Journal of the Korean Physical Society, Vol. 73, No. 4, August 2018, pp. 449∼465 Review Articles

Particle, Astroparticle Physics and Cosmology in Dark Matter Models with Dark Gauge Symmetries

Pyungwon Ko*
School of Physics, Korea Institute for Advanced Study, Seoul 02455, Korea

(Received 9 May 2018)

In this review, I describe a class of dark matter (DM) models where DM is either stable due to unbroken dark gauge symmetries, or long-lived due to accidental global symmetries resulting from the underlying dark gauge symmetries. Within these models, I discuss various topics on DM in the context of particle and astroparticle physics and cosmology: DM thermal relic density and (in)direct detection, Higgs inflation assisted with Higgs portal interaction, dark radiation (DR), DM-DR interaction and suppression of the matter power spectrum at large k, strong first order phase transition and gravitational wave production, etc.. Especially I emphasize the importance of dark gauge symmetries, unitarity and renormalizability, and the limitation of the DM effective field theory (EFT) or simplified models for DM searches at high energy colliders, including the role of dark Higgs boson and dark gauge bosons.

PACS numbers: 12.60.−i, 12.90.+b, 13.90.+i, 95.35.+d
Keywords: Dark matter, Dark gauge symmetry, Dark Higgs, Dark gauge boson
DOI: 10.3938/jkps.73.449

I. INTRODUCTION

The long-sought-for SM Higgs boson has been finally discovered at the Large Hadron Collider (LHC) in 2012 [1,2]. So far the standard model (SM) has been extremely successful in explaining various experimental data from low energy atomic scale up to $\sim O(1)$ TeV scale except for a few places. The only unexplored territories are the Higgs self couplings, the Yukawa couplings of the SM fermions (especially the 1st and the 2nd generation fermions), and nonperturbative aspects of the SM such as QCD instanton and electroweak sphaleron. Future colliders such as ILC, CEPC, FCC-hh, FCC-ee, FCC-he etc. will probe these territories to some extent.

As of writing this paper, there are a few anomalies which are not fully accounted for within the SM. Let me just list them, relegating the details to the recent literature:

- Muon $g-2$: see Ref. [3] and references therein
- B physics anomalies: see, for example, Refs. [4–8]
- Proton radius puzzle: see Refs. [9–14]

It is amusing to notice that most of them except for $R(D^{(*)})$ anomalies involves muons, but there are no good theoretical resolutions of all them at once. It remains to be seen whether some of these anomalies will be gone or not, when more data are accumulated and systematic uncertainties become under better control.

Even if the aforementioned puzzles turn out to be due to statistical or systematic uncertainties and they are not signals of new physics beyond the SM (BSM), there are observational facts which definitely call for new physics beyond the SM (BSM):

- Baryon number asymmetry of the universe (BAU)
- Nonzero neutrino masses and mixings
- Nonbaryonic dark matter (DM)
- Inflation in the early universe
- Dark energy.

There are huge literature for each issue in many different directions. It would be interesting and important if some of the resolutions can ever be verified or falsified by some terrestrial experiments or astrophysical/cosmological observations within foreseeable time.

In this review article, I will concentrate on the issue of DM, mainly based on a series of my works during the past several years with a number of collaborators [15–49]. While I consider DM models with local dark gauge symmetries (the rationale for which will be discussed shortly in the following section), there appear some natural connections between DM models with local dark gauge symmetries and other issues in particle physics and cosmology, such as neutrino masses and mixings, baryo/lepto

---

*E-mail: pko@kias.re.kr; Fax: +82-2-958-3870

genesis, the strong 1st order phase transition and gravitational wave (GW) productions, and inflation, *etc.* I will try to touch upon these briefly in the following sections.

I will start with the basic assumptions for DM models I have constructed, emphasizing the role of dark gauge symmetry, renormalizability, unitarity and limitation of DM effective field theory (EFT) and simplified DM models. Then I discuss the simple Higgs portal DM models, both in EFT and in renormalizable and unitary models, and the limitations of EFT approach. Then I give specific examples where (i) DM is absolutely stable because of some unbroken dark gauge symmetry or (ii) topological reason, and (iii) DM is long-lived because of some accidental global symmetry that is the result of the underlying dark gauge symmetry. One of the common features of these models is the existence of a new neutral scalar boson from dark sector, which I will call dark Higgs boson. I show that dark Higgs boson can play a new key role in Higgs inflation, EW vacuum stability, light mediator generating self-interaction of DM, and explaining the galactic center γ-ray excess. I will also discuss the dark radiation and its interaction with DM, and their impact on the matter power spectrum at high k region. I also show that the dark Higgs field can affect the electroweak (EW) phase transition in the early universe, and discuss the strong 1st order phase transition and gravitational wave (GW) production. Then I summarize the paper.

II. BASIC ASSUMPTIONS FOR DM MODELS

1. Relevant questions for DM

So far the existence of DM was confirmed through the astrophysical and cosmological observations where only gravitational force plays an important role. Let us first list the relevant questions we have to answer for better understanding of DM from the viewpoint of particle physics described by quantum field theory:

- How many species of DM are there in the universe ?
- What are their masses and spins ?
- Are they absolutely stable or very long-lived ?
- How do they interact among themselves and with the SM particles ?
- Where do their masses come from ?

In order to answer (some of) these questions, we have to observe DM signals through non-gravitational observations such as colliders and/or various (in)direct detection experiments.

So far, SUSY models have been the (arguably) leading candidate for BSM, because it addresses the fine tuning problem of the Higgs mass, is consistent with the idea of grand unification, and provides good CDM candidates once the R-parity is imposed as an exact symmetry. The lightest SUSY particle (LSP) such as neutralino or gravitino can make a good candidate of CDM. However, there are no hints for SUSY at the LHC so far. Therefore it would be better for us to be open-minded about the BSM, especially regarding the new physics models regarding the DM. In principle physics of DM does not have to be connected with the fine tuning problem of Higgs mass parameter.

