Gauged $U(1)$ Clockwork

Hyun Min Lee[*][†]

Department of Physics, Chung-Ang University, Seoul 06974, Korea.
E-mail: hminlee@cau.ac.kr

We review the gauged $U(1)$ clockwork models as portals for new physics to obtain effectively hierarchical couplings of the Standard Model (SM) particles from genuine couplings of order one. Nearest-neighbor mass mixings between multiple $U(1)$ gauge bosons lead to the localization of the zero mode gauge boson, realizing the milicharged gauge boson in the effective theory. We discuss some of interesting examples utilizing the gauged clockwork beyond the SM.

Corfu Summer Institute 2017 'School and Workshops on Elementary Particle Physics and Gravity'
2-28 September 2017
Corfu, Greece

---

[*]Speaker.

1. Introduction

Particle physics experiments at intensity and energy frontiers have been complementary direct probes for new physics with light or heavy particles beyond the Standard Model. New symmetry for solving the Higgs mass hierarchy problem suggests new particles at weak scale to cancel the quadratic sensitivity to UV physics, but there is no direct hint for new physics yet. However, there are plenty of hints for new physics such as dark matter, strong CP problem, origin of baryon asymmetry, neutrino masses, etc. Searching in the unknown territory, minimal approaches for new physics are taken, such as portal models of the following type,

$$\mathscr{L}_{\text{portal}} = |H|^2 \mathscr{O}_{\text{new}}, \quad F_{\mu\nu}F'^{\mu\nu}, \quad N_R \mathscr{O}_{\text{new}}, \quad a_{\text{ALP}}F_{\mu\nu}F^{\mu\nu}, \quad \frac{1}{\Lambda}G_{\mu\nu}T^{\mu\nu}, \cdots \tag{1.1}$$

which are portals via Higgs, hypercharge gauge boson, right-handed neutrino, axion-like particle, massive spin-2 particle, etc.

Recently there have been new interests in the clockwork mechanism [1, 2], where hierarchical couplings between the mediators and the SM particles are naturally generated due to the nearest neighbor interactions of the mediators. In this framework, for order one couplings of similar strength to start with, the lightest mass eigenstate of the mediators can have exponentially small couplings to the SM particles due to the localization in the flavor space. The gravity dual of the clockwork scenarios in extra dimensional models with dilaton background have been discussed [2, 3, 4, 5, 6] and has drawn new attention recently. In this article, we review the basic idea of the clockwork mechanism and discuss the gauged $U(1)$ clockwork in specific and illustrate some interesting examples to resolve some of hierarchy problems in particle physics.

2. Clockwork mechanism and continuum limit

Suppose that there are multiple real scalar fields, $\phi_i (i = 0, 1, 2, \cdots, N)$, whose shift symmetries are broken into one shift symmetry, by the following nearest-neighbor interactions [1, 2],

$$\mathscr{L}_{SCW} = \frac{1}{2} \sum_{j=0}^{N} \partial_\mu \phi_i \partial^\mu \phi_j - \frac{m^2}{2} \sum_{j=0}^{N-1} (\phi_i - q\phi_{j+1})^2. \tag{2.1}$$

Then, the mass matrix for scalar fields is given by

$$M_\phi^2 = m^2 \begin{pmatrix} 1 & -q & 0 & \cdots & & 0 \\ -q & 1+q^2 & -q & \cdots & & 0 \\ 0 & -q & 1+q^2 & \cdots & & 0 \\ \vdots & \vdots & \vdots & \ddots & & \vdots \\ & & & & 1+q^2 & -q \\ 0 & 0 & 0 & \cdots & -q & q^2 \end{pmatrix}. \tag{2.2}$$

