Gravitational Production of Hidden Photon Dark Matter in light of the XENON1T Excess

Kazunori Nakayama\textsuperscript{(a,b)} and Yong Tang\textsuperscript{(c,d,e)}

\textsuperscript{(a)}Department of Physics, Faculty of Science, The University of Tokyo, Bunkyo-ku, Tokyo 113-0033, Japan
\textsuperscript{(b)}Kavli IPMU (WPI), The University of Tokyo, Kashiwa, Chiba 277-8583, Japan
\textsuperscript{(c)}School of Astronomy and Space Sciences, University of Chinese Academy of Sciences (UCAS), Beijing, China
\textsuperscript{(d)}National Astronomical Observatories, Chinese Academy of Sciences, Beijing, China
\textsuperscript{(e)}School of Fundamental Physics and Mathematical Sciences, Hangzhou Institute for Advanced Study, UCAS, Hangzhou 310024, China

Abstract

Recently, the XENON1T experiment has reported an excess in the electronic recoil events. The excess is consistent with the interpretation of absorption of 3 keV bosonic dark matter, for example, hidden photon dark matter with kinetic mixing of the order of $10^{-15}$. We point out that the minimally gravitational production provides a viable mechanism for obtaining a correct relic hidden photon abundance. We present parameter dependence of the hidden photon dark matter abundance on the inflationary scale $H_{\text{inf}}$ and also the reheating temperature $T_R$. We show that the inflationary Hubble scale and reheating temperature are both bounded from below, $H_{\text{inf}} \gtrsim 7 \times 10^{11}$ GeV, $T_R \gtrsim 10^2$ GeV. In particular, the high-scale inflation is consistent with 3 keV hidden photon dark matter.
Recently, the XENON1T collaboration has reported excess events in the electronic recoil with the recoil energy around 2–7 keV [1]. The excess may be interpreted as a contribution from the solar axion [1], but this interpretation is inconsistent with the stellar cooling constraint, in particular the observation of white dwarfs and red giants [2]. On the other hand, various connections of the XENON1T excess with particle physics models, constraints and implications have been investigated in [3–18]. Absorption of bosonic dark matter (DM) may also give similar signals [19, 20]. Fig. 1 shows the signal event shape along with the XENON1T data points for the hidden photon mass $m_V = 2.7$ keV and the kinetic mixing parameter $\epsilon = 7 \times 10^{-16}$ (see Eq. (1) for definition), which apparently shows that the XENON1T excess events can be well fitted by the hidden photon DM. The indicated hidden photon mass $m_V \simeq 3$ keV and the kinetic mixing parameter $\epsilon \sim 10^{-15}$ is also shown to be consistent with the anomalous cooling of horizontal branch stars [5].

The viable production of keV DM is not trivial since the constraints from astrophysical observations are severe for keV DM produced from thermal plasma. For example, the recent Lyman-$\alpha$ gives the low bound $\gtrsim 5.3$ keV [21] if it is thermally produced. Hence, a keV DM candidate would require other viable production mechanism. In this short note, we focus on the hidden photon DM interpretation and show that the gravitational production mechanism is consistent with such a $\sim 3$ keV hidden photon DM with explicit parameter dependence on the inflationary energy scale and the reheating temperature.

The most relevant action of the hidden photon $V_\mu$ and the Standard Model (SM) elec-
tromagnetic photon $A_\mu$ for our discussions is given by

$$S = \int d^4x \sqrt{-g} \left( \frac{-1}{4} V_{\mu\nu} V^{\mu\nu} - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{2} m^2 V_\mu V^\mu + \epsilon V_{\mu\nu} F^{\mu\nu} \right)$$

$$= \int d^4x \sqrt{-g} \left( \frac{-1}{4} V_{\mu\nu} V^{\mu\nu} - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{2} m^2 V_\mu V^\mu \right),$$