From particle physics point of view, the most unique and important property of DM would be that DM particle should be absolutely stable or long-lived enough, similarly to the case of electron and proton in the SM. Let us recall that electron stability within the SM is accounted for by electric charge conservation (which is an exact symmetry), and this implies that there should be massless photon, the gauge boson of unbroken $U(1)_{\rm em}$ gauge symmetry. On the other hand, the longevity of proton is ascribed to the baryon number that is an accidental global symmetry of the SM, and is broken only by dim-6 operators. Note that this is also related with proton being a composite of 3 valence quarks, and not a fundamental fermion.

We would like to have DM models where DM is absolutely stable or long-lived enough in the similar way that electron is stable and proton is long-lived in the SM. And this special property of DM has to be realized in the fundamental Lagrangian for DM in a proper way in QFT, similarly to QED and the SM. In this regard, local dark gauge symmetry will play important roles, by guaranteeing the stability/longevity of DM, as well as determine dynamics of DM and SM particles in a complete and mathematically consistent manner.

2. Hidden sector DM and local dark gauge symmetry

If one introduces new particles with nonzero SM charges and weak scale masses, there are very strong constraints from electroweak precision test and CKM phenomenology. The simplest way to evade these two strong constraints is to assume a weak scale hidden sector, which consists of particles without SM gauge charges. Note that hidden sector particles could be a good cold DM (CDM) candidates of the universe, if they are absolutely stable or long lived. Note that hidden sectors are very generic in many BSMs, including SUSY models and superstring theories.

The hidden sector matters may have their own gauge interactions (which we call dark gauge interaction) associated with local dark gauge symmetry $G_{\rm hidden}$. This local dark gauge symmetry $G_{\rm hidden}$ in the hidden sector is to stabilize the weak scale DM particle by dark charge conservation laws, if it is unbroken or it has unbroken subgroup, in the same way electron is absolutely stable in the SM because it is the lightest charged particle and electric charge is absolutely conserved. Dark gauge symmetry could be helpful for longevity of DM too if there are accidental global symmetries that prohibits DM decaying operators up to dimension 5 opearators. Hidden sector DM can be easily thermalized through some messengers connecting the SM and the hidden sectors. We shall assume all the singlet operators such as Higgs portal, right-handed neutrinos (if it is a gauge singlet) and $U(1)$ gauge kinetic mixing play the role of messengers.

Finally note that all the particles observed so far in Nature feel some gauge interactions in addition to gravity. Therefore it sounds very natural to assume that dark matter of the universe (at least some of the DM species) may also feels some (new) gauge force, in addition to gravity.

III. EFT VS. RENORMALIZABLE THEORIES: HIGGS PORTAL DM MODELS AS EXAMPLES

1. Higgs portal DM models

In order to demonstrate the limitation of the EFT or simplified models for DM phenomenology, let us start with the Higgs portal DM models, which is the simplest DM models in terms of the number of new degrees of freedom. In the literature, three types of Higgs portal DM models have been studied comprehensively [50–53]:

$$\mathcal{L}_{\rm scalar} = \frac{1}{2}\partial_\mu S\partial^\mu S - \frac{1}{2}m_S^2 S^2 - \frac{1}{2}\lambda_{HS}S^2 H^\dagger H, \quad (1)$$

$$\mathcal{L}_{\rm fermion} = \overline{\psi}\left[i\not{\partial} - m_\psi\right]\psi - \frac{\lambda_{H\psi}}{\Lambda}H^\dagger H\overline{\psi}\psi, \quad (2)$$

$$\mathcal{L}_{\rm VMD} = -\frac{1}{4}V_{\mu\nu}V^{\mu\nu} + \frac{1}{2}m_V^2 V_\mu V^\mu - \frac{\lambda_{HV}}{2}V_\mu V^\mu |H|^2 - \frac{\lambda_V}{4!}V^4, \quad (3)$$

where one imposes dark Z_2 symmetries in order to stabilize the DM particles. Under dark Z_2, the DM fields are assumed to transform as

$$S \to -S, \quad \psi \to -\psi, \quad V_\mu \to -V_\mu .$$

These models are very convenient to analyze, since DM phenomenology depends only on two extra parameters, the DM mass and the Higgs portal coupling. However one should be careful with these models.

The scalar DM Lagrangian is renormalizable, and has no problem in principle. On the other hand, the Lagrangians for fermion and vector DM are not renormalizable or unitary. Effective field theory (EFT) approaches such as Eqs. (2) and (3) are often adopted for DM physics, which however could lead to unphysical results, especially at high energy colliders. Therefore it is safer to consider their UV completions, where unitarity and renormalizability are respected. Let us discuss these two cases one by one.

2. Fermionic DM with Higgs portal

In order to illustrate the main point clearly, let us start with a singlet fermion DM model with Higg portal in EFT, Eq. (2). This simple model is nice for phenomenology, since one can study DM physics with just two new parameters, $\lambda_{H\psi}/\Lambda$ and m_ψ. This is why this model has been widely discussed in literature. However this model has to be improved since it is not renormalizable and thus violates unitarity at high energy scale.

Let us consider one of its UV completions [18, 19], which is probably the simplest UV completion in terms of the number of new degrees of freedom:

$$\mathcal{L}_{\rm DM} = \frac{1}{2}(\partial_\mu S\partial^\mu S - m_S^2 S^2) - \mu_S^3 S - \frac{\mu_S'}{3}S^3 - \frac{\lambda_S}{4}S^4 + \overline{\psi}(i\not{\partial} - m_\psi)\psi - \lambda S\overline{\psi}\psi - \mu_{HS}SH^\dagger H - \frac{\lambda_{HS}}{2}S^2 H^\dagger H. \quad (4)$$

We have introduced a singlet scalar S in order to make the model (2) renormalizable, and call this model the singlet fermion DM model (SFDM). Then there will be two neutral scalar bosons H_1 and H_2 (two independent linear combinations of H and S) in our model. The additional scalar S makes the DM phenomenology completely different from those from Eq. (2). This is also true for vector DM models [20,26].

