Contributions to ΔN_{eff} from the dark photon of $U(1)_{T3R}$

Bhaskar Dutta,^{1,*} Sumit Ghosh^{1,†} and Jason Kumar^{2,‡}

¹*Mitchell Institute for Fundamental Physics and Astronomy, Department of Physics and Astronomy,
Texas A&M University, College Station, Texas 77843, USA*

²*Department of Physics, University of Hawaii, Honolulu, Hawaii 96822, USA*

(Received 10 February 2020; accepted 30 June 2020; published 16 July 2020)

We consider the effect on early Universe cosmology of the dark photon associated with the gauging of $U(1)_{T3R}$, a symmetry group under which only right-handed Standard Model fermions transform nontrivially. We find that cosmological constraints on this scenario are qualitatively much more severe than on other well-studied cases of a new $U(1)$ gauge group, because the dark photon couples to chiral fermions. In particular, the dark photon of $U(1)_{T3R}$ is always produced and equilibrates in the early Universe, no matter how small the gauge coupling, unless the symmetry-breaking scale is extremely large. This occurs because, no matter how the weak the coupling, the Goldstone mode (equivalently, the longitudinal polarization) does not decouple. As a result, even the limit of an extremely light and weakly coupled dark photon of $U(1)_{T3R}$ is effectively ruled out by cosmological constraints, unless the symmetry-breaking scale is extremely large. We also discuss the possibility of ameliorating the Hubble tension in this model.

DOI: [10.1103/PhysRevD.102.015013](https://doi.org/10.1103/PhysRevD.102.015013)

I. INTRODUCTION

In recent times, precise measurements of the cosmic microwave background (CMB) by the Planck experiment have placed tight constraints on the number of effective relativistic degrees of freedom in the early Universe, encoded in the quantity ΔN_{eff} [1]. These constraints can rule out models of new physics with new low-mass particles. Recent work has considered the constraints imposed on models of new physics in which a low-mass dark photon (A') couples to Standard Model (SM) (see, e.g., [2–7]). These works have focused on scenarios in which the dark photon is either secluded (coupling to SM particles only via kinetic mixing) or couples to the charges $B - L$ or $L_i - L_j$ [8–10]. But another well-studied anomaly-free choice of new $U(1)$ gauge group is $U(1)_{T3R}$; in this scenario, only one or more complete generations of right-handed SM fermions are charged, with up-type and down-type fermions having opposite charge. This scenario was originally considered in the context of left-right models [11–13], in which $U(1)_{T3R}$ is the diagonal subgroup of $SU(2)_R$, under which right-handed fermions transform as

doublets. Recently, $U(1)_{T3R}$ has been investigated for the purpose of building a well-motivated model of sub-GeV dark matter [14]. This model explains the hierarchies among the light fermion masses and contains a light gauge boson and a light scalar particle. In this brief paper, we point out that the dark photon of $U(1)_{T3R}$ contributes to ΔN_{eff} in a manner which is qualitatively different than the dark photon of other well-studied examples, such as $B - L$, $L_i - L_j$, a secluded $U(1)$, etc.

In these other well-studied examples, there are generally two ways in which one can ensure that the contribution of the dark photon to ΔN_{eff} is negligible; either the dark photon can be heavy enough that its abundance is negligible due to Boltzmann suppression at the time of neutrino decoupling, or its coupling can be so weak that it is never produced in the early Universe, again leading to a negligible abundance. But if the dark photon is the gauge boson of $U(1)_{T3R}$, then this second option is foreclosed; the dark photon is always produced in the early Universe, no matter how weak the coupling unless the symmetry-breaking scale is $\gtrsim 10^6$ GeV.

This result might at first seem counterintuitive. But one way to see this is to note that $m_{A'} \propto gV$, where g is the gauge coupling and V is the expectation value of the field which breaks the $U(1)$ gauge symmetry. Thus, for fixed V , as the gauge coupling gets weaker, the mass of the dark photon becomes smaller. If we then consider an inverse decay process like $\bar{f}f \rightarrow \gamma A'$, the sum over A' polarizations yields a factor of $-(g^{\mu\nu} - k^\mu k^\nu / m_{A'}^2)$, where the second term arises due to contribution of the longitudinal polarization. The $m_{A'}^2$ factor in the denominator cancels the g^2

*dutta@physics.tamu.edu

†ghosh@tamu.edu

‡jkumar@hawaii.edu

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

factor in the squared matrix element, potentially leaving a finite term even at arbitrarily small coupling.