The Lagrangian (2.1) respects the unbroken shift symmetry, under $\phi_j \to \phi_j + \frac{c}{q^j}$. Then, solving the equations for ϕ_j leads to $\phi_N = \frac{1}{q}\phi_{N-1} = \cdots = \frac{1}{q^N}\phi_0$, corresponding to a sequential decrease of field

values in going from site 0 toward site N [1]. This is the so called clockwork mechanism and the fields in the flavor space are clock gears. Indeed, the zero mode has a profile, $\tilde{\phi}_0 \sim \sum_j \phi_j/q^j$, localized toward the site at $j = 0$, so it has position-dependent couplings to external fields, in particular, suppressed couplings to external fields containing interactions only to ϕ_0 by $\mathscr{L}_{\text{int}} = \phi_0 \mathscr{O}_{\text{ext}}$. On the other hand, the massive modes of the clockwork, $\tilde{\phi}_k (k \neq 0)$ with nonzero mass M_k, are given by the rest of linear combinations of the clock gears, with a mass gap m from the zero mode and the squeezed mass spectrum, $\delta M_k/M_k \sim 1/N$, for a large N [2].

One can take the continuum limit of the clockwork by introducing a five-dimensional scalar field, $\phi(x,y) = \phi_j(x)$, with $y = ja$ where a is the lattice distance. In the limit of $a \to 0$ and $N \to \infty$, we keep $\pi R \equiv Na$ finite. Then, additionally taking $m \to \infty$ with $ma \to 1$ and $q \to 1$, the clockwork Lagrangian (2.1) with interactions to external fields, $\mathscr{L}_{\text{int}} = \phi_0 \mathscr{O}_{\text{ext}}$, becomes [2]

$$\mathscr{L}_{5D,SCW} = \int_0^{\pi R} dy \left(\frac{1}{2}(\partial_\mu \phi)^2 - \frac{1}{2}(\partial_y \phi + k\phi)^2 \right) + \int_0^{\pi R} dy \, \delta(y - \pi R)\phi \, \mathscr{O}_{\text{ext}} \tag{2.3}$$

where $k \equiv (q-1)/(qa)$ and $q^N = e^{k\pi R}$. Then, the continuum clockwork can be described by a bulk scalar field ϕ with bulk and mass terms proportional to k [7], so the zero mode of ϕ has an exponential profile, $\phi_0 \sim e^{-ky}$, such that it has suppressed couplings to the external fields. With the field redefinition by $\chi = e^{ky}\phi$ and the dilaton background $S = -2ky$, the above Lagrangian becomes [3, 4]

$$\mathscr{L}_{5D,SCW} = \int_0^{\pi R} dy \, e^S \left[\frac{1}{2}\partial_M \chi \partial^M \chi \right] + \int_0^{\pi R} dy \, \delta(y - \pi R) e^{\frac{1}{2}S} \chi \, \mathscr{O}_{\text{ext}}. \tag{2.4}$$

Thus, one bulk scalar field χ does not have induced mass terms but it has position-dependent kinetic terms. In this picture, the zero mode of the bulk scalar field has a constant bulk profile so it has suppressed couples to the external fields via the dilaton factor.

We remark the five-dimensional realization for the continuum clockwork. The 5D gravity with a dilaton S [9, 2] was considered as in the following,

$$S = \int d^5x \sqrt{-g} \frac{M_5^3}{2} e^S \left(R + (\partial_M S)^2 + 4k^2 \right)$$
$$- \int d^5x \sqrt{-g} \, e^S \left(\frac{\delta(y)}{\sqrt{g_{55}}} \Lambda_0 + \frac{\delta(y - \pi R)}{\sqrt{g_{55}}} \Lambda_\pi \right) \tag{2.5}$$

where M_5 is the 5D Planck mass, k is the 5D curvature and $\Lambda_{0,\pi}$ are the brane tensions. Then, the scale symmetry with $S \to S + c$ and $g_{MN} \to e^{-2c/3} g_{MN}$ is broken explicitly by a nonzero k^2 as well as brane tensions. With a Weyl rescaling of the metric, $g_{MN} \to e^{2S/3} g_{MN}$, the above action becomes in Einstein frame

$$S = \int d^5x \sqrt{-g} \frac{M_5^3}{2} \left(R - \frac{1}{3}(\partial_M S)^2 + 4k^2 e^{-2S/3} \right)$$
$$- \int d^5x \sqrt{-g} \, e^{-S/3} \left(\frac{\delta(y)}{\sqrt{g_{55}}} \Lambda_0 + \frac{\delta(y - \pi R)}{\sqrt{g_{55}}} \Lambda_\pi \right) \tag{2.6}$$