where $V_{\mu\nu} \equiv \partial_\mu V_\nu - \partial_\nu V_\mu$ and $F_{\mu\nu} \equiv \partial_\mu A_\nu - \partial_\nu A_\mu$ are the field strength tensor of the hidden photon and SM photon, respectively. In the second line we have defined $	ilde{A}_\mu = A_\mu - \epsilon V_\mu$, $\tilde{V}_\mu = \sqrt{1 - \epsilon^2} V_\mu$ and $m^2 V_\mu = m V_\mu/\sqrt{1 - \epsilon^2}$.\(^\#1\) In the second basis, the kinetic terms are diagonalized while the SM particles are coupled with the hidden photon with coupling suppressed by $\epsilon$. The mass term is understood as a result of the Stuckelberg mechanism or the Higgs mechanism with the Higgs excitation being assumed to be heavy enough so that its dynamics is safely ignored. The massive hidden photon with kinetic mixing has rich phenomenological implications [22, 23].

For such a hidden photon to be DM, some production mechanisms are required. Conversion of thermal SM photons to the hidden photon does not lead to enough amount of relic hidden photon for the parameters we are interested in [24–26]. There are several proposed mechanisms so far: tachyonic instability due to the axionic scalar coupling to the hidden photon [27–29], production from the dark Higgs dynamics [30], from the cosmic strings associated with the spontaneous breaking of hidden $U(1)$ symmetry [31], vector coherent oscillation [32–37] and the gravitational production [38–40]. Among them we focus on the gravitational production mechanism since it is ubiquitous: such a contribution is unavoidable as far as we consider the inflationary universe and actually it is enough to reproduce the DM abundance for $m_V \sim \mathcal{O}(1)$ keV, as shown below.

The gravitational production of hidden photon DM was first discussed in Ref. [38] in the case of $m_V \ll H_{\text{inf}}$ with the assumption of instant reheating, where $H_{\text{inf}}$ denotes the Hubble scale during inflation. In Ref. [39] it was extended to the case of delayed reheating $H_R \ll H_{\text{inf}}$ and also the case of heavy hidden photon: $m_V \gtrsim H_{\text{inf}}$, where $H_R$ is the Hubble scale at the completion of reheating.

The hidden photon abundance from the gravitational production, in terms of the energy density-to-entropy density ratio, is given by [39]

$$\frac{\rho_{\text{HP}}}{s} \sim \begin{cases} 
3 \frac{m_V T_R H_{\text{inf}}}{2048\pi M_{\text{Pl}}^2} & \text{for } H_{\text{inf}} < m_V \\
1 \frac{T_R H_{\text{inf}}^2}{32\pi^2 M_{\text{Pl}}^2} & \text{for } H_R < m_V < H_{\text{inf}}, \\
\left( \frac{90}{\pi^2 g_*} \right)^{1/4} \frac{1}{32\pi^2} \frac{m_V^{1/2} H_{\text{inf}}^2}{M_{\text{Pl}}^3/2} & \text{for } m_V < H_R.
\end{cases}$$

\(^\#1\) If we start from the hypercharge photon, instead of the electromagnetic photon, and introduces the kinetic mixing of $\epsilon_Y$, we end up with (1) after the electroweak symmetry breaking with $\epsilon = \epsilon_Y \cos \theta_W$ where $\theta_W$ is the Weinberg angle.
where $M_{\text{Pl}}$ is the reduced Planck scale and we have defined the reheating temperature $T_R$ through $T_R = (90/\pi^2 g_*)^{1/4} \sqrt{H_R M_{\text{Pl}}}$. As usual, the universe is assumed to be matter-dominated during the reheating since the inflaton harmonic oscillation behaves as non-relativistic matter. Practically the case of $H_{\text{inf}} < m_V$ is irrelevant since it predicts too small hidden photon abundance and it cannot be DM for $m_V \sim O(1)$ keV. Below we briefly summarize how to obtain (3) for $m_V < H_{\text{inf}}$. Since the kinetic mixing is so tiny, we do not distinguish $V$ and $\bar{V}$ below. In the Appendix we will discuss small effect from nonzero $\epsilon$.