Of course, these considerations apply for any choice of $U(1)$. For any choice of the $U(1)$ gauge group, $m_{A'} \propto gV$, and the longitudinal polarization thus always receives an enhancement which is proportional to $1/g$. But the enhanced term in the polarization sum is $\propto k^\mu k^\nu$; in cases where A' couples to SM fermions through a purely vector interaction, the resulting term in the matrix element is zero due to the Ward identity. However, if the gauge group is $U(1)_{T3R}$, then the longitudinal polarization is contracted with a combination of vector and axial vector currents, and the axial vector term does not vanish. This feature causes the A' production from the SM fermions to be nonzero even in the limit where the A' gauge coupling is taken to be very small.

Another way to see this result is to note that, in the weakly coupled limit, the $U(1)$ gauge group essentially becomes a global symmetry group, and the transverse polarizations of the A' manifestly decouple. But the longitudinal polarization instead becomes the massless Goldstone mode of the spontaneously broken global symmetry, which need not decouple. Again, these considerations apply for any choice of the $U(1)$ gauge group. But the relevant question is how does the Goldstone mode couples to SM fermions. The coupling of the Goldstone mode derives from the complex scalar whose vacuum expectation value (VEV) breaks the $U(1)$ symmetry; the Goldstone is the real excitation orthogonal to direction of the symmetry-breaking VEV. Since an unbroken $U(1)_{T3R}$ would forbid a SM fermion mass, the coupling of the Goldstone boson to any SM fermion charged under $U(1)_{T3R}$ must scale as m_f/V . But if the dark photon instead couples to $B - L$ or $L_i - L_j$, there is no reason why the symmetry-breaking field need to have a sizeable coupling to SM fields at the era of neutrino decoupling. As a result of these considerations, we will find that the scenario in which $U(1)_{T3R}$ is gauged is much more tightly constrained by cosmological observations than other recently studied scenarios. We will see explicitly that these stringent constraints emerge when the relativistic new gauge bosons are produced directly from on-shell muons. As a result, if only second-generation fermions are charged under $U(1)_{T3R}$, then collider and other astrophysical constraints are largely unaffected by these considerations, whereas constraints arising from early Universe cosmology become much more severe.

II. PRODUCTION OF A' IN THE EARLY UNIVERSE

For simplicity, we assume that only second-generation right-handed fermions are charged under $U(1)_{T3R}$, with up-type and down-type fermions having opposite charge ($Q_{c_R} = Q_{\nu_R} = 1$, $Q_{s_R} = Q_{\mu_R} = -1$). One can verify that

this choice is anomaly free. We will assume that $g \ll 1$, where g is the coupling of $U(1)_{T3R}$. In that case, the dominant processes by which A' can be produced in the early Universe are inverse decay processes, in which only one factor of g appears in the matrix element. In [5], it was argued that the dominant production process is $\bar{\mu}\mu \rightarrow \gamma A'$. For our purpose, it will be sufficient to consider this process in order to demonstrate that A' is always produced in the early Universe, provided this process is kinematically allowed and the symmetry-breaking scale is not extremely large.

The relevant Lagrangian for the gauge boson A' is

$$\mathcal{L} = -\frac{1}{4}B_{\mu\nu}B^{\mu\nu} + ig(\phi\partial_\mu\phi^* - \phi^*\partial_\mu\phi)A'^\mu + g^2\phi\phi^*A'_\mu A'^\mu - g\sum_f Q_f \bar{f}_R \gamma_\mu f_R A'^\mu, \quad (1)$$

where $B_{\mu\nu} = \partial_\mu A'_\nu - \partial_\nu A'_\mu$, ϕ is complex scalar field charged under $U(1)_{T3R}$, and $\langle\phi\rangle = V$. The condensation of ϕ spontaneously breaks $U(1)_{T3R}$, giving the dark photon a mass $m_{A'}^2 = 2g^2V^2$.

We may express the excitation of ϕ about its VEV in terms of two real fields, ϕ' and ϕ_I , yielding $\phi = V + (1/\sqrt{2})\phi' + (i/\sqrt{2})\phi_I$. ϕ' is the dark Higgs and is a physical real scalar excitation. ϕ_I is the Goldstone mode, which is absorbed by dark photon in order to provide the third physical polarization of A' .

