Dark sector shining through 750 GeV dark Higgs boson at the LHC

P. Ko and Takaaki Nomura
School of Physics, KIAS, Seoul 02455, Korea
(Dated: May 23, 2016)

We consider a dark sector with $SU(3)_C \times U(1)_Y \times U(1)_X$ and three families of dark fermions that are chiral under dark $U(1)_X$ gauge symmetry, whereas scalar dark matter $X$ is the SM singlet. $U(1)_X$ dark symmetry is spontaneously broken by nonzero VEV of dark Higgs field ($\Phi$), generating the masses of dark fermions and dark photon $Z'$. The resulting dark Higgs boson $\phi$ can be produced at the LHC by dark quark loop (involving 3 generations) and will decay into a pair of photon through charged dark fermion loop. Its decay width can be easily $\sim 45$ GeV due to its possible decays into a pair of dark photon, which is not strongly constrained by the current LHC searches $pp \rightarrow \phi \rightarrow Z'Z'$ followed by $Z'$ decays into the SM fermion pairs. The scalar DM can achieve thermal relic density without conflict with direct detection bound or the invisible $\phi$ decay into a pair of DM.

INTRODUCTION

Recently both ATLAS and CMS Collaborations announced that there are some excess in the diphoton channel around $m_{\gamma\gamma} \approx 750$ GeV \cite{1,2}:\begin{equation}
\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) = (6.2^{+2.4}_{-2.0}) \text{ fb} \quad \text{(ATLAS)} \tag{1}
\end{equation}
\begin{equation}
\Gamma_{\text{tot}}(\phi) \sim 45\text{GeV}(\text{ATLAS}) \quad \text{(2)}
\end{equation}
whereas the CMS data prefers a smaller decay width \cite{3}. Furthermore, at Moriond 2016, ATLAS and CMS have reported that the local (global) significances of the diphoton excess are about 3.9(2.0)$\sigma$ and 3.4(1.6)$\sigma$, respectively, where CMS added 0.6 fb$^{-1}$ new data to the 13 TeV analysis and combined with 8 TeV data \cite{3,4}.

This excess motivated a lot of phenomenological study on possible scenarios of new physics beyond the Standard Model (BSM) which include models related to DM physics \cite{5–44}, new gauge symmetry models \cite{13,20,23} and other models \cite{58–161}. It is not easy to generate a large enough width $\sim 45$ GeV with large $BR(\phi \rightarrow \gamma\gamma)$, maintaining relevant cross section of $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)$ $\sim O(10)$ fb and evading various collider search bounds.

In this letter, we solve these problems by introducing dark $U(1)_X$ gauge symmetry, dark photon $Z'$, three generations of dark fermions with $SU(3)_C \times U(1)_Y$ charges and singlet scalar DM $X$. Dark photon $Z'$ can decay into SM fermions via a small $\sim 45$ GeV with large $BR(\phi \rightarrow \gamma\gamma)$, maintaining relevant cross section of $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)$ $\sim O(10)$ fb and evading various collider search bounds.

Let us introduce a dark sector with new dark fermions which carry both the SM $SU(3)_C \times U(1)_Y$ quantum numbers as well as dark $U(1)_X$ gauge charges, and a SM singlet complex scalar field $X$ as summarized in Table I. In this model, every right-handed fermion $f_R$ in the SM has its partner fermion $F_L$ with nonzero dark charge in the dark sector. Then the $T^L_L f_R$ operator becomes invariant under the SM gauge transformation. Its nonzero dark charge is cancelled by the dark charge of scalar DM $X$ in such a way that $T^L_L f_R X$ becomes gauge invariant operator. And $F_L$ becomes vectorlike under the SM gauge group by introducing its chiral partner $F_R$. The model is very simple and free from gauge anomalies for arbitrary $a$ and $b$. A novel feature of this model is that the new fermions $F_L$ and $F_R$ are chiral under dark $U(1)_X$ gauge symmetry so that they are massless before spontaneous symmetry breaking. And their effects on $\phi \rightarrow gg, \gamma\gamma$ through triangle diagram evades from the decoupling theorem as their mass becomes heavy.