---

[1]We can introduce the nearest-neighbor interactions by $(q\phi_i - \phi_{j+1})^2$ instead. In this case, the field values are related by $\phi_0 = \frac{1}{q}\phi_1 = \cdots = \frac{1}{q^N}\phi_N$ in the minimum of the potential so they increase in going from site 0 toward site N

Thus, the brane tensions couple to the dilaton in Einstein frame, stabilizing the radius of the extra dimension. A 5D dilaton background with 4D flat space was obtained in Einstein frame for the above action, as follows [8, 9, 2, 6],

$$ds^2 = e^{\frac{4}{3}k|y|}(\eta_{\mu\nu}dx^\mu dx^\nu + dy^2); \quad S = 2k|y|, \tag{2.7}$$

with the tuning relation between brane tensions,

$$\Lambda_\pi = -\Lambda_0 = 4kM_5^3. \tag{2.8}$$

In this model, the 4D effective Planck mass is given by

$$M_P^2 = M_5^3 \int_{-\pi R}^{\pi R} dy \sqrt{-g}\, e^{-\frac{4}{3}k|y|} = \frac{M_5^3}{3} L_5\, e^{\frac{1}{3}k\pi R} \tag{2.9}$$

with $L_5 \equiv \int_{-\pi R}^{\pi R} dy\, e^{\frac{4}{3}k|y|}$ being the proper length of the extra dimension. Thus, in the 5D dilaton background with branes, the hierarchy problem between the Planck scale and the weak scale can be solved in combination of the warp factor and the large fifth dimension. The bulk theory is known to be a holographic dual of Little String Theory [10, 8].

3. Gauged $U(1)$ clockwork

The clockwork mechanism can be simply applied to the case with gauge bosons for multiple local $U(1)$'s. We break $N+1$ copies of local $U(1)$'s, $U(1)_0 \times U(1)_1 \times \cdots \times U(1)_N$, down to one $U(1)$ by introducing gauge boson masses via Stueckelberg mechanism or Higgs mechanism [2, 3, 7], as follows,

$$\mathscr{L}_{GCW} = \sum_{j=0}^{N-1} \frac{1}{2} m^2 \left(A_\mu^j - q A_\mu^{j+1} + \frac{1}{m} \partial_\mu \pi^j \right)^2 \tag{3.1}$$

where π_j transforms under the gauge transformations as $\pi^j \to -m(\alpha^j - q\alpha^{j+1})$ for $A_\mu^j \to A_\mu^j + \partial_\mu \alpha^j$. Here, we note that the remaining $U(1)$ transforms by $A_\mu^j \to A_\mu^j + \frac{1}{q^j}\partial_\mu \alpha$ under which π^j is invariant. In the case with Higgs mechanism, we get $m^2 = g^2 f^2$ for nonzero VEVs of ϕ_j carrying charge $(1,-q)$ under $U(1)_j \times U(1)_{j+1}$, i.e. $\phi_j = \frac{1}{\sqrt{2}}(f + h_j)\, e^{i\pi^j/f}$ [7]. The Higgs perturbations of ϕ_j turn out to be decoupled in the continuum limit [7], but they should be kept for the renormalizability of the discrete clockwork.