The massive hidden photon is decomposed into the transverse and longitudinal mode:

$$V_\mu(\vec{x}, \tau) = \int \frac{d^3k}{(2\pi)^3} V_\mu(\vec{k}, \tau) e^{i\vec{k} \cdot \vec{x}}; \quad \bar{V}(\vec{k}, \tau) = \bar{V}_T(\vec{k}, \tau) + \frac{\vec{k}}{|\vec{k}|} V_L(\vec{k}, \tau),$$

(4)

where the transverse mode satisfies $\vec{k} \cdot \bar{V}_T = 0$. The transverse action is

$$S_T = \int \frac{d^3k d\tau}{(2\pi)^3 \cdot 2} \left( |\bar{V}_T'|^2 - (k^2 + a^2 m_V^2) |\bar{V}_T|^2 \right),$$

(5)

and the longitudinal action is

$$S_L = \int \frac{d^3k d\tau}{(2\pi)^3 \cdot 2} \left( \frac{a^2 m_V^2}{k^2 + a^2 m_V^2} |\bar{V}_L'|^2 - a^2 m_V^2 |\bar{V}_L|^2 \right)$$

$$= \int \frac{d^3k d\tau}{(2\pi)^3 \cdot 2} \left( |\bar{V}_L'|^2 - \left[ k^2 + a^2 m_V^2 - \frac{k^2}{k^2 + a^2 m_V^2} \left( \frac{a''}{a} - \frac{a^2}{a^2 k^2 + a^2 m_V^2} \right) \right] |\bar{V}_L|^2 \right),$$

(6)

(7)

where $a$ denotes the cosmic scale factor and we have defined $\bar{V}_L \equiv \sqrt{a^2 m_V^2/(k^2 + a^2 m_V^2)} V_L$ and the prime denotes the derivative with respect to the conformal time $\tau$. The transverse mode production is negligible compared with the longitudinal one and we focus on the latter. The longitudinal action reduces to the same action as a minimal scalar for $k > a m_V$ and hence the quantum fluctuation in inflationary epoch results in $\langle |\bar{V}_L|^2 \rangle \sim a_{\text{end}}^2 H_{\text{inf}}^2 / (2k^3)(2\pi)^3 \delta(k - k')$ for the superhorizon modes $a_{\text{end}} H_{\text{inf}} < k < a_{\text{end}} m_V$ at the end of inflation, where $a_{\text{end}}$ denotes the scale factor at the end of inflation. Now we want to estimate the abundance when the Hubble parameter $H$ becomes equal to the hidden photon mass: $H = m_V$. It is found that the high frequency mode $k > k_* \equiv a_* m_V$ is redshifted away rather rapidly, where $a_*$ denotes the scale factor at $H = m_V$, while the low frequency mode $k < k_*$ has an initially suppressed energy density already at $a = a_{\text{end}}$. The dominant contribution to the final energy density comes from the mode $k \sim k_*$. Thus the energy density at $H = m_V$ is evaluated as

$$\rho_{\text{HP}}(a_*) \simeq \frac{1}{2} \left( \frac{k_*}{a_*} \right)^2 \left( \frac{H_{\text{inf}}}{2\pi} \right)^2 = \frac{m_V^2 H_{\text{inf}}^2}{8\pi^2}.$$ 

(8)

When $H < m_V$, the energy density scale as $\propto a^{-3}$ as an ordinary non-relativistic matter. It leads to the second and third line of (3).\textsuperscript{2}

\textsuperscript{2} This expression of the hidden photon energy density is similar to the scalar coherent oscillation energy density as if the initial amplitude of the scalar field is $H_{\text{inf}}/(2\pi)$. 