The Yukawa interactions and the scalar potential including new fields in the dark sector are described by

\begin{equation}
L_{\text{Yukawa}} = y^E \bar{E}_L E_R \Phi + y^N \bar{N}_L N_R \Phi^\dagger + y^U \bar{U}_L U_R \Phi^\dagger + y^D \bar{D}_L D_R \Phi + y^e \bar{E}_L e_R X + y^u \bar{U}_L u_R X^\dagger + y^d \bar{D}_L d_R X + h.c.,
\end{equation}

where $E, N, U, D, L, R, X$ are SU(3)$_C \times$SU(2)$_L \times$U(1)$_Y \times$U(1)$_X$. We consider three families of dark fermions.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|c|c|c|}
\hline
& $E_L$ & $E_R$ & $N_L$ & $N_R$ & $U_L$ & $U_R$ & $D_L$ & $D_R$ & $\Phi$ & $X$ \\
\hline
SU(3) & 1 & 1 & 1 & 1 & 3 & 3 & 3 & 3 & 1 & 1 \\
SU(2) & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 \\
U(1)$_Y$ & $-1$ & $-1$ & 0 & 0 & $\frac{2}{3}$ & $\frac{2}{3}$ & $-\frac{1}{3}$ & $-\frac{1}{3}$ & 0 & 0 \\
U(1)$_X$ & $a$ & $-b$ & $-a$ & $b$ & $a$ & $-b$ & $a$ & $b$ & $a$ & $a$ \\
\hline
\end{tabular}
\caption{Contents of new fermions and scalar fields and their charge assignments under the gauge symmetry SU(3)$_C \times$SU(2)$_L \times$U(1)$_Y \times$U(1)$_X$.}
\end{table}
\[ V = \mu^2 H^\dagger H + \lambda (H^\dagger H)^2 + \mu^0 \Phi^\dagger \Phi + \mu_X X^\dagger X \]
\[ + \lambda_X (X^\dagger X)^2 + \lambda_H (H^\dagger H)(\Phi^\dagger \Phi) \]
\[ + \lambda_H (H^\dagger H)(X^\dagger X) + \lambda_X (X^\dagger X)(\Phi^\dagger \Phi). \]

where \( H \) denote the SM Higgs field. We have suppressed the generation indices on the SM and the dark fermions for simplicity. The Yukawa interactions provide mass terms for the dark fermions \( F_i \), which decay through \( F \to X f_j \). \( X \) is the SM singlet and can be a good DM candidate. Note that there is an accidental \( Z_2 \) symmetry, \( X \to -X, F_L \to -F_L \) and \( F_R \to -F_R \) which make \( X \) stable at renormalizable level. There could be gauge invariant operators that break this accidental \( Z_2 \) symmetry: \( X^\dagger \Phi^a \) and/or \( X \Phi^a \) which would generate nonzero VEV for \( X \) after \( U(1)_X \) symmetry breaking by nonzero \( \langle \Phi \rangle \neq 0 \). Gauge invariance requires that \( \pm a/(a+b) = n \) to be an integer. We can forbid this type of operators by making a judicious choice of \( a, b \) so that \( \pm a/(a+b) \) is not an integer. Or we can make \( n \) very large so that even if \( X \) develops a nonzero VEV, the lifetime of \( X \) becomes long enough (\( \tau_X \gtrsim 10^{28} \) sec) to be a good DM candidate. This model can be considered as a generalization of the singlet portal extensions of the SM where dark matter lives in the dark sector \([166]\), but the dark sector now contains dark fields which are charged under the SM gauge group as well as dark gauge group, unlike the earlier models \([166]\).

