1 is consistent with Lyman-\(\alpha\) forest observation if $A'_\mu$ is the WDM produced based on the mechanism we discussed in our model.

\(^9\) The same thing can be done for the case of 2 keV.

E. Combined constraints on $V(\Phi)$

In Fig. 1, for the exemplary value of $m_{A'} = 3$ keV, we show the resultant parameter space of $(m_\Phi, \lambda)$ obtained by applying the constraints on $\Delta N_{\text{eff,BBN}}$ and the free-streaming criterion. The gray shaded area shows the region where $\Delta N_{\text{eff,BBN}} < 0.364$ is satisfied. Each of the solid lines of different colors is a set of points on $(m_\Phi, \lambda)$ plane producing each specified free-streaming length of $A'_\mu$ based on Eq. (15). The region below the green dotted line makes the dark thermal bath formation possible while the region above the red dotted line makes it sure that $\phi$ becomes non-relativistic before it decays to $A'_\mu$'s. Finally, the region below the dotted blue line produces the lower bounds of $m_{A'}$ smaller than 3 keV via the mapping of the lower bound on thermal WDM provided by the Lyman-\(\alpha\) forest data, i.e. $m_{\text{thermal}} > 5.3$ keV [11].

When combined together, indeed the applied three constraints denoted by the dotted lines produce the overlapped region for $\lambda \gtrsim 10^{-2}$ and $m_\Phi \gtrsim 10^{14}$ GeV. The points residing in the overlapped region enables $A'_\mu$ to resolve the small scale problems as the WDM by producing $\lambda_{\text{FS}} \approx (0.05 - 0.08)\text{Mpc}$. As such, $A'_\mu$ in the model is really shown to be a non-thermally originated WDM causing XENON1T anomaly, consistent with the various existing cosmological constraints.

V. DISCUSSION

In this letter, motivated by the recently reported XENON1T excess, we propose a minimal model where the massive gauge boson of $U(1)_{B-L}$ gauge theory with the kinetic mixing $\sim 10^{-15}$ can play the role of the dark photon dark matter with its mass $m_{A'} \in (2$ keV, 3 keV) inducing XENON1T anomaly. As the DM candidate with the free-streaming length of order $\sim \mathcal{O}(0.1)\text{Mpc}$, $A'_\mu$ is classified as the WDM which can address the small scale problems that $\Lambda$CDM is suffering from.

Relying on $U(1)_{B-L}$ gauge theory which is the well motivated extension of the SM, we introduced a scalar and three heavy right-handed neutrinos charged under $U(1)_{B-L}$ in addition to the SM particle content, which is the minimal set-up for ensuring the successful operation of the seesaw mechanism and leptogenesis. Imposing masses to the three-right handed neutrinos, the scalar was shown to be able to serve as the parent particle for the dark photon warm dark matter. Due to the non-thermal production mechanism we assumed in the model, the gauge boson mass of $m_{A'} \in (2$ keV, 3 keV) can be consistent with the Lyman-\(\alpha\) forest observation. Intriguingly, we figured out that for the self-quartic interaction $\lambda \gtrsim 10^{-2}$ and the scalar mass $m_\Phi \gtrsim 10^{14}$ GeV, the model can produce the non-thermally originated dark photon warm dark matter consistent without running a foul of existing cosmological and astrophysical constraints as a source of XENON1T anomaly. Although...
we focus on $U(1)_{B-L}$ gauge theory, our DM production mechanism can be easily applied to a hidden Abelian gauge theory.

**Note added:** After we finished this paper, we became aware of arXiv:2007.02898 [hep-ph] [48]. However, the mechanism to make the $U(1)_{B-L}$ gauge boson DM there is different from ours. See also references in their paper for proposals related to XENON1T anomaly.

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