For example, the direct detection experiments such as XENON100 and LUX exclude thermal DM within the EFT model (2), but this is not true within the UV completion (4), because of generic cancellation mechanism in the direct detection due to destructive interference between H_1 and H_2 contributions for fermion or vector DM [18,20]. The direct detection cross section in the UV completion is related with that in the EFT by [28]

$$\sigma_{\rm SI}^{\rm ren} = \sigma_{\rm SI}^{\rm EFT}\left(1 - \frac{m_{125}^2}{m_2^2}\right)^2 \cos^4\alpha \ . \quad (5)$$

Here m_2 is the mass of the singlet-like scalar boson and m_{125} is the Higgs mass found at the LHC. Note that the EFT result is recovered when $\alpha \to 0$ and $m_2 \to \infty$. This expression (5) includes the cancellation mechanism in the DM direct detection, and corrects the results reported by ATLAS and CMS (see Fig. 1). And it turns out that the same cancellation mechanism works for unitary and gauge invariant model for vector DM with Higgs portal [26].

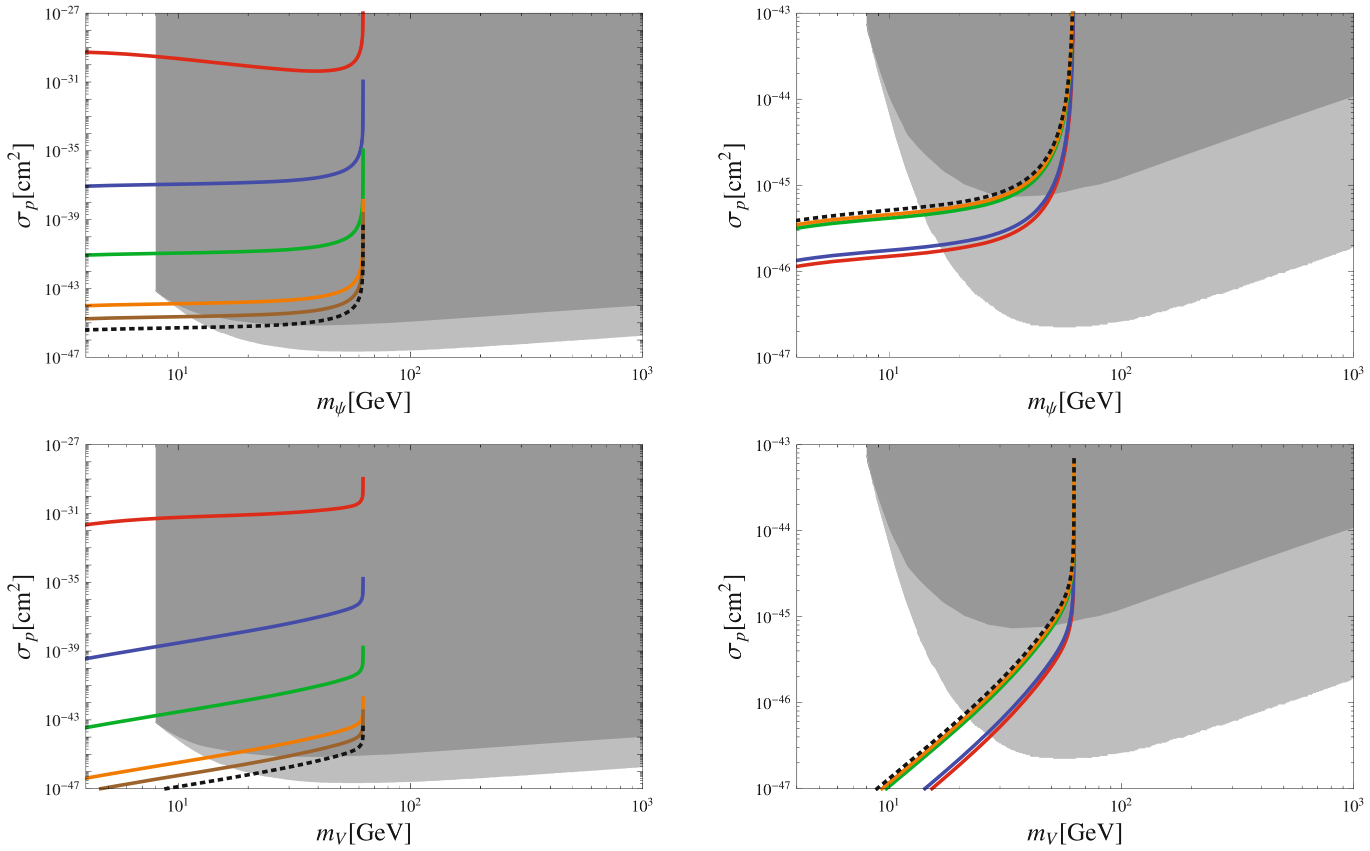

Fig. 1. (Color online) σ_p^{SI} as a function of the mass of dark matter for SFDM (top) and VDM (bottom) for a mixing angle $\alpha = 0.2$. Left panel: $m_2 = 10^{-2}, 1, 10, 50, 70$ GeV for solid lines from top to bottom. Right panel: $m_2 = 100, 200, 500, 1000$ GeV for dashed lines from bottom to top. The black dotted line is EFT predictions presented by ATLAS and CMS [54,55]. Dark-gray and gray region are the exclusion regions of LUX and projected XENON1T (gray).