The gauge symmetry is broken after \( H \) and \( \Phi \) get non-zero VEVs:

\[ H = \left( \begin{array}{c} G^+ \\ \frac{1}{\sqrt{2}}(v + h + iG^0) \end{array} \right), \quad \Phi = \frac{1}{\sqrt{2}}(v_\phi + \phi + iG_\phi), \]

where \( G^\pm, G^0 \) and \( G_\phi \) are NG bosons which are absorbed by \( W^\pm, Z \) and \( Z' \) respectively. We shall call \( \phi \) as dark Higgs boson, since it appears as a result of spontaneous breaking of dark \( U(1)_X \) gauge symmetry.

We assume \( \lambda_{H\Phi} \) is negligible and the mixing between SM Higgs boson \( h \) and \( \phi \) is negligibly small which is consistent with the current Higgs data analysis \([167]\). Then the scalar VEVs are given approximately by

\[ v \approx \sqrt{\frac{-\mu^2}{\lambda}}, \quad v_\phi \approx \sqrt{\frac{-\mu^0}{\lambda_\phi}}. \]

The masses of new fermions are generated such that

\[ M_F = \frac{g_F}{\sqrt{2}} v_\phi, \]

where \( F = E, N, U \) and \( D \).

We consider kinetic mixing of the \( U(1)_Y \) and \( U(1)_X \) gauge fields which are denoted respectively as \( B_\mu \) and \( \tilde{X}_\mu \):

\[ \mathcal{L}_{\text{kin}} = -\frac{1}{4} W_{\mu \nu}^a W^{a \mu \nu} - \frac{1}{4} (\tilde{B}_{\mu \nu}, \tilde{X}_{\mu \nu}) \left( \begin{array}{c} 1 \\ s_\chi \end{array} \right) \left( \begin{array}{c} 1 \\ s_\chi \end{array} \right) \tilde{Z}^{\mu \nu} \]

where \( s_\chi \equiv \sin \chi \). The kinetic terms are diagonalized by the following non-unitary transformation:

\[ \left( \begin{array}{c} \tilde{B}_\mu \\ \tilde{X}_\mu \end{array} \right) = \left( \begin{array}{cc} 1 & -t_\chi \\ 0 & 1/t_\chi \end{array} \right) \left( \begin{array}{c} B_\mu \\ X_\mu \end{array} \right), \]

where \( t_\chi = \tan \chi \). After \( \Phi \) and \( H \) develop non-zero VEVs, the mass matrix for neutral gauge field is approximately given by

\[ \frac{1}{8} \left( \begin{array}{c} \tilde{Z} \\ X \end{array} \right)^T \left( \begin{array}{cc} (g^2 + g'^2)v^2 & t_\chi g' \sqrt{g^2 + g'^2}v^2 \\ t_\chi g' \sqrt{g^2 + g'^2}v^2 & 4(a + b)^2 g^2 v_\phi^2 \end{array} \right) \left( \begin{array}{c} \tilde{Z} \\ X \end{array} \right), \]

where \( W^3 = \cos \theta_W Z + \sin \theta_W A \) and \( B_\mu = -\sin \theta_W + \cos \theta_W A_\mu \) are used. Assuming \( \chi \ll 1 \) \([178]\), neutral gauge boson masses are

\[ m_{Z}^2 \approx \frac{1}{4} (g^2 + g'^2)v^2, \quad m_{Z'}^2 \approx (a + b)^2 g^2 v_\phi^2. \]

The mass eigenstates are given by

\[ \left( \begin{array}{c} Z_\mu \\ Z'_\mu \end{array} \right) = \left( \begin{array}{cc} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{array} \right) \left( \begin{array}{c} \tilde{Z}_\mu \\ X_\mu \end{array} \right), \]

and the small \( Z - Z' \) mixing angle is given by

\[ \tan 2\theta \approx \frac{g' \sqrt{g^2 + g'^2}v^2}{2(m_{Z}^2 - m_{Z'}^2)} t_\chi. \]

In Fig. we show the branching ratios of \( Z' \) as a function of its mass. Here \( q = u, d, s, c, b, \) and \( \nu \bar{\nu} \) includes all the three flavors. Note that \( Z' \) decays into the SM through the kinetic mixing so that \( \Gamma(Z')/m_{Z'} \approx O(\chi^2) \lesssim 10^{-4} \). Therefore \( Z' \) would be a very narrow resonance.