Similarly as in the scalar clockwork [2], the zero mode of the gauge clockwork is localized at $j = 0$ for $q > 1$, given [2, 7] by

$$\tilde{A}_\mu^0 = \sum_{j=0}^{N} a_{j0} A_\mu^j(x) \tag{3.2}$$

where $a_{j0} = N_0/q^j$ with $N_0 = \sqrt{(q^2-1)/(q^2-a^{-2N})}$. The massive modes of the gauge clockwork are also given by

$$\tilde{A}_\mu^k(x) = \sum_{j=0}^{N} a_{jk} A_\mu^j(x), \quad k = 1, 2, \cdots, N, \tag{3.3}$$

with the mass eigenvalues being

$$M_k^2 = m^2 \left(1 + q^2 - 2q \cos \frac{k\pi}{N+1} \right) \equiv m^2 \lambda_k, \tag{3.4}$$

and the wave functions being

$$a_{jk} = N_k \left[q \sin \left(\frac{jk\pi}{N+1} \right) - \sin \left(\frac{(j+1)k\pi}{N+1} \right) \right], \quad N_k = \sqrt{\frac{2}{(N+1)\lambda_k}}. \tag{3.5}$$

The interacting gauge fields are are invertible to get

$$A_\mu^j(x) = \sum_{j=0}^N a_{jk} \tilde{A}_\mu^j(x), \quad j = 0, 1, 2, \cdots, N. \tag{3.6}$$

This relation is important to identify the gauged clockwork couplings to the external fields localized at a site.

We can also take the continuum limit of the gauged clockwork by taking $a \to 0$ and $f \to \infty$ and $q \to 1$ with $agfq = 1$. Then, with a five-dimensional gauge field, $A_\mu(x,y) \equiv A_\mu^j(x)$ and $A_y(x,y) \equiv \pi^j(x)$, the gauged clockwork Lagrangian becomes [7]

$$\mathscr{L}_{5D,GCW} = \int_0^{\pi R} dy \left[-\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{2} (\partial_y A_\mu - \partial_\mu A_y + k A_\mu)^2 \right] \tag{3.7}$$

with $k \equiv (q-1)/(qa)$ again. The resulting Lagrangian contains bulk and brane mass terms proportional to k as in the scalar clockwork, leading to the localization of the zero mode of the gauged clockwork [7] such as

$$\psi_0(y) = N_0 e^{-k|y|}, \quad N_0 = \sqrt{\frac{k\pi R}{1 - e^{-2k\pi R}}} \tag{3.8}$$

in the covering space S^1 with $-\pi R < y < \pi R$ where $-\pi R < y < 0$ is identified with $0 < y < \pi R$ by a Z_2 symmetry. Furthermore, the wave functions for massive modes are

$$\psi_n(y) = N_n \left(\cos \frac{ny}{R} - \frac{kR}{n} \sin \frac{n|y|}{R} \right), \quad N_n = \frac{n}{m_n R}, \tag{3.9}$$

with mass eigenvalues being

$$m_n^2 = k^2 + \frac{n^2}{R^2}. \tag{3.10}$$

The results can be comparable to the case with discrete clockwork in eqs. (3.4) and (3.5). For $kR \gg 1$, the warp factor of the zero mode, $e^{-y\pi R}$, gets small, so the spectrum of massive modes is squeezed at the scale of k.

With the field redefinition by $B_M = e^{ky} A_M$, the above Lagrangian can be rewritten as

$$\mathscr{L}_{5D,GCW} = \int_0^{\pi R} dy \, e^S \left[-\frac{1}{4} F_{MN} F^{MN} \right] \tag{3.11}$$

with $F_{MN} = \partial_M B_N - \partial_N B_M$. Thus, the continuum gauged clockwork is equivalent to a five-dimensional massless $U(1)$ theory with the dilaton background.