The matrix element for the process $\bar{f}(p_2)f(p_1) \rightarrow \gamma(k_2)A'(k_1)$ is given by

$$i\mathcal{M}_{A'} = -i\frac{eQ_f m_{A'}}{\sqrt{2}V}\bar{v}(p_2)\left[\frac{\gamma^\mu \not{k}_2 \gamma^\nu - 2p_2^\mu \gamma^\nu}{-2p_2 \cdot k_2 + k_2^2} + \frac{2p_1^\mu \gamma^\nu - \gamma^\nu \not{k}_2 \gamma^\mu}{-2p_1 \cdot k_2 + k_2^2}\right] \times \frac{1+\gamma^5}{2}u(p_1)\epsilon_\nu^*(k_1)\epsilon_\mu^*(k_2), \quad (2)$$

where $\epsilon(k_1)$ and $\epsilon(k_2)$ are the polarization vectors of the A' and γ , respectively. The $P_R = (1 + \gamma^5)/2$ projector appears because A' only couples to f_R .

One can easily verify that the matrix element vanishes under the replacement $\epsilon^\mu(k_2) \rightarrow k_2^\mu$, as required by the Ward Identity. But one can also verify that, under the replacement $\epsilon^\nu(k_1) \rightarrow k_1^\nu$, the only nonvanishing term is proportional to the one proportional to γ^5 . This is also a result of the Ward identity. If the γ^5 term had been removed, then the coupling of f to A' would have been a pure vector interaction, and contracting the external momentum into the vector current necessarily yields zero.

This result immediately indicates that, in the case where the A' coupling to SM fermions is a pure vector interaction, the longitudinal polarization yields no parametric enhancement to the matrix element. The squared matrix element is contracted with an A' polarization sum factor given by $-(g^{\mu\nu} - k_1^\mu k_1^\nu/m_{A'}^2)$. In the weak

coupling limit ($m_{A'}/V \propto g \rightarrow 0$), the second term receives a parametric enhancement, but vanishes identically when contracted into a purely vector current.

We are interested in squared matrix element in limit where $g \ll 1$. In this case, only the $k_1^\mu k_1^\nu / m_{A'}^2$ term in the polarization sum is relevant, as this is the only term which can yield a nonzero contribution which contracted with a matrix element that scales as g^2 . From the Ward identity, we see that we need only consider the term in the matrix element proportional to γ^5 . Summing over the polarizations of the A' , we thus find

$$\begin{aligned} & \sum_{A' \text{ pols}} |\mathcal{M}_{A'}|^2 \\ &= \left(\frac{em_f}{2\sqrt{2}V} \right)^2 \left| \bar{v}(p_2) \left[\frac{\gamma^\mu k_1}{p_2 \cdot k_2} - \frac{2p_1^\mu - k_2 \gamma^\mu}{p_1 \cdot k_2} \right] \gamma^5 u(p_1) \epsilon_\mu^*(k_2) \right|^2, \end{aligned} \quad (3)$$

where we have set $k_2^2 = 0$ and $Q_f = -1$. It is thus clear that a finite piece is left, even in the limit $g \rightarrow 0$, when the dark photon couples to a chiral fermion.

One can verify this result straightforwardly by considering the limit where $g = 0$, in which case the A' is exactly massless, and $U(1)_{T3R}$ becomes effectively a global symmetry. In this case, the transverse polarizations of the A' must decouple, but the coupling of the massless Goldstone mode should reproduce the above squared matrix element. Indeed, this intuition is easily verified. The coupling of the Goldstone mode to f is induced from the coupling of the symmetry-breaking field ϕ to f , which is required in order for the fermion mass to be generated from a gauge-invariant Yukawa coupling. In the effective field theory defined below the electroweak symmetry-breaking scale, we find

$$\begin{aligned} \mathcal{L}_{\text{yuk}} &= \lambda_f \phi \bar{f} \left(\frac{1 + \gamma^5}{2} \right) f + \lambda_f \phi^* \bar{f} \left(\frac{1 - \gamma^5}{2} \right) f \\ &= m_f \bar{f} f + \frac{m_f}{\sqrt{2}V} \phi' \bar{f} f + i \frac{m_f}{\sqrt{2}V} \phi_I \bar{f} \gamma^5 f, \end{aligned} \quad (4)$$

implying that the Goldstone mode ϕ_I couples to f as a pseudoscalar with coupling $m_f/\sqrt{2}V$.