We also find that Yukawa coupling of new dark fermions and \( \lambda_\phi \) can be written in terms of \( g_X \) and \( m_{Z'} \);

\[ y_F = \sqrt{2} \frac{(a + b) g_X M_F}{m_{Z'}^2}, \quad \lambda_\phi = \frac{(a + b)^2 m_{Z'}^2 g_X^2}{2m_{Z'}^2}. \]

In our analysis, we require these couplings are perturbative as \( y_F < 4\pi \) and \( \lambda_\phi < 4\pi \).

PHENOMENOLOGY

750 GeV Diphoton Excess

In this section, we analyze the production of \( \phi \) and its decays at the LHC 13 TeV. The production of \( \phi \) is
Similarly the partial decay width for $\phi \to \gamma \gamma$ is given via dark fermion loops such that

$$\Gamma_{\phi \to \gamma \gamma} = \frac{\alpha^2 m_\phi^3}{256 \pi^3} \left| \sum_F N_c^F (a + b) g_X Q_F^2 m_{Z'} \right|^2,$$

(19)

On the other hand, the decay widths of $\phi$ into $Z'Z'$, $X^*X$
is negligible. Then the dominant annihilation processes so that their contribution to thermal relic calculation of DM in our model are \(XX^*(NN) \rightarrow Z'Z'\) assuming \(m_{X,N} > m_{Z'}\). We have included the \(t(u)\) channel processes mediated by virtual \(F\) exchange as well as the \(s\)-channel process mediated by \(\phi\) exchange. Note that the \(Z'\)-exchanging processes are suppressed since interactions between \(Z'\) and SM particles are small due to the small \(Z - Z'\) mixing we assume.

The thermal relic density is numerically estimated with \texttt{micrOMEGAs 4.1.5}\footnote{\url{http://micromegas.ippcfr.fr}} to solve the Boltzmann equation by implementing relevant interactions relevant for the DM pair annihilation processes. In calculating the relic density we assume \(a \simeq b \simeq 1\) (but \(a \neq b\)). We find that the DM relic density is given dominantly by scalar DM \(X\) in the parameter region where one can explain the 750 GeV diphoton excess. It turns out that the relic density of \(N\) is small due to large Yukawa coupling \(y^N\) which makes the amplitude for the \(\bar{N}N \rightarrow \phi \rightarrow Z'Z'\) process large. Thus the thermal relic density of scalar DM \(X\) is calculated with fixed parameter set

\[\Gamma_{\phi \rightarrow Z'Z'} = \frac{(a + b)^2 g_X^2 m_{Z'}^2}{32\pi m_{\phi}} \times \frac{m_{\phi}^4 - 4m_{Z'}^2 m_{Z'}^2 + 12m_{Z'}^4}{m_{Z'}^2} \sqrt{1 - \frac{4m_{Z'}^2}{m_{\phi}^2}},\]

\[\Gamma_{\phi \rightarrow X^+X^-} = \frac{\lambda_X^2 m_{Z'}^2}{16\pi(a + b)^2 g_X^2 m_{\phi}} \sqrt{1 - \frac{4m_X^2}{m_{\phi}^2}},\]

\[\Gamma_{\phi \rightarrow FF} = \frac{g_X^2 M_F^2}{4\pi m_{Z'}^2} m_{\phi} \sqrt{1 - \frac{4m_F^2}{m_{Z'}^2}},\]

Fig. 6 shows the total decay width of \(\phi\) in the \(m_{Z'} - g_X\) plane where the same parameter set as in Fig. 2 is used. The branching fractions of \(\phi\) decay can be obtained by partial decay widths, which is shown as a function of \(g_X\) in Fig. 1 for \(m_{Z'} = 300\) GeV with the above parameter setting. Finally Fig. 5 shows contours of \(\sigma(gg \rightarrow \phi) \text{BR}(\phi \rightarrow \gamma\gamma)\) in the \(m_{Z'} - g_X\) plane. We therefore find that \(3 - 10\) fb cross section for diphoton mode can be obtained in the region of \(g_X \simeq 0.2 - 0.5\) and \(m_{Z'} < m_S/2\), simultaneously with a rather large decay width of \(\phi\): \(\Gamma_{\phi} \simeq 5 - 40\) GeV.