3. Dark Higgs mechanism for the vector DM (VDM) and galactic center (GC) γ-ray excess

One can also consider Higgs portal vector dark matter (VDM) both in EFT and in a unitary and renormalizable model [20], where dark Higgs is naturally introduced. The Higgs portal VDM model within the EFT is usually described by Eq. (3). Although all the operators are either dim-2 or dim-4, this Lagrangian breaks gauge invariance and is neither unitary nor renormalizable when we include the Higgs portal interaction. One has to consider its UV completion.

We can consider the renormalizable version of the Higgs portal VDM by introducing a dark Higgs Φ that generates nonzero mass for VDM by the usual Higgs mechanism [20]:

$$\mathcal{L}_{VDM} = -\frac{1}{4}X_{\mu\nu}X^{\mu\nu} + (D_\mu\Phi)^\dagger(D^\mu\Phi) - \lambda_\Phi\left(|\Phi|^2 - \frac{v_\Phi^2}{2}\right)^2 - \lambda_{\Phi H}\left(|\Phi|^2 - \frac{v_\Phi^2}{2}\right)\left(|H|^2 - \frac{v_H^2}{2}\right) , \quad (6)$$

Then the dark Higgs from Φ mixes with the SM Higgs boson in a similar manner as in SFDM. And there should be a generic cancellation again in the DM direct detection cross section. Therefore one can have a wider range of VDM mass compatible with both thermal relic density and direct detection cross section (see Ref. [20] for more details). In particular the dark Higgs can play an important and crucial role in DM phenomenology (see below the discussion about the GeV scale γ-ray excess from the GC).

An important observable in Higgs portal DM models is the Higgs invisible decay width. The invisible Higgs decay width in the EFT VDM model is given by

$$(\Gamma_h^{\mathrm{inv}})_{\mathrm{EFT}} = \frac{\lambda_{VH}^2}{128\pi}\frac{v_H^2 m_h^3}{m_V^4} \times \left(1 - \frac{4m_V^2}{m_h^2} + 12\frac{m_V^4}{m_h^4}\right)\left(1 - \frac{4m_V^2}{m_h^2}\right)^{1/2} . \quad (7)$$

Note that the invisible decay rate in the EFT becomes arbitrarily large as $m_V \to 0$, which is not physically sensible. Let us compare this with the invisible Higgs decay rate in the renormalizable and unitary Higgs portal VDM

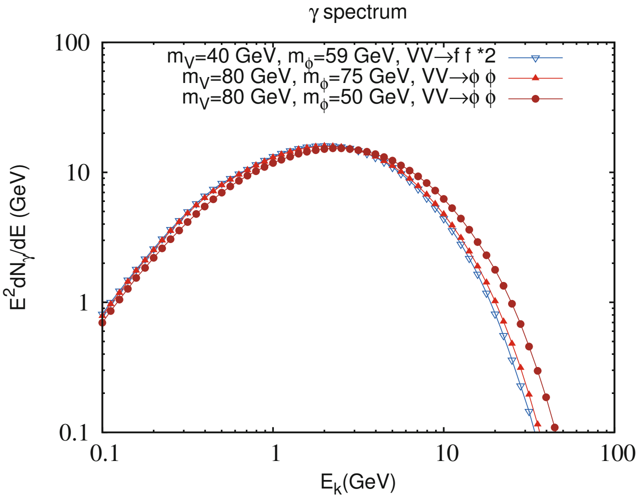

Fig. 2. (Color online) Illustration of the γ-ray spectra from different channels. The first two cases give almost the same spectra while in the third case γ is boosted so the spectrum is shifted to higher energy.

model, which is given by [28]

$$\Gamma_i^{\rm inv} = \frac{g_X^2}{32\pi}\frac{\kappa_i m_i^3}{m_V^2}\left(1-\frac{4m_V^2}{m_i^2}+12\frac{m_V^4}{m_i^4}\right)\left(1-\frac{4m_V^2}{m_i^2}\right)^{1/2}, \tag{8}$$

where m_V is the mass of VDM and $\kappa_1 = \cos^2\alpha$ and $\kappa_2 = \sin^2\alpha$ (we assume H_2 is the observed 125 GeV scalar boson). In this case $m_V = g_X v_\Phi$ so that the invisible decay width does not blow up when $m_V \to 0$, unlike the EFT VDM case. This is another example demonstrating the limitation of the EFT calculation.

Having the dark Higgs boson can be very important in DM phenomenology. Let me demonstrate it in the context of the GeV scale γ-ray excess from the GC. In the Higgs portal VDM with dark Higgs boson, one can have a new channel for γ-rays: namely, $VV \to H_2 H_2$ followed by $H_2 \to b\bar{b}, \tau\bar{\tau}$ through a small mixing between the SM Higgs and the dark Higgs. As long as V is slightly heavier than H_2 with $m_V \sim 80$ GeV, one can reproduce the γ-ray spectrum similar to the one obtained from $VV \to b\bar{b}$ with $m_V \sim 40$ GeV (see Fig. 2 and Ref. [26] for more detail). Note that this mass range for VDM was not allowed within the EFT approach based on Eq. (4), where the dark Higgs boson was not present from the beginning. It would have been simply impossible to accommodate the GC γ-ray excess within the Higgs portal VDM within EFT, simply because there is no dark Higgs boson in the EFT. Note that this mechanism would be generically possible in hidden sector DM models [29].

4. Collider Search for DM : Beyond the DM EFT and simplified models

Finally let us discuss the collider search for the dark Higgs boson and DM particles. A classic signature for DM search would be mono $X + \not{E}_T$. In early 2015, both ATLAS and CMS reported such studies in the monojet $+\ \not{E}_T$ [56] and the $t\bar{t}\ +\ \not{E}_T$ [57], respectively. Their analyses were based on the simplified model which was neither renormalizable nor unitary.

Let us consider the following example:

$$\mathcal{L}_{SS} \equiv \frac{1}{\Lambda_{dd}^2}\bar{q}q\bar{\chi}\chi \quad \text{or} \quad \frac{m_q}{\Lambda_{dd}^3}\bar{q}q\bar{\chi}\chi. \tag{9}$$

Here χ is a Dirac fermion DM that is stabilized by some conserved quantum number. A lot of results have been obtained on the scale Λ_{dd} of this operator in literature, assuming the complementarity among direct detection, collider search and indirect detection (or thermal relic density) [34].