We consider the gauged clockwork couplings to external fields with nonzero $U(1)$ charges. General couplings to a fermion ψ localized at site l are given [7] by

$$\mathscr{L}_{\text{fermion}} = -g\bar{\psi}\gamma^{\mu}(v_{\psi} + a_{\psi}\gamma^{5})\psi A^{l}_{\mu}, \tag{3.12}$$

where v_{ψ}, a_{ψ} are vectorial and axial couplings, respectively, while those to a scalar ϕ localized at site p are given by

$$\mathscr{L}_{\text{scalar}} = -ig(\phi^{*}\partial^{\mu}\phi - \phi\partial^{\mu}\phi^{*})A^{p}_{\mu} + g^{2}(A^{p}_{\mu})^{2}|\phi|^{2}. \tag{3.13}$$

Then, from eq. (3.6), the gauged clockwork couplings to fermions or scalars can be easily identified. In particular, even if fermions or scalars carry order one charge under their local $U(1)$, the effective couplings to the zero mode of the gauged clockwork can be exponentially suppressed as $g_{\text{eff}} = N_{0}g/q^{l,p}$, realizing the milicharged $U(1)$.

Remarkably, when a dark Higgs ϕ at site p breaks the remaining $U(1)$ of the gauge clockwork by $\langle\phi\rangle = \frac{1}{\sqrt{2}}v_{\phi} \ll f$, the mass of the zero mode can be naturally suppressed by $M_{0}^{2} = \frac{g^{2}N_{0}^{2}}{q^{2p}}v_{\phi}^{2} \ll v_{\phi}^{2}$ [7]. From the point view of the continuum clockwork, a localized Higgs mechanism with nonzero VEV v_{IR} on the IR brane at $y = \pi R$ leads to the gauge boson mass which is much smaller than the IR scale by the warp factor as $M_{0} = e^{-k\pi R}gv_{\text{IR}} \ll v_{\text{IR}}$ [7].

4. Examples of gauged $U(1)$ clockwork

Weakly Interacting Massive Particles (WIMPs) have been the main paradigm for dark matter, having weak-scale mass and neutrino-like interactions to the SM particles. But, there has been no conclusive hint for WIMP dark matter yet and the bounds from direct detection experiments are getting more stringent. On the other hand, orbital velocity curves of galaxies hint at a core dark matter profile, unlike the cuspy profile obtained from the N-body simulation with WIMP dark matter, causing the cusp-core problem. Thus, together with too-big-to-fail problem, there are small-scale problems at galaxy scales. A large self-scattering cross section for dark matter about $\sigma_{\text{self}}/m_{\text{DM}} = 0.1 - 10\,\text{cm}^{2}/g$ would be required to solve the small-scale problems [11].

Various applications of the gauged clockwork have been proposed in Ref. [7], such as self-interacting dark matter with suppressed couplings to the SM particles, flavor-dependent $U(1)'$ interactions for explaining the B-meson anomalies in semi-leptonic decays at LHCb [12], and suppressed D-term supersymmetry breaking due to gauge kinetic mixings between multiple $U(1)$ gauge multiplets, etc.

Among those concrete models for the gauged clockwork, we consider the gauged clockwork as a mediator between dark matter and the SM particles to explain the smallness of interactions of dark matter to the SM particles while allowing for large self-interactions of dark matter, due to the localization of the light zero mode of a mediator gauge boson towards dark matter. There are anomaly-free $U(1)'$ without extra charged fermions beyond the SM. For instance, $U(1)_{B-L}$ and $U(1)_{L_{i}-L_{j}}$ with $i,j = 1,2,3$ and $i \neq j$ are such examples [7]. The LHC searches for extra gauge bosons such as dimuon searches, however, have put stringent constraints on the Z' mass for $U(1)'$ with SM hypercharge-like gauge coupling to be heavier than multi-TeV. Thus, in this case, thermal dark matter mediated by one of those gauge interactions would be heavy near the half the gauge boson mass and have very tiny self-interactions.