\[\text{FIG. 3: The total decay width of } \phi \text{ in unit of GeV with same parameter setting as Fig. 2}\]

\[\text{FIG. 4: Branching fraction for decay of } \phi.\]

\[\text{FIG. 5: The } \sigma(gg \rightarrow \phi) \text{BR}(\phi \rightarrow \gamma\gamma) \text{ in unit of fb with same parameter setting as Fig. 2}\]

**Dark Matter Phenomenology**

The DM candidates of our model are \(X\) and \(N\). We assume that the Higgs portal coupling \(\lambda_{HX} = 0\) for simplicity, since this case is studied in great detail\footnote{\url{http://micromegas.ippcfr.fr}}. We also assume that the Yukawa couplings involving the DM \(X\) and SM fermions in Eq. (4) are small enough so that their contribution to thermal relic calculation is negligible. Then the dominant annihilation processes
of \{M_{U,D}, M_{E,N}, m_X\} = \{800 \text{ GeV}, 400 \text{ GeV}, 350 \text{ GeV}\}
and by taking \{g_X, \lambda_{X\Phi}, m_{Z'}\} as free parameters. We
then search for the parameter region which give the right thermal relic density, i.e. \(\Omega h^2 = 0.1199 \pm 0.0027\) as re-
ported by Planck Collaboration \[174\]. The upper figure
in Fig. 6 shows the parameter region in the \((m_{Z'}, g_X)\)
plane providing the observed relic density for \(\lambda_{X\Phi} = 0\).
On the other hand, the lower figure in Fig. 6 shows the
corresponding parameter region in the \((m_{Z'}, \lambda_{X\Phi})\) plane
for \(g_X = 0.1\) and \(0.3\). We find that interference between \(t(u)\) channel processes and \(\phi\) exchanging \(s\) channel
process makes \(\lambda_{X\Phi}\) dependence of the relic density non-
trivial. For smaller \(\lambda_{X\Phi}\) and \(g_X\), small amount of Higgs
portal coupling \(\lambda_{HX}\) can help us to achieve the correct
thermal relic density.

In this model, DM-nucleon scattering occurs through \(h, \phi\) and \(Z'\) exchanges. The amplitude for \(Z'\) exchange
will be small since it involves \(Z - Z'\) mixing which can be
sufficiently small. Also the Higgs contribution can be made small enough if we take a small \(\lambda_{HX}\). For \(\phi\)
exchange, we have contribution to DM-nucleon scattering
amplitude from \(\phi\)-gluon-gluon coupling in Eq. \[17\] and
\(\phi - X - X\) coupling even if we suppress \(\phi - h\) mixing.
The relevant effective coupling is given by

\[
\mathcal{L}_{XXGG} = \frac{\alpha_S}{4\pi} \left( \sum_{F=U,D} \frac{\lambda_{X\Phi}}{m_{Z'}^2} A_{1/2}(\tau_F) \right) X^\dagger X G^{a\mu\nu} G^a_{\mu\nu}
\]

Then the spin-independent DM-nucleon scattering cross
section is obtained as \[173\]

\[
\sigma_{SI} = \frac{m_N^2}{\pi(m_X + m_N)^2} f_N^2
\]

\[
f_N = \frac{2}{9} C_g f_{T_G}^{(N)}
\]

where \(m_N\) is the nucleon mass and \(f_{T_G}^{(N)}\) is the mass fraction
of gluonic operators in the nucleon mass. For the
numerical values for these parameters, we adopt values
in Ref. \[176\]. We find that DM-nucleon scattering cross
section is small as \(\sigma_{SI} \lesssim 10^{-48} \text{cm}^2\) for the \(\lambda_{X\Phi}\)
providing the observed relic density in Fig. 6. Therefore it is
difficult to observe the DM-nucleon scattering in direct
detection experiment.