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relic density for \( m_V < H_R \) as

\[
\rho_{\text{HP}} \sim \rho_{\text{DM}} \times \frac{m_V}{3\,\text{keV}} \times \left( \frac{H_{\text{inf}}}{7 \times 10^{11}\,\text{GeV}} \right)^2,
\]

where \( \rho_{\text{DM}} \) is the average energy density of DM at present. When the reheating temperature after inflation is low, i.e. \( H_R < m_V < H_{\text{inf}} \), we have

\[
\rho_{\text{HP}} \sim \rho_{\text{DM}} \times \frac{T_R}{10^6\,\text{GeV}} \times \left( \frac{H_{\text{inf}}}{7 \times 10^{11}\,\text{GeV}} \right)^2,
\]

which is independent of the mass \( m_V \).

In Fig. 2, we illustrate how the relic density of hidden photon depends on the reheating temperature \( T_R \) and inflation scale \( H_{\text{inf}} \), when fixing the mass of hidden photon as \( m_V = 3\,\text{keV} \). The solid line gives the right relic abundance of DM, while the dotted and dashed lines correspond to ten times larger and one tenth smaller, respectively. In the large \( T_R \), the production is independent of \( T_R \), which shows that even the instant reheating is valid. When \( T_R \) is small enough such that \( H_R \) is smaller than the hidden photon’s mass, we would need large inflation scale \( H_{\text{inf}} \) to compensate the production loss because the relic density now depends on \( T_R \) linearly. The turning points of these curves occur around \( H_R \approx m_V \). Since the inflation scale \( H_{\text{inf}} \) is bounded from above, \( H_{\text{inf}} \lesssim 10^{14}\,\text{GeV} \), which is based on the non-observation of primordial gravitational waves by the Planck satellite [41], the correct relic abundance of hidden photon would give the lower bound on the Hubble scale and reheating temperature,

\[
H_{\text{inf}} \gtrsim 7 \times 10^{11}\,\text{GeV}, \quad T_R \gtrsim 10^2\,\text{GeV}.
\]

The reheating temperature is also bounded from above, \( T_R \lesssim 10^{15}\,\text{GeV} \), corresponding to the instant reheating limit.

Interestingly, the delayed reheating scenario \( (H_R \ll H_{\text{inf}}) \) opens up a possibility for high-scale inflation to be consistent with hidden photon DM. For maximally possible inflation scale \( H_{\text{inf}} \sim 10^{14}\,\text{GeV} \), the reheating temperature is predicted to be around the weak scale: \( T_R \sim 10^2\,\text{GeV} \). In this case, it is possible to probe the primordial gravitational waves through the observation of the B-mode polarization in the cosmic microwave background anisotropy. On the other hand, it is below the sensitivity of future space-based direct gravitational wave detectors due to too low \( T_R \) [42–44].

One of the good aspects of the hidden photon DM with gravitational production is that it is not constrained from the limit on the DM isocurvature fluctuation. In contrast to light scalar DM, the isocurvature fluctuation spectrum is strongly blue for the hidden photon DM [38] due to the peculiar behavior of the longitudinal component for \( k < a m_V \) and hence practically there is no effect on the cosmic microwave background anisotropy on the cosmological scale. Thus even the high-scale inflation does not suffer from the isocurvature constraint. Another aspect is that it is an unavoidable contribution since the gravity is universal. In this sense it gives lower bound on the hidden photon abundance. It is possible
Figure 2: The illustration of the energy density of dark photon $\rho_{\text{HP}}$ for $m_V = 3$ keV. The solid line shows how the correct relic abundance would require the proper values of inflation scale $H_{\text{inf}}$ and reheating temperature $T_R$. The dotted and dashed lines correspond to ten times larger and smaller, respectively. The turning points indicate $H_R \simeq m_V$.

that hidden photon has interactions with other sector and the production is more efficient, which, however, is highly model-dependent.