It is then straightforward to compute the squared matrix element for the process $\bar{f}(p_2)f(p_1) \rightarrow \gamma(k_2)\phi_I(k_1)$, yielding

$$\begin{aligned} & |\mathcal{M}_{\text{Gold}}|^2 \\ &= \left(\frac{em_f}{2\sqrt{2}V} \right)^2 \left| \bar{v}(p_2) \left[\frac{\gamma^\mu k_1}{p_1 \cdot k_1} - \frac{2p_1^\mu - k_2 \gamma^\mu}{p_1 \cdot k_2} \right] \gamma^5 u(p_1) \epsilon_\mu^*(k_2) \right|^2. \end{aligned} \quad (5)$$

In the limit $m_{A'} = 0$, we find $p_1 \cdot k_1 = p_2 \cdot k_2$, implying that the cross section for producing the massless A' in the

weakly coupled limit is equal to the cross section for producing the massless Goldstone boson, as required by the Goldstone equivalence theorem.

From here on, it is convenient to proceed in the Goldstone limit, where we take $g = 0$. If we choose simple kinematics for the incoming SM fermions, $p_1^\mu = (E, \vec{p})$, $p_2^\mu = (E, -\vec{p})$, defining $p = |\vec{p}|$, we find

$$\sigma v = \frac{\alpha_{em} m_f^2}{4E^2 V^2} \left[\frac{(2E^2 + p^2) \tanh^{-1}(p/E)}{Ep} - 1 \right]. \quad (6)$$

As expected, the cross section scales as $\alpha m_f^2/V^2$, since the coupling of the Goldstone mode to f is inherited from the coupling of the symmetry-breaking field, which necessarily scales as m_f/V , since $U(1)_{T3R}$ protects the fermion mass. We find that the thermally averaged cross section is given by

$$\langle \sigma v \rangle_{T \sim m_f} \sim 0.18 \frac{\alpha_{em}}{V^2}. \quad (7)$$

To determine the range of V for which A' equilibrates in the early Universe, we explicitly solve the Boltzmann equation for the A' abundance. But following [5], we find an approximate criterion for A' to not have equilibrated in the early Universe,

$$\eta_{f,\bar{f}}(T = m_f) \langle \sigma v \rangle_{T \sim m_f} \lesssim H = \sqrt{\frac{g_* \rho_{\text{rad}}(T = m_f)}{3M_{pl}^2}}, \quad (8)$$

where M_{pl} is the reduced Planck mass and g_* is the effective number of Standard Model relativistic degrees of freedom at $T = m_f$, yielding

$$\langle \sigma v \rangle_{T \sim m_f} \lesssim \frac{2.2 \sqrt{g_*}}{2m_f M_{pl}}. \quad (9)$$

We then find that A' will have equilibrated in the early Universe unless

$$\begin{aligned} V &\gtrsim \left(\frac{0.18 2\alpha_{em} m_f M_{pl}}{2.2 \sqrt{g_*}} \right)^{1/2}, \\ &\gtrsim (9 \times 10^6 \text{ GeV}) \left[\left(\frac{m_f}{m_\mu} \right) \left(\frac{g_*}{16.02} \right)^{-1/2} \left(\frac{\alpha_{em}}{1/137} \right) \right]^{1/2}. \end{aligned} \quad (10)$$

A solution of the Boltzmann equation yields a similar result.

III. ΔN_{eff}

Given that A' is produced and equilibrates in the early Universe, we must now determine how its abundance at the time of recombination corrects N_{eff} . For this purpose, we will assume that the neutrino mixing angle is small (the sterile neutrino mass eigenstate, ν_s , is almost entirely ν_R)

and that $m_{\nu_s} > 10$ MeV. If $m_{A'} > 10$ MeV, then the A' abundance is heavily Boltzmann suppressed at the time of neutrino decoupling, and its impact on N_{eff} is negligible [5].