**Muon \((g-2)\)\(_\mu\)**

It is interesting to note that this model can also solve
the muon \((g-2)\)\(_\mu\) through the dark muon and dark mat-
ter loop. For \(m_X = 350\) GeV and \(m_{E_i} = 400\) GeV,
we can account for the deficit in the \(a_\mu = 8 \times 10^{-10}\)
if \(g^{E_i\mu} \sim 2 - 3\) assuming the universal \(g^{E\mu}_{E\mu}\) and \(m_{E_i}\).
If we assume flavor conserving Yukawa, \(y \sim 5\) is needed.

**Stability of the potential**

Here we briefly discuss the stability of the scalar potential.
The one-loop beta functions of the Yukawa coupling
\(y^F\) and \(\lambda_\Phi\) are given by \[128\]

\[
\beta_{y^F} = y^F \left[ 3(2N_{F_i}^c + 1)(y^F)^2 - \frac{18}{5} Q_F^2 g_{1(3)}^2 - 8g_{3}^2 \right],
\]

\[
\beta_{\lambda_\Phi} = 8\lambda_\Phi \sum_F N_{F_i}^c (y^F)^2 + 18\lambda_\Phi^2 - 8\sum_F N_{F_i}^c (y^F)^4
\]

where \(g_{1(3)}\) are gauge couplings for \(\text{SU}(1)_Y(\text{SU}(3))\) and
the \(\overline{\text{MS}}\) scheme is applied. As a rough estimation, we ig-
more the running of gauge couplings in the energy range of $O(1)$ TeV to $O(10)$ TeV since the moderate running of gauge couplings in the RHS of Eq. (28) does not make significant changes for the running behavior of $y^F$ and $\lambda_\Phi$. In Fig. 7 we show the renormalization group running of $\lambda_\Phi$ where we took $\lambda_\Phi = \{1.3, 1.4, 1.5\}$ as reference points at $\mu = 1$ TeV and assumed universal Yukawa couplings $y^F = 1.2$ at the same $\mu$ for simplicity. We thus find that $\lambda_\Phi$ cannot be too small or too large to stabilize the potential. Also relative magnitude between $y^F$ and $\lambda_\Phi$ changes the running property significantly, which can be tuned by changing $U(1)_X$ charge of $\Phi$, $a + b$, according to Eq. (15) and (16). By tuning the parameters, the stability of the potential can be achieved up to $\sim 10$ TeV. The complete analysis is beyond the scope of this letter and we left it as future work.

**Future Tests of This Model**

The model presented in this letter can be tested at the upcoming LHC experiments by searching for a pair of dark photons around $m_{Z'} \sim 750$ GeV in the following channels:

$$pp \to \phi \to Z' Z', \quad Z' Z' \to 4 j, 2 j + l l, 2 j + \bar{E}_T, 4 l, 2 l + \bar{E}_T,$$

where $\bar{E}_T$ is from $\nu \bar{\nu}$ pair. Note that the total decay width of dark photon $Z'$ should be very narrow, $\Gamma_{\text{tot}}(Z')/m_{Z'} \lesssim 10^{-4}$. If the current ATLAS result on $\Gamma_{\phi} \sim 45$ GeV is confirmed in the future, our model predicts that the main decay channel of dark Higgs $\phi$ should be a pair of dark photon, with a large cross section, $\sigma(\phi \to Z' Z') \approx O(5 - 40) \text{ pb}$ (see Fig. 1) at the LHC@$\sqrt{s} = 13$ TeV. Therefore a dedicated search for dark photon pair could confirm or exclude our model.

Our model also opens widely a new window for DM model building, especially the Higgs portal DM. By assuming that the dark sector matter fields carry nonzero SM charges, the collider signatures become richer and also the Higgs signal strength can be different from the usual Higgs portal DM models in the presence of the mixing between the dark Higgs and the SM Higgs bosons. Our model can satisfy all the constraints from (in)direct search bounds as well as DM searches at colliders.