In summary, hidden photon DM with $m_V \sim 3$ keV and $\epsilon \sim 10^{-15}$ is a good candidate to explain the XENON1T excess events. The gravitational production works well for such a mass region for reasonable inflationary scale and the wide range of the reheating temperature. In particular, taking account of the effect of delayed reheating, the high-scale inflation can be consistent with the hidden photon DM scenario.

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Appendix

In this Appendix we show that the effect of small kinetic mixing $\epsilon$ on the gravitational hidden photon production is small enough. In the mass basis (2), the SM fermions are coupled to the hidden photon with a coupling of $O(\epsilon e)$ with $e$ denoting the SM gauge coupling constant. Therefore, at the quantum level, there is a trace anomaly that violates the conformality of the
massless vector boson and some amount of particle production is expected. The longitudinal component is already highly non-conformal even for $\epsilon = 0$ as highlighted in the main text, and hence we focus on the transverse mode below.

Let us suppose that the SM fermions are light during inflation. It has been found that the equation of motion of the transverse mode is modified as follows \cite{45,46}:

$$
\dddot{V}_T + \kappa \frac{a'}{a} \ddot{V}_T + \left( k^2 + a^2 m_V^2 \right) \dot{V}_T = 0,
$$

where $\kappa \simeq O(e^2 \epsilon^2)$. Following the standard procedure of the quantization we introduce the creation and annihilation operator that satisfies $[a_{\vec{k},h}, a^\dag_{\vec{k}',h'}] = (2\pi)^3 \delta(\vec{k} - \vec{k}') \delta_{hh'}$ and so on:

$$
\ddot{V}_T(\vec{k}, \tau) = \sum_{h=\pm} \left[ \mathcal{V}_T(\vec{k}, \tau) \epsilon_h a_{\vec{k},h} + \mathcal{V}_T^*(\vec{k}, \tau) \epsilon^*_h a^\dag_{\vec{k},h} \right],
$$

where $h = \pm$ denotes the two helicity states and the polarization vector satisfies $\epsilon_h \cdot \epsilon^*_h = \delta_{hh'}$. One can solve the equation during inflation with the Bunch-Davies boundary condition as

$$
\mathcal{V}_T(\vec{k}, \tau) = e^{i(2\nu + 1)\pi/4} \frac{1}{\sqrt{2k}} \sqrt{-\frac{\pi k \tau}{2}} H^{(1)}_{\nu}(k \tau), \quad \nu^2 \equiv \frac{(1 + \kappa)^2}{4} - \frac{m_V^2}{H_{\text{inf}}^2},
$$

where $H^{(1)}_{\nu}(x)$ denotes the Hankel function of the first kind. Since both $|\kappa|$ and $m_V^2/H_{\text{inf}}^2$ are much smaller than unity, we have $2\nu - 1 \simeq (\kappa - 2m_V^2/H_{\text{inf}}^2)$. It characterizes the two origins of the breaking of scale invariance, the quantum anomaly effect $\kappa$ and the mass term $m_V^2/H_{\text{inf}}^2$. If $|\kappa| \ll m_V^2/H_{\text{inf}}^2$, one can neglect $\kappa$ and in this case it is already known that the transverse mode production is much less efficient than the longitudinal one \cite{38,39}. Thus let us consider the opposite case, $\kappa \gg m_V^2/H_{\text{inf}}^2$. In this case, the spectrum at the end of inflation $\tau = \tau_{\text{end}}$ (or $k \tau \to 0$) is given by

$$
\left| \mathcal{V}_T(\vec{k}) \right|^2 \sim \frac{1}{k} \left( \frac{a_{\text{end}} H_{\text{inf}}}{k} \right)^\kappa.
$$