In the framework of the gauged clockwork, we introduce general $U(1)'$ interactions to a Dirac fermion dark matter χ at site 0 and the SM fermions f at site N [7], as follows,

$$\mathscr{L}_{Z'} = -g_{Z'}\bar{\chi}\gamma^\mu(v_\chi + a_\chi\gamma^5)\chi Z_\mu'^0 - g_{Z'}\bar{\chi}\gamma^\mu(v_f + a_f\gamma^5)\chi Z_\mu'^N \tag{4.1}$$

where $v_{\chi,f}$ and $a_{\chi,f}$ are vectorial and axial couplings for dark matter and SM fermions, respectively. Then, expanding the gears of the gauged clockwork in terms of mass eigenstates $\tilde{Z}_\mu^k$ as in eq. (3.6),

$$Z_\mu'^j(x) = \sum_{j=0}^{N} a_{jk}\tilde{Z}_\mu^j(x), \quad j = 0, 1, 2, \cdots, N, \tag{4.2}$$

eq. (4.1) becomes

$$\mathscr{L}_{Z'} = -g_{Z'}\bar{\chi}\gamma^\mu(v_\chi + a_\chi\gamma^5)\chi \left(N_0\tilde{Z}_\mu^0 + \sum_{k=1}^{N} a_{0k}\tilde{Z}_\mu^k(x) \right)$$
$$- g_{Z'}\bar{f}\gamma^\mu(v_f + a_f\gamma^5)f \left(\frac{N_0}{q^N}\tilde{Z}_\mu^0 + \sum_{k=1}^{N} a_{Nk}\tilde{Z}_\mu^k(x) \right). \tag{4.3}$$

As a consequence, for $m_\chi \ll M_0 \ll M_k$ with $k \neq 0$, we can integrate out the gauged clockwork to obtain the effective interactions for dark matter,

$$\mathscr{L}_{\text{eff}} = -\frac{N_0^2 g_{Z'}^2}{2M_0^2}\left(\bar{\chi}\gamma^\mu(v_\chi + a_\chi\gamma^5)\chi \right)^2 - \frac{N_0^2 g_{Z'}^2}{2q^{2N}M_0^2}\left(\bar{f}\gamma^\mu(v_f + a_f\gamma^5)f \right)^2$$
$$- \frac{N_0^2 g_{Z'}^2}{q^N M_0^2}\left(\bar{\chi}\gamma^\mu(v_\chi + a_\chi\gamma^5)\chi \right)\left(\bar{f}\gamma^\mu(v_f + a_f\gamma^5)f \right). \tag{4.4}$$

As a consequence, the effective self-interactions of dark matter can be large due to the light zero mode of the gauged clockwork, while the interactions between dark matter and SM fermions are suppressed for $q^N \gg 1$ due to a strong localization of the zero mode towards dark matter. Therefore, a self-interacting dark matter can be realized with small couplings (or effective milicharges) to the SM particles in the gauged clockwork models [7].

For a light zero mode of the gauged clockwork, the above effective Lagrangian is not valid and we need to include the light zero mode explicitly in the calculations of dark relic density, direct detection and self-scattering cross section, etc [7, 13]. Dark photon searches at LHCb, Belle II and SHiP experiments [14] could probe the zero mode of the gauged clockwork in relation to the self-interacting dark matter.

Furthermore, there are additional effective interactions induced due to massive states of the gauged clockwork, but they are negligible for $M_k \gg q^{N/2}M_0$ for $k \neq 0$. Otherwise, the sum of those massive states contribute sizably to the effective interactions [7, 15]. In this case, collider searches for massive states with squeezed spectrum [6] would be complementary for testing the gauged clockwork.

5. Conclusions

We have given a review on clockwork scenarios where multiple global or local symmetries are broken into a single symmetry due to nearest-neighbor interactions between multiple fields. We

discussed the gauged clockwork mechanism in more detail and showed that the localization of the zero mode in the gauged clockwork can explain the hierarchy of couplings in external fields for genuine couplings of order one. Various interesting applications of the clockwork scenarios are suggested in concrete models with extra $U(1)$ interactions.

References

[1] K. Choi and S. H. Im, JHEP **1601** (2016) 149 doi:10.1007/JHEP01(2016)149 [arXiv:1511.00132 [hep-ph]]; D. E. Kaplan and R. Rattazzi, Phys. Rev. D **93** (2016) no.8, 085007 doi:10.1103/PhysRevD.93.085007 [arXiv:1511.01827 [hep-ph]].