However, the above operator is not suitable for DM search study at high energy colliders since it is not invariant under the full SM gauge symmetry. Therefore this operator has to be mended in order that the full SM gauge symmetry could be respected. Note that the operator $\bar{q}q$ can be written into $\overline{Q}_L H d_R$ and $\overline{Q}_L \widetilde{H} u_R$ for down-type and up-type quarks (nothing but the Yukawa couplings) resepctively, in a way invariant under the full SM gauge symmetry. Here $Q_L \equiv (u_L, d_L)^T$. Likewise, the singlet fermion χ cannot have renormalizable couplings to the SM Higgs boson ($h\bar{\chi}\chi$ where h is the Higgs field after electroweak symmetry breaking), since χ is a singlet whereas the Higgs field comes from a doublet. Similarly, the quark bilinear $\bar{q}q$ can not have renormalizable couplings to a singlet scalar field S.

All these problems can be resolved if we introduce a real singlet scalar field S and write down a renormalizable operator that is invariant under the full SM gauge group [18,20]. The SM Higgs will mix with the S Higgs fields after EWSB. Then one can generate Eq. (9) by $s\bar{\chi}\chi \times h\bar{q}q \to \frac{1}{m_s^2}\bar{\chi}\chi\bar{q}q$ through the $h - s$ mixing, which results in two physical neutral scalars H_1 and H_2 with the mixing angle α. Exchange of these two H_1 and H_2 for DM direct detection scattering result in a generic cancellation between two contributions from two neutral scalars, which cannot be seen within EFT approach [18, 20].

Such a model for a singlet fermion DM χ and a singlet scalar S was already discussed in Sec. III.2, Eq. (4), with identifying ψ with χ. And one can calculate the $\chi q \to \chi q$ scattering amplitude therein. The interaction Lagrangian of H_1 and H_2 with the SM fields and DM χ

is given by

$$\mathcal{L}_{\rm int} = -(H_1\cos\alpha + H_2\sin\alpha) \times \left[\sum_f \frac{m_f}{v_H}\bar{f}f - \frac{2m_W^2}{v_H}W^+_\mu W^{-\mu} - \frac{m_Z^2}{v_H}Z_\mu Z^\mu\right] + \lambda(H_1\sin\alpha - H_2\cos\alpha)\bar\chi\chi\ , \quad (10)$$

The observed 125 GeV scalar boson is denoted by H_1. The mixing between h and s leads to the universal suppression of the Higgs signal strengths at the LHC, independent of production and decay channels [18].

The DM-quark scattering amplitude can be calculated in the renormalizable model, Eq. (4): $\chi(p) + q(k) \to \chi(p') + q(k')$, the parton level amplitude of which is given by

$$\mathcal{M} = -\overline{u(p')}u(p)\overline{u(k')}u(k)\ \frac{m_q}{v_H}\lambda\sin\alpha\cos\alpha \times \left[\frac{1}{t - m_{H_1}^2 + im_{H_1}\Gamma_{H_1}} - \frac{1}{t - m_{H_2}^2 + im_{H_2}\Gamma_{H_2}}\right] \quad (11)$$

$$\to \overline{u(p')}u(p)\overline{u(k')}u(k)\ \frac{m_q}{2v_H}\lambda\sin 2\alpha\left[\frac{1}{m_{H_1}^2} - \frac{1}{m_{H_2}^2}\right] \equiv \frac{m_q}{\Lambda_{dd}^3}\overline{u(p')}u(p)\overline{u(k')}u(k), \quad (12)$$

where $t \equiv (p' - p)^2$ is the (4-momentum transfer)^2 to the nucleon. In the second line, we assumed $t \to 0$, keeping the DM-nucleon scattering in mind. Then two scalar bosons H_1 and H_2 destructively interfere in the amplitude for the DM direct detection cross section [18] The scale of the dim-7 effective operator, $m_q\,\bar{q}q\,\bar\chi\chi$ in Eq. (8), is defined in terms of Λ_{dd}:

$$\Lambda_{dd}^3 \equiv \frac{2m_{H_1}^2 v_H}{\lambda\sin 2\alpha}\left(1 - \frac{m_{H_1}^2}{m_{H_2}^2}\right)^{-1}, \quad (13)$$

$$\bar\Lambda_{dd}^3 \equiv \frac{2m_{H_1}^2 v_H}{\lambda\sin 2\alpha}, \quad (14)$$

where $\bar\Lambda_{dd}$ is derived from Λ_{dd} assuming $m_{H_2} \gg m_{H_1}$. Since the amplitude (11) was derived from renormalizable and unitary Lagrangian with the full SM gauge symmetry, it can be used for studying DM searches at high energy colliders.

The amplitude for the monojet with missing transverse energy($\not{E}_T$) signature at hadron colliders is connected to the amplitude (11) by crossing symmetry $s \leftrightarrow t$, and the effective scale $\Lambda_{\rm dd}^3$ should be replaced by

$$\frac{1}{\Lambda_{dd}^3} \to \frac{1}{\bar\Lambda_{dd}^3}\left[\frac{m_{H_1}^2}{\hat{s} - m_{H_1}^2 + im_{H_1}\Gamma_{H_1}} - \frac{m_{H_1}^2}{\hat{s} - m_{H_2}^2 + im_{H_2}\Gamma_{H_2}}\right], \quad (15)$$