In the limit $m_{A'}/V \rightarrow 0$, the transverse polarizations of the A' completely decouple, and we are left with a massless Goldstone mode, which thermalizes in the early Universe and decouples before neutrino decoupling, and which does not decay. As a result, the Goldstone degree of freedom is at the same temperature as the neutrinos, and its energy density at recombination contributes as $\Delta N_{\text{eff}} = 4/7$.

If $m_{A'}$ is non-negligible, but $m_{A'} < 1$ MeV, then the A' can decay to $\nu_A \nu_A$ through a one-loop process (decay to $\gamma\gamma$ is forbidden by the Landau-Yang Theorem). As the temperature drops well below $m_{A'}$, A' decays will heat the neutrino population, leading to an even larger value of ΔN_{eff} [5,6].

But if $m_{A'}$ lies in the range $\sim 1\text{--}10$ MeV, the analysis is model dependent. In particular, A' can also decay to e^+e^- through a one-loop kinetic-mixing process. The relative branching fractions for A' decay to $\nu_A\nu_A$ and e^+e^- are determined by the details of the neutrino mass matrix. This yields two relevant effects. First, electrons and neutrinos can remain coupled via decays and inverse decays of A' , delaying the time neutrino decoupling. As shown in [5], this can yield an $\mathcal{O}(1)$ correction to the allowed mass range for $m_{A'}$. But an even more significant effect arises if the branching fraction for $A' \rightarrow e^+e^-$ can be large. If the dominant decay of A' is to $\nu_A\nu_A$, then little changes from the above analysis. But if the dominant decay of A' is to e^+e^- , then when the temperature drops well below $m_{A'}$, the photon temperature increases, yielding a negative contribution to ΔN_{eff} . With an appropriate choice of branching fraction, ΔN_{eff} can be tuned to be arbitrarily small.

In Fig. 1(a), we plot the excluded region of parameter space in the $(m_{A'}, g)$ plane for the case where A' couples to $U(1)_{T3R}$ (blue), along with similar results from [5] (purple)

for the case where A' couples to $L_\mu - L_\tau$. Note, the A' abundance produced via inverse decay is computed by solving the Boltzmann equation. To facilitate comparison with [5], we will treat as excluded models for which $\Delta N_{\text{eff}} \geq 0.5$. The red dashed line indicates the parameter space for which A' will not fully equilibrate, yielding $\Delta N_{\text{eff}} \sim 0.2$. In Fig. 1(b), we plot the excluded regions of parameter space in the $(m_{A'}, V)$ plane. The larger V values correspond to regions where A' will not be in equilibrium.

In both Figs. 1(a) and 1(b), we show the excluded region of parameter space constrained by relevant fixed target experiments and different astrophysical processes. In electron beam dump experiments, such as SLAC E137 [15–18], A' can be produced via e -bremsstrahlung and decay to a e^+e^- pair through loop-suppressed mixing with photon. Due to the loop suppression, other fixed target experiments are irrelevant in the parameter space of our interest [19]. Note, E137 can only provide bounds in the model-dependent region, i.e., $m_{A'} > 1$ MeV. The A' can be produced inside the core of a supernova through the mixing with the photon, and can subsequently escape, resulting in energy loss. Constraints on this process are found by observing the energy loss of SN1987A [20,21]. The A' can also be produced in the Sun and can contribute to the solar cooling process. By requiring that the luminosity due to the dark photon be sufficiently small compared to the luminosity due to the photon, bounds can be derived [21,22]. Bounds can be found from the cooling of stars in Globular clusters in a similar way [21]. The green region in Figs. 1(a) and 1(b) shows the combined excluded region, considering the cooling of supernovae, the Sun and globular clusters. Note the parameter space for the lower mass range is not constrained by the astrophysical constraints, but is tightly constrained by the cosmological bounds we have found.

FIG. 1. The excluded regions of parameter space ($\Delta N_{\text{eff}} \geq 0.5$) in the $(m_{A'}, g)$ plane (left panel) and $(m_{A'}, V)$ plane (right panel). In blue is the excluded region if A' couples to $U(1)_{T_{3R}}$, under which second-generation Standard Model fermions are charged. For the case where A' couples to $L_\mu - L_\tau$, the excluded region in purple is reproduced from [5]. In both cases, the range $1 \text{ MeV} \leq m_{A'} \leq 10 \text{ MeV}$ is shaded, as exclusion contours in this mass range depend on details of the model.