The energy density will be dominated by the high-momentum mode, $k \approx a_{\text{end}} H_{\text{inf}}$. Taking account of the subtraction of the UV divergence as usual, we finally obtain

$$
\rho_{\text{HF}}^{(T)} \sim \begin{cases} 
\kappa \frac{m_V T_R H_{\text{inf}}}{M_P^2} & \text{for } m_V < H_R \\
\kappa \frac{m_V^{3/2} H_{\text{inf}}}{M_P^{3/2}} & \text{for } m_V > H_R
\end{cases},
$$

Since we are interested in the parameter region of $\kappa = O(e^2 \epsilon^2) \sim 10^{-30}$, it is evident that it is much smaller than the longitudinal contribution (3).
References

[1] E. Aprile et al. [XENON], [arXiv:2006.09721 [hep-ex]].

[2] M. Giannotti, I. G. Irastorza, J. Redondo, A. Ringwald and K. Saikawa, JCAP 10, 010 (2017) [arXiv:1708.02111 [hep-ph]].

[3] F. Takahashi, M. Yamada and W. Yin, [arXiv:2006.10035 [hep-ph]].

[4] K. Kannike, M. Raidal, H. Veermae, A. Strumia and D. Teresi, [arXiv:2006.10735 [hep-ph]].

[5] G. Alonso-Alvarez, F. Ertas, J. Jaeckel, F. Kahlhoefer and L. Thormaehlen, [arXiv:2006.11243 [hep-ph]].

[6] B. Fornal, P. Sandick, J. Shu, M. Su and Y. Zhao, [arXiv:2006.11264 [hep-ph]].

[7] C. Boehm, D. G. Cerdeno, M. Fairbairn, P. A. Machado and A. C. Vincent, [arXiv:2006.11250 [hep-ph]].

[8] G. Paz, A. A. Petrov, M. Tammaro and J. Zupan, [arXiv:2006.12462 [hep-ph]].

[9] J. Buch, M. A. Buen-Abad, J. Fan and J. S. C. Leung, [arXiv:2006.12488 [hep-ph]].

[10] D. A. Sierra, V. De Romeri, L. Flores and D. Papoulias, [arXiv:2006.12457 [hep-ph]].

[11] G. Choi, M. Suzuki and T. T. Yanagida, [arXiv:2006.12348 [hep-ph]].

[12] N. F. Bell, J. B. Dent, B. Dutta, S. Ghosh, J. Kumar and J. L. Newstead, [arXiv:2006.12461 [hep-ph]].

[13] Y. Chen, J. Shu, X. Xue, G. Yuan and Q. Yuan, [arXiv:2006.12447 [hep-ph]].

[14] L. Di Luzio, M. Fedele, M. Giannotti, F. Mescia and E. Nardi, [arXiv:2006.12487 [hep-ph]].

[15] M. Du, J. Liang, Z. Liu, V. Tran and Y. Xue, [arXiv:2006.11949 [hep-ph]].

[16] L. Su, W. Wang, L. Wu, J. M. Yang and B. Zhu, [arXiv:2006.11837 [hep-ph]].

[17] A. Bally, S. Jana and A. Trautner, [arXiv:2006.11919 [hep-ph]].

[18] K. Harigaya, Y. Nakai and M. Suzuki, [arXiv:2006.11938 [hep-ph]].

[19] M. Pospelov, A. Ritz and M. B. Voloshin, Phys. Rev. D 78 (2008), 115012 [arXiv:0807.3279 [hep-ph]].

[20] K. Arisaka, P. Beltrame, C. Ghag, J. Kaidi, K. Lung, A. Lyashenko, R. Peccei, P. Smith and K. Ye, Astropart. Phys. 44 (2013), 59-67 [arXiv:1209.3810 [astro-ph.CO]].
[21] N. Palanque-Delabrouille, C. Yeche, N. Schoneberg, J. Lesgourgues, M. Walther, S. Chabanier and E. Armengaud, JCAP 04 (2020), 038 [arXiv:1911.09073 [astro-ph.CO]].