**CONCLUSION**

In this letter, we proposed a new dark matter model with 3 generations of dark fermions that are chiral under new dark $U(1)_X$ gauge symmetry. Both dark photon and the dark fermions get their masses entirely from spontaneous breaking of $U(1)_X$ gauge symmetry from the nonzero VEV of $\Phi$, and dark Higgs boson $\phi$ appears as a result. Then the diphoton excess at $750$ GeV is identified as the dark Higgs boson from $U(1)_X$ symmetry breaking. The main decay mode of $\phi$ is a pair of dark photon ($\phi \to Z' Z'$) and could be probed at the LHC by searching for $4 j, 2 j + l l, 2 j + \bar{E}_T, 4 l, 2 l + \bar{E}_T$. It is remained to be seen if the $750$ GeV diphoton excess survives in the future data accumulation. If it does, the model presented in this letter would be an interesting possibility without conflict with the known experimental constraints even for large decay width of $\phi$. In particular the production and the decay of the dark Higgs boson $\phi$ involves dark fermions in the triangle loops, opening a new window to the dark sector.

We are grateful to Jack Kai-Feng Chen, Sung Won Lee and Hwidong Yoo for discussions on the experimental status on the dark photon searches. We also thank the anonymous referee for valuable suggestions. This work is supported in part by National Research Foundation of Korea (NRF) Research Grant NRF-2015R1A2A1A05001869, and by SRC program of NRF Grant No. 20120001176 funded by MEST through Korea Neutrino Research Center at Seoul National University (PK).
1602, 186 (2016) arXiv:1512.08507 [hep-ph].

[140] N. Bizot, S. Davidson, M. Frigerio and J.-L. Kneur, JHEP 1603, 073 (2016) arXiv:1512.08508 [hep-ph].

[141] L. E. Ibanez and V. Martin-Lozano, arXiv:1512.08777 [hep-ph].

[142] S. K. Kang and J. Song, arXiv:1512.08963 [hep-ph].

[143] Y. Hamada, T. Noumi, S. Sun and G. Shiu, arXiv:1512.08984 [hep-ph].

[144] S. Kanemura, N. Machida, S. Odori and T. Shindou, arXiv:1512.09053 [hep-ph].

[145] I. Low and J. Lykken, arXiv:1512.09089 [hep-ph].

[146] L. Marzola, A. Racioppi, M. Raidal, F. R. Urban and H. Veermae, JHEP 1603, 190 (2016) arXiv:1512.09136 [hep-ph].

[147] E. Ma, arXiv:1512.09159 [hep-ph].

[148] S. Jung, J. Song and Y. W. Yoon, arXiv:1601.00006 [hep-ph].

[149] C. T. Potter, arXiv:1601.00240 [hep-ph].

[150] E. Palti, arXiv:1601.00285 [hep-ph].

[151] X. F. Han, L. Wang, L. Wu, J. M. Yang and M. Zhang, Phys. Lett. B 756, 309 (2016) arXiv:1601.00534 [hep-ph].

[152] U. Danielsson, R. Enberg, G. Ingelman and T. Mandal, arXiv:1601.00624 [hep-ph].

[153] W. Chao, arXiv:1601.00633 [hep-ph].

[154] C. Csaki, J. Hubisz, S. Lombardo and J. Terning, arXiv:1601.00635 [hep-ph].

[155] A. E. C. Hernandez, I. d. M. Varzielas and E. Schumacher, arXiv:1601.00691 [hep-ph].

[156] B. Dutta, Y. Gao, T. Ghosh, I. Gogoladze, T. Li, Q. Shafi and J. W. Walker, arXiv:1601.00866 [hep-ph].

[157] H. Ito, T. Moroi and Y. Takaesu, Phys. Lett. B 756, 147 (2016) arXiv:1601.01144 [hep-ph].