[2] G. F. Giudice and M. McCullough, JHEP **1702** (2017) 036 doi:10.1007/JHEP02(2017)036 [arXiv:1610.07962 [hep-ph]].

[3] N. Craig, I. Garcia Garcia and D. Sutherland, JHEP **1710** (2017) 018 doi:10.1007/JHEP10(2017)018 [arXiv:1704.07831 [hep-ph]].

[4] G. F. Giudice and M. McCullough, arXiv:1705.10162 [hep-ph].

[5] K. Choi, S. H. Im and C. S. Shin, arXiv:1711.06228 [hep-ph].

[6] G. F. Giudice, Y. Kats, M. McCullough, R. Torre and A. Urbano, arXiv:1711.08437 [hep-ph].

[7] H. M. Lee, Phys. Lett. B **778** (2018) 79 doi:10.1016/j.physletb.2018.01.010 [arXiv:1708.03564 [hep-ph]].

[8] I. Antoniadis, S. Dimopoulos and A. Giveon, JHEP **0105** (2001) 055 doi:10.1088/1126-6708/2001/05/055 [hep-th/0103033].

[9] M. Baryakhtar, Phys. Rev. D **85** (2012) 125019 doi:10.1103/PhysRevD.85.125019 [arXiv:1202.6674 [hep-ph]]; P. Cox and T. Gherghetta, JHEP **1205** (2012) 149 doi:10.1007/JHEP05(2012)149 [arXiv:1203.5870 [hep-ph]].

[10] O. Aharony, M. Berkooz, D. Kutasov and N. Seiberg, JHEP **9810** (1998) 004 doi:10.1088/1126-6708/1998/10/004 [hep-th/9808149]; O. Aharony, A. Giveon and D. Kutasov, Nucl. Phys. B **691** (2004) 3 doi:10.1016/j.nuclphysb.2004.05.015 [hep-th/0404016].

[11] S. Tulin and H. B. Yu, Phys. Rept. **730** (2018) 1 doi:10.1016/j.physrep.2017.11.004 [arXiv:1705.02358 [hep-ph]].

[12] L. Bian, S. M. Choi, Y. J. Kang and H. M. Lee, Phys. Rev. D **96** (2017) no.7, 075038 doi:10.1103/PhysRevD.96.075038 [arXiv:1707.04811 [hep-ph]]; L. Bian, H. M. Lee and C. B. Park, arXiv:1711.08930 [hep-ph].

[13] S. M. Choi and H. M. Lee, JHEP **1509** (2015) 063 doi:10.1007/JHEP09(2015)063 [arXiv:1505.00960 [hep-ph]]; S. M. Choi, Y. J. Kang and H. M. Lee, JHEP **1612** (2016) 099 doi:10.1007/JHEP12(2016)099 [arXiv:1610.04748 [hep-ph]].

[14] P. Ilten, Y. Soreq, J. Thaler, M. Williams and W. Xue, Phys. Rev. Lett. **116** (2016) no.25, 251803 doi:10.1103/PhysRevLett.116.251803 [arXiv:1603.08926 [hep-ph]]; R. Essig, J. Mardon, M. Papucci, T. Volansky and Y. M. Zhong, JHEP **1311** (2013) 167 doi:10.1007/JHEP11(2013)167 [arXiv:1309.5084 [hep-ph]]; S. Alekhin *et al.*, Rept. Prog. Phys. **79** (2016) no.12, 124201 doi:10.1088/0034-4885/79/12/124201 [arXiv:1504.04855 [hep-ph]].

[15] D. K. Hong, D. H. Kim and C. S. Shin, Phys. Rev. D **97** (2018) no.3, 035014 doi:10.1103/PhysRevD.97.035014 [arXiv:1706.09376 [hep-ph]].