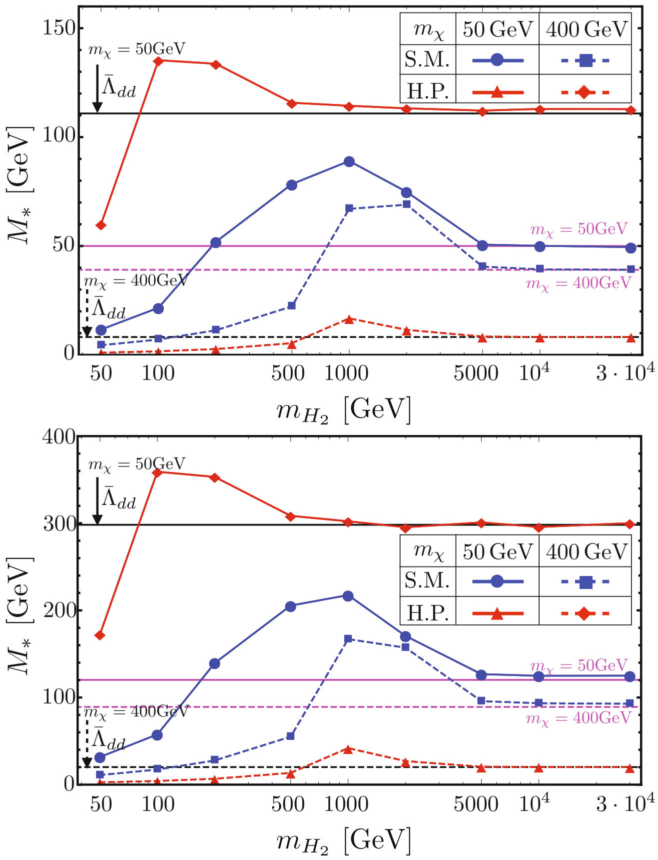

Fig. 3. (Color online) Observed exclusion limits in terms of m_χ and M_* with 90% CL. from mono-jet+$\not{E}_T$ search (left) and $t\bar{t}$+$\not{E}_T$ search (right). Blue lines are the results obtained by ATLAS [56] and CMS [57] collaborations, whereas red lines are the results obtained from renormalizable and gauge invariant Higgs portal models [34].

where $\sqrt{\hat{s}} \equiv m_{\chi\chi}$ is the DM pair invariant mass. Note that we have to include two scalar propagators, one for the 125 GeV Higgs boson and the other for the dark Higgs boson, which can not be seen in the usual DM EFT or simplified DM models. This is the result of our request for the model to be renormalizable and unitary [1]. Note that there is only a single propagator introduced to replace $1/\Lambda^2$ in the usual simplified DM models, and such presciption would break gauge invariance and unitarity in general.

The two propagators would interfere destructively for very high $\hat{s}$ or small t (direct detection), but constructively for $m_{H_1}^2 < \hat{s} < m_{H_2}^2$. If one can fix $\hat{s}$ and $m_{H_2}^2 \gg \hat{s}$, we can ignore the 2nd propagator. But at hadron colliders, $\hat{s}$ is not fixed, except for the kinematic condition $4m_\chi^2 \leq \hat{s} \leq s$ (with s = 14TeV for example at the LHC@14TeV). Therefore we cannot say clearly when we

---

[1] In fact, having two independent propagators for the mediators are very generic because the SM fermions have two different chiralities, and the SM gauge interactions are chiral. See Ref. [37] for more discussions on this point.

can ignore $\hat{s}$ compared with $m_{H_2}^2$ at hadron colliders, unless $m_{H_2}^2 > s$ (not $\hat{s}$).

One can derive the bound on the effective mass scale M_* within the full renormalizable and unitary models and compared with the bounds derived with the EFT approaches, with the same $\overline{\Lambda}_{dd}$. The results are shown in Fig. 3: the left panel on the monojet + $\not{E}_T$ from ATLAS data and the right panel on the $t\bar{t} + \not{E}_T$ from the CMS data. The blue lines are the results from the simplified model with a singlet scalar propagator, and the red lines are those from the renormalizable and unitary (and gauge invariant for the VDM) models. Note that the bounds depend very much on the underlynig model assumption, and are sensitive to the 2nd scalar boson, which does not appear in the EFT or the usual simplified model. These plots show that it is very important to analyze the monojet + $\not{E}_T$ and the $t\bar{t} + \not{E}_T$ data from the LHC within well-defined renormalizable, unitary and gauge invariant DM models. The usual EFT and the simplified models without the full SM gauge symmetry do not describe DM physics at high energy colliders properly.

Finally, the Higgs portal DM searches at the ILC and at the 100 TeV *pp* colliders have been studied within the renormalizable and unitary models in Ref. [36,44,47]. Readers are invited to the original papers on this issue.

5. Further Comments

It may sound strange that I emphasized the importance of the renormalizability of DM models, since it looks against our current understanding of effective field theory (EFT). We are working within EFT framework in most cases. This is especially fine for DM direct detections, but not for the indirect DM signatures (where the relevant energy scale is $\mu_{indirect} \sim 2m_{DM}$ for DM pair annihilation and $\mu_{indirect} \sim m_{DM}$ for DM decays) or DM collider searches (where the relevant energy scale is the collider CM energy, $\mu_{collider} \sim \sqrt{s}$). This is because we do not know the relevant degrees of freedom in the dark sectors at and below the relevant energy scales $\mu_{indirect}$ or $\mu_{collider}$. If the DM is the lightest particle in the dark sector, the story is rather simple. However there could be mediators in the dark sector which are lighter than the DM. In this case we can not integrate them out, and construct DM EFT in terms of the dark matter and the SM fields alone. The meaning of the renormalizability and unitarity is that the model should include all the relevant degrees of freedom and all the relevant and marginal (renormalizable) interactions including them as well as the SM fields. Effects of nonrenormalizable operators should be suppressed by ether $(E/\mu_{indirect})^n$ or $(E/\mu_{collider})^n$ with some positive integer n.