Note the longitudinal polarization of A' has negligible effect on the collider and astrophysical bounds. The reason is because, in all of those cases, A' is produced through kinetic mixing with the photon, and its longitudinal polarization necessarily decouples. The more stringent constraints on $U(1)_{T3R}$ which arise from production of the longitudinal polarization come into play only when relativistic A' 's are produced directly from on-shell muons. As a result, these cosmological constraints are uniquely constraining.

It has been noted (see, e.g., [5,23–25]) that the tension between the determination of H_0 from low- z measurements [26,27] and from the CMB [1] can potentially be resolved if $\Delta N_{\text{eff}} \sim 0.2\text{--}0.5$. This range of ΔN_{eff} can arise in this model for a large range of $m_{A'}$. In Fig. 1, we show the parameter space where $\Delta N_{\text{eff}} \sim 0.2\text{--}0.5$ by solving the Boltzmann equation. In the model-dependent part of the parameter space, $m_{A'} \sim 1\text{--}10$ MeV, ΔN_{eff} can be set to $\sim 0.2\text{--}0.5$ by appropriately choosing the mixing between the active and sterile component which determines the branching fraction for A' decay to $\nu_A \nu_A$. In this case, ΔN_{eff} can receive both positive and negative contributions which can be tuned against each other by tuning the branching fraction for A' decay to $\nu_A \nu_A$ and $e^+ e^-$. For $m_{A'} \gtrsim 10$ MeV, we choose the gauge coupling appropriately to obtain the correct ΔN_{eff} .

IV. CONCLUSION

We have considered the effect of the dark photon of $U(1)_{T3R}$ on cosmology in the early Universe. We have found that, unlike other recently studied cases, such as $B-L$ and $L_i - L_j$, if the dark photon is the gauge boson of $U(1)_{T3R}$, cosmological constraints are much tighter. In particular, A' is always produced and equilibrates in the early Universe, not matter how small the gauge coupling is, provided the symmetry-breaking scale is $\lesssim 10^6$ GeV (for the case where second-generation right-handed fermions are charged under $U(1)_{T3R}$). Even if the gauge coupling is made arbitrarily small, this suppression of the A' production cross section is compensated by the enhancement of the longitudinal polarization when there is an axial vector coupling. This amounts to saying that, even in the limit when coupling becomes negligible and the symmetry becomes global, the Goldstone mode remains coupled to the charged fermions. We calculated ΔN_{eff} from the A' abundance by solving Boltzmann equation for this model and showed contours of $\Delta N_{\text{eff}} = 0.2, 0.5$ along with various constraints, e.g., collider, beam dump, cooling of

supernova, Sun and globular clusters, etc. We found that the cosmological constraints obtained in this work can exclude a large region of parameter space which is allowed by all other laboratory or astrophysical constraints.

We could consider the same scenario in the case where right-handed first-generation fermions are instead charged under $U(1)_{T3R}$. The considerations described above are largely unchanged; in this case, A' is produced and equilibrates in the early Universe unless the symmetry-breaking scale is $> 10^5$ GeV. One difference occurs if $m_{A'}$ lies in the 1–10 MeV range. In this case, assuming the sterile neutrino is heavy, one finds that the $A' \rightarrow \nu_A \nu_A$ decay process is one loop suppressed, while $A' \rightarrow e^+ e^-$ decay occurs at tree level. Thus, one would generally expect A' decay to inject energy into the photon gas, yielding a negative contribution to N_{eff} .

We see that regions of parameter space at very small $m_{A'}$ found in [14] are in fact in tension with cosmological constraints. In particular, this would rule out the scenarios described in [14] in which the dark photon coupled to electrons. Models in which $m_{A'} > 10$ MeV are still consistent with cosmological constraints, but if A' couples to right-handed electrons, then they are in tension with atomic parity violation experiments. But it may be possible to relax the tension with atomic parity violation experiments with a modest fine-tuning against additional sources of new physics; it would be interesting to investigate this further.