[158] H. Zhang, arXiv:1601.01355 [hep-ph].

[159] I. Sahin, arXiv:1601.01676 [hep-ph].

[160] S. Fichet, G. von Gersdorff and C. Royon, arXiv:1601.01712 [hep-ph].

[161] D. Stolarski and R. Vega-Morales, Phys. Rev. D 93, no. 5, 055008 (2016) arXiv:1601.02004 [hep-ph].

[162] S. Baek, P. Ko and W. I. Park, Phys. Lett. B 747, 255 (2015) arXiv:1407.6588 [hep-ph].

[163] P. Ko and Y. Tang, JCAP 1405, 047 (2014) arXiv:1402.6449 [hep-ph].

[164] P. Ko and Y. Tang, JCAP 1501, 023 (2015) arXiv:1407.5492 [hep-ph].

[165] J. Guo, Z. Kang, P. Ko and Y. Orikasa, Phys. Rev. D 91, no. 11, 115017 (2015) arXiv:1502.00508 [hep-ph].

[166] S. Baek, P. Ko and W. I. Park, JHEP 1307, 013 (2013) arXiv:1303.4280 [hep-ph].

[167] S. Choi, S. Jung and P. Ko, JHEP 1310, 225 (2013) arXiv:1307.3948 [hep-ph]; K. Cheung, P. Ko, J. S. Lee and P. Y. Tseng, JHEP 1510, 057 (2015) arXiv:1507.06158 [hep-ph]; K. Cheung, P. Ko, J. S. Lee, J. Park and P. Y. Tseng, arXiv:1512.07853 [hep-ph].

[168] J. Jaeckel, M. Jankowiak and M. Spannowsky, Phys. Dark Univ. 2, 111 (2013) arXiv:1212.3620 [hep-ph].

[169] A. Belyaev, N. D. Christensen and A. Pukhov, Comput. Phys. Commun. 184, 1729 (2013) arXiv:1207.6082 [hep-ph].

[170] P. M. Nadolsky, H. L. Lai, Q. H. Cao, J. Huston, J. Pumplin, D. Stump, W. K. Tung and C.-P. Yuan, Phys. Rev. D 78, 013004 (2008) arXiv:0802.0007 [hep-ph].

[171] J. F. Gunion, H. E. Haber, G. L. Kane and D. Wang, Front. Phys. 80, 1 (2000).

[172] J. M. Cline, K. Kainulainen, P. Scott and C. Weniger, Phys. Rev. D 88, 055025 (2013) Phys. Rev. D 92, no. 3, 039906 (2015) arXiv:1306.3710 [hep-ph].

[173] G. Belanger, F. Bondjema, A. Pukhov and A. Semenov, arXiv:1407.0129 [hep-ph].

[174] P. A. R. Ade et al. [Planck Collaboration], Astron. Astrophys. (2014) arXiv:1303.5076 [astro-ph.CO].

[175] F. Giacchino, A. Ibarra, L. L. Honorez, M. H. G. Tytgat and S. Wild, JCAP 1602, no. 02, 002 (2016) arXiv:1507.06158 [hep-ph].

[176] J. Hisano, R. Nagai and N. Nagata, JHEP 1505, 037 (2015) arXiv:1502.02241 [hep-ph].

[177] For $a = b = 1$, there appears an extra term $\Phi^\dagger X^2$ in the potential, which breaks $U(1)_X$ down to $Z_2$ subgroup after $S$ develops nonzero VEV. Likewise, for $3a = (a + b)$, there appears an extra term $\Phi^\dagger X^3$, which breaks $U(1)_X$ down to $Z_3$ subgroup after $S$ develops nonzero VEV. In this paper, we do not consider these possibilities, relegating the readers to Ref. [162] and Refs. [163–165] for $Z_2$ and $Z_3$ cases, respectively.

[178] The upper bound on the kinetic mixing is roughly $\lesssim 0.01$ in the dark photon mass range $m_{Z'} \lesssim 350$ GeV considered in this letter [168].