[22] J. Jaeckel and A. Ringwald, Ann. Rev. Nucl. Part. Sci. 60, 405 (2010) [arXiv:1002.0329 [hep-ph]].

[23] M. Fabbrichesi, E. Gabrielli and G. Lanfranchi, [arXiv:2005.01515 [hep-ph]].

[24] J. Redondo and M. Postma, JCAP 02, 005 (2009) [arXiv:0811.0326 [hep-ph]].

[25] H. An, M. Pospelov and J. Pradler, Phys. Lett. B 725, 190-195 (2013) [arXiv:1302.3884 [hep-ph]].

[26] J. Redondo and G. Raffelt, JCAP 08, 034 (2013) [arXiv:1305.2920 [hep-ph]].

[27] P. Agrawal, N. Kitajima, M. Reece, T. Sekiguchi and F. Takahashi, Phys. Lett. B 801, 135136 (2020) [arXiv:1810.07188 [hep-ph]].

[28] R. T. Co, A. Pierce, Z. Zhang and Y. Zhao, Phys. Rev. D 99, no. 7, 075002 (2019) [arXiv:1810.07196 [hep-ph]].

[29] M. Bastero-Gil, J. Santiago, L. Ubaldi and R. Vega-Morales, JCAP 1904, no. 04, 015 (2019) [arXiv:1810.07208 [hep-ph]].

[30] J. A. Dror, K. Harigaya and V. Narayan, Phys. Rev. D 99, no. 3, 035036 (2019) [arXiv:1810.07195 [hep-ph]].

[31] A. J. Long and L. Wang, Phys. Rev. D 99, no.6, 063529 (2019) [arXiv:1901.03312 [hep-ph]].

[32] A. E. Nelson and J. Scholtz, Phys. Rev. D 84, 103501 (2011) [arXiv:1105.2812 [hep-ph]].

[33] P. Arias, D. Cadamuro, M. Goodsell, J. Jaeckel, J. Redondo and A. Ringwald, JCAP 1206, 013 (2012) [arXiv:1201.5902 [hep-ph]].

[34] G. Alonso-Alvarez, J. Jaeckel and T. Hugle, JCAP 02, no.02, 014 (2020) [arXiv:1905.09836 [hep-ph]].

[35] K. Nakayama, JCAP 10, no.10, 019 (2019) [arXiv:1907.06243 [hep-ph]].

[36] K. Nakayama, [arXiv:2004.10036 [hep-ph]].

[37] Y. Nakai, R. Namba and Z. Wang, [arXiv:2004.10743 [hep-ph]].

[38] P. W. Graham, J. Mardon and S. Rajendran, Phys. Rev. D 93, no. 10, 103520 (2016) [arXiv:1504.02102 [hep-ph]].

[39] Y. Ema, K. Nakayama and Y. Tang, JHEP 07, 060 (2019) [arXiv:1903.10973 [hep-ph]].
[40] A. Ahmed, B. Grzadkowski and A. Socha, [arXiv:2005.01766 [hep-ph]].

[41] Y. Akrami et al. [Planck], [arXiv:1807.06211 [astro-ph.CO]].

[42] K. Nakayama, S. Saito, Y. Suwa and J. Yokoyama, Phys. Rev. D 77, 124001 (2008) [arXiv:0802.2452 [hep-ph]].

[43] K. Nakayama, S. Saito, Y. Suwa and J. Yokoyama, JCAP 06, 020 (2008) [arXiv:0804.1827 [astro-ph]].

[44] S. Kuroyanagi, T. Chiba and N. Sugiyama, Phys. Rev. D 79, 103501 (2009) [arXiv:0804.3249 [astro-ph]].

[45] A. Dolgov, Phys. Rev. D 48, 2499-2501 (1993) [arXiv:hep-ph/9301280 [hep-ph]].

[46] T. Prokopec, [arXiv:astro-ph/0106247 [astro-ph]].