We also discussed the possibilities of ameliorating Hubble parameter measurements in this model which requires $\Delta N_{\text{eff}} \sim 0.2\text{--}0.5$. This range of ΔN_{eff} can arise in this model for a large range of $m_{A'}$. We showed that for $m_{A'} < 1$ MeV, some parts of the required ΔN_{eff} range are allowed by all other astrophysical constraints. In the model-dependent part of the parameter space, $m_{A'} \sim 1\text{--}10$ MeV, ΔN_{eff} can be set to $\sim 0.2\text{--}0.5$ by appropriately choosing the mixing between the active and sterile component and for $m_{A'} \gtrsim 10$ MeV, the gauge coupling can be appropriately chosen to obtain the required ΔN_{eff} .

We have focused in particular on the case where $U(1)_{T3R}$ is gauged. But the general result is valid in any scenario in which the dark photon has a chiral coupling to SM fermions. One would expect any such model to be tightly constrained by early Universe cosmology.

ACKNOWLEDGMENTS

The work of B. D. and S. G. are supported in part by the DOE Grant No. DE-SC0010813. The work of J. K. is supported in part by DOE Grant No. DE-SC0010504.

- [1] N. Aghanim *et al.* (Planck Collaboration), [arXiv:1807.06209](#).
- [2] A. Kamada and H.-B. Yu, *Phys. Rev. D* **92**, 113004 (2015).
- [3] A. Kamada, K. Kaneta, K. Yanagi, and H.-B. Yu, *J. High Energy Phys.* **06** (2018) 117.
- [4] S. Knapen, T. Lin, and K. M. Zurek, *Phys. Rev. D* **96**, 115021 (2017).
- [5] M. Escudero, D. Hooper, G. Krnjaic, and M. Pierre, *J. High Energy Phys.* **03** (2019) 071.
- [6] N. Sabti, J. Alvey, M. Escudero, M. Fairbairn, and D. Blas, *J. Cosmol. Astropart. Phys.* **01** (2020) 004.
- [7] R. Foot and S. Vagnozzi, *Phys. Rev. D* **91**, 023512 (2015).
- [8] R. Foot, *Mod. Phys. Lett. A* **06**, 527 (1991).
- [9] X. G. He, G. C. Joshi, H. Lew, and R. R. Volkas, *Phys. Rev. D* **43**, R22 (1991).
- [10] X.-G. He, G. C. Joshi, H. Lew, and R. R. Volkas, *Phys. Rev. D* **44**, 2118 (1991).
- [11] J. C. Pati and A. Salam, *Phys. Rev. D* **10**, 275 (1974); **11**, 703(E) (1975).
- [12] R. N. Mohapatra and J. C. Pati, *Phys. Rev. D* **11**, 2558 (1975).
- [13] G. Senjanovic and R. N. Mohapatra, *Phys. Rev. D* **12**, 1502 (1975).
- [14] B. Dutta, S. Ghosh, and J. Kumar, *Phys. Rev. D* **100**, 075028 (2019).
- [15] E. Riordan *et al.*, *Phys. Rev. Lett.* **59**, 755 (1987).
- [16] J. Bjorken, S. Ecklund, W. Nelson, A. Abashian, C. Church, B. Lu, L. Mo, T. Nunamaker, and P. Rassmann, *Phys. Rev. D* **38**, 3375 (1988).
- [17] J. D. Bjorken, R. Essig, P. Schuster, and N. Toro, *Phys. Rev. D* **80**, 075018 (2009).
- [18] S. Andreas, C. Niebuhr, and A. Ringwald, *Phys. Rev. D* **86**, 095019 (2012).
- [19] M. Bauer, P. Foldenauer, and J. Jaeckel, *J. High Energy Phys.* **07** (2018) 094.
- [20] J. B. Dent, F. Ferrer, and L. M. Krauss, [arXiv:1201.2683](#).
- [21] R. Harnik, J. Kopp, and P. A. Machado, *J. Cosmol. Astropart. Phys.* **07** (2012) 026.
- [22] J. Redondo, *J. Cosmol. Astropart. Phys.* **07** (2008) 008.
- [23] J. L. Bernal, L. Verde, and A. G. Riess, *J. Cosmol. Astropart. Phys.* **10** (2016) 019.
- [24] J. Alcaniz, N. Bernal, A. Masiero, and F. S. Queiroz, [arXiv:1912.05563](#).
- [25] S. Vagnozzi, [arXiv:1907.07569](#).
- [26] A. G. Riess *et al.*, *Astrophys. J.* **826**, 56 (2016).
- [27] A. G. Riess *et al.*, *Astrophys. J.* **861**, 126 (2018).