IV. STABLE DM WITH UNBROKEN DARK GAUGE SYMMETRIES

1. Local Z_2 scalar DM model

In order to highlight the idea of local dark gauge symmetry, let us revisit the scalar DM S with Higgs portal described by Eq. (1). This model is the simplest DM model in terms of the number of new degrees of freedom beyond the SM, and its phenomenology has been studied comprehensively (see Ref. [58] for the most recent comprehensive analysis). However the origin and the nature of Z_2 symmetry has not been specified at all in the literature.

If this Z_2 symmetry is a global symmetry, it could be broken by gravitation effects with Z_2-breaking dim-5 operator [21,30]:

$$\frac{\lambda}{M_{\rm Planck}} S F_{\mu\nu} F^{\mu\nu} \ , \quad \frac{\lambda}{M_{\rm Planck}} S \overline{Q}_L H d_R \ , \quad etc. \tag{16}$$

Then the decay rate of S due to these Z_2-breaking dim-5 operators is given by

$$\Gamma(S) \sim \frac{\lambda^2 m_S^3}{M_{Planck}^2} \sim \lambda^2 \left(\frac{m_S}{100 \text{ GeV}} \right)^3 10^{-37} \text{ GeV} \ . \tag{17}$$

Therefore the EW scale CDM 'S' would decay very fast and cannot be a good CDM candidate, unless the coefficient of this Z_2-breaking dim-5 operator λ is far less than 10^{-8}. This is one possibility, but another possibility is to implement the global Z_2 symmetry as an unbroken subgroup of some local dark gauge symmetry [21,30].

In fact, one can construct local Z_2 model á la Krauss and Wilczek [59], by assuming that a DM X and a dark Higgs ϕ_X carry $U(1)_X$-charges equal to 1 and 2, respectively. The renormalizable Lagrangian of this model is given by [30]

$$\begin{aligned} \mathcal{L} = {} & \mathcal{L}_{\rm SM} - \frac{1}{4} \hat{X}_{\mu\nu} \hat{X}^{\mu\nu} - \frac{1}{2} \sin\epsilon \hat{X}_{\mu\nu} \hat{B}^{\mu\nu} \\ & + D_\mu \phi_X D^\mu \phi_X + D_\mu X^\dagger D^\mu X - \mu \left(X^2 \phi_X^\dagger + H.c. \right) \\ & - m_X^2 |X|^2 - \lambda_X |X|^4 - \lambda_\phi \left(|\phi_X|^2 - \frac{v_\phi^2}{2} \right)^2 \\ & - \lambda_{\phi X} |X|^2 |\phi_X|^2 - \lambda_{\phi H} |\phi_X|^2 |H|^2 - \lambda_{HX} |X|^2 |H|^2, \end{aligned} \tag{18}$$

which is much more complicated than the original Z_2 scalar DM model, Eq. (1). After $U(1)_X$ symmetry is broken by the nonzero $\langle \phi_X \rangle = v_X$, there still remains a Z_2 symmetry, $X \to -X$, which guarantees the scalar DM to be absolutely stable even if we consider higher dimensional operators. The $U(1)_X$ breaking also lifts the degeneracy between the real and the imaginary parts of X, X_R and X_I, respectively. Compared with the global Z_2 scalar DM model described by Eq. (1), the local

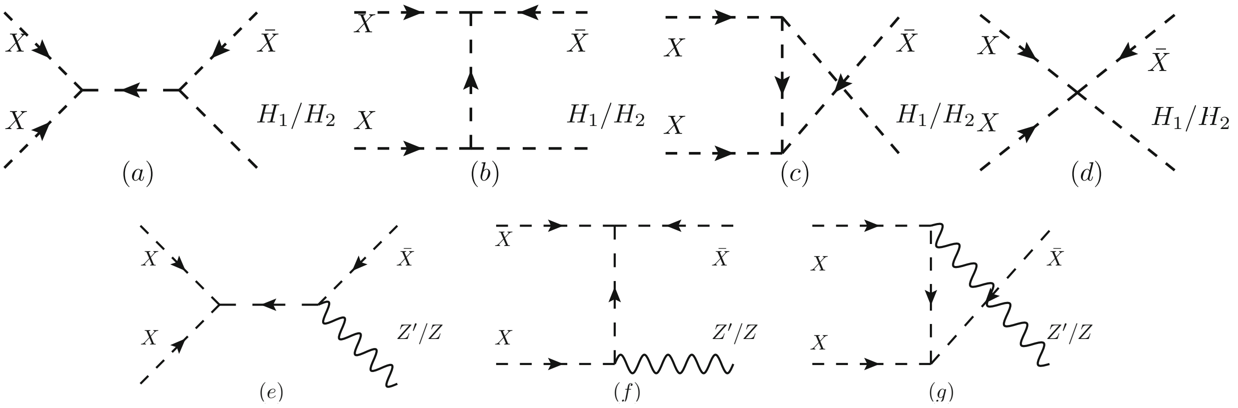

Fig. 4. Feynman diagrams for dark matter semi-annihilation which are not present in the Z_2 model discussed in IV.A. There are only (a), (b), and (c) with H_1 as final state appear in the global Z_3 model [60], whereas all diagrams could contribute in local Z_3 model [24,29].

Z_2 model has three more fields: dark photon $Z^{'}$, dark Higgs ϕ_X and the excited real scalar DM X_R, assuming X_I is lighter than X_R and makes DM. Then the DM phenomenology would be much richer than the global Z_2 scalar DM model. For example, one can consider $X_I X_I \to \phi_X \phi_X$ and the subsequent decay of ϕ_X into the SM particles through the small mixing between dark Higgs ϕ_X and the SM Higgs boson h, as a possible explanation of the galactic center γ-ray excess (see Ref. [30] for more detail).

2. Local Z_3 scalar DM model

In this subsection, we discuss another model with spontaneous $U(1)_X \to Z_3$ breaking á la Krauss and Wilczek [59] again. This can be achieved with two complex dark scalars ϕ_X and X with $U(1)_X$ charges being equal to 1 and 1/3, respectively [24, 29]. Here ϕ_X is the dark Higgs that breaks $U(1)_X$ into its Z_3 subgroup by nonzero VEV. Then the most general renormalizable Lagrangian for the SM fields and the dark sector fields, $\tilde{X}_\mu, \phi_X$ and X is given by

$$\begin{aligned}\mathcal{L} = \mathcal{L}_{\rm SM} - \frac{1}{4}\tilde{X}_{\mu\nu}\tilde{X}^{\mu\nu} - \frac{1}{2}\sin\epsilon \tilde{X}_{\mu\nu}\tilde{B}^{\mu\nu} \\ + D_\mu\phi_X^\dagger D^\mu\phi_X + D_\mu X^\dagger D^\mu X - V(H, X, \phi_X)\end{aligned} \tag{19}$$

$$\begin{aligned}V = -\mu_H^2|H|^2 + \lambda_H|H^\dagger H|^4 - \mu_\phi^2|\phi_X|^2 + \lambda_\phi|\phi_X|^4 \\ +\mu_X^2|X|^2+\lambda_X|X|^4+\lambda_{\phi H}|\phi_X|^2|H|^2+\lambda_{\phi X}|X|^2|\phi_X|^2 \\ + \lambda_{HX}|X|^2|H|^2 + \left(\lambda_3 X^3\phi_X^\dagger + H.c.\right)\end{aligned} \tag{20}$$

with $D_\mu \equiv \partial_\mu - i\tilde{g}_X Q_X \tilde{X}_\mu$.

Let us consider the phase with the following VEVs for the scalar fields in the model:

$$\langle H\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_h \end{pmatrix}, \ \langle\phi_X\rangle = \frac{v_\phi}{\sqrt{2}}, \ \langle X\rangle = 0. \tag{21}$$

This vacuum will break electroweak symmetry into $U(1)_{\rm em}$, and $U(1)_X \to Z_3$, thereby guaranteeing the stability of the scalar DM X even if we consider higher dimensional nonrenormalizable operators which are invariant under $U(1)_X$. This can be compared with the global Z_3 model in Ref. [60]. Also the particle contents and the resulting DM phenomenology in two models will be very different as summarized in Table 1.

In Fig. 4, I show the Feynman diagrams relevant for thermal relic density of local Z_3 DM X. If we worked in global Z_3 DM model instead, we would have diagrams only with H_1 in (1), (b) and (c) [60]. For local Z_3 model, there are two more new fields compared with global Z_3 model: the dark Higgs boson H_2 and the dark photon $Z^{'}$, which can make the phenomenology of local Z_3 case completely difference from that of global Z_3 case.

In fact, this can be observed immediately in Fig. 5, where the open circles are allowed points in global Z_3 model, whereas the triangles are allowed in local Z_3 case. The main difference is that in global Z_3 case, the same Higgs portal coupling λ_{HX} enters both thermal relic density and direct detections. And the stringent constraint from direct detection forbids the region for DM below 120 GeV. On the other hand this no longer true in local Z_3 case, and there are more options to satisfy all the constraints [24,29]. The color codes represent the fraction of the contribution from the semi-annihilation described in terms of the following parameter:

$$r \equiv \frac{1}{2}\frac{v\sigma^{XX\to X^*Y}}{v\sigma^{XX^*\to YY} + \frac{1}{2}v\sigma^{XX\to X^*Y}}. \tag{22}$$

We can derive the low energy EFT of this model in the limit of very heavy $Z^{'}$ and H_2, which would be nothing but the global Z_3 model plus an infinite tower of

Table 1. Comparison between the global and the local Z_3 scalar dark matter models. Here X is a complex scalar DM, H is the observed SM-HIggs like boson, and ϕ is the dark Higgs from $U(1)_X$ breaking into Z_3 subgroup.

	Global Z_3	Local Z_3
Extra fields	X	$X, Z^{'}, \phi$
Mediators	H	$H,\ Z^{'},\ \phi$
Constraints	Direct detection	Can be relaxed
	Vacuum stability	Can be relaxed
DM mass	$m_X \gtrsim 120$ GeV	$m_X < m_H$ allowed

higher dimensional operators with $U(1)_X$ gauge symmetry. However, if we started from global Z_3 model from the beginning with higher dimensional operators, the stability of DM X would not be guaranteed in general. Also one can drive the low energy EFT and discuss its limitation, the details of which can be found in Ref. [24]. The main message is that the EFT cannot enjoy the advantages of having the full particles spectra in the gauge theories, namely not-so-heavy dark Higgs and dark gauge bosons, which could be otherwise helpful for explaining the GC γ-ray excess or the self-interacting DM if either H_2 or Z' is light enough [24,29]. Therefore it is important to know which symmetry stabilizes the DM particles.

3. Other possibilities

Sterile neutrinos including the RH neutrinos are natural candidates for hidden sector fermions with dark gauge charges. In fact there have been some attempt to construct models for CDM interacting with sterile neutrinos in order to solve the some puzzles in the standard CDM paradigm as well as to reconcile the amount of dark radiation reported by Planck observation and the sterile neutrino masses and mixings that fit the neutrino oscillation data [25]. One can also consider unbroken $U(1)_X$ dark gauge symmetry with scalar DM and the RH neutrinos decay both into the SM and into the dark sector particles [21]. If $U(1)_X$ is broken, the lightest RH neutrino could be the origin of matter and DM asymmetries of the universe [21,64].

V. STABLE DM DUE TO TOPOLOGY: HIDDEN SECTOR MONOPOLE AND VECTOR DM, DARK RADIATION (DR)

In field theory there could be a topologically stable classical configurations such as domain walls and strings, *etc.* Another renowned example is the 't Hooft-Polyakov magnetic monopole [61,62]. This object in fact puts a serious problem in cosmology, and was one of the motivations for inflationary paradigm. In Ref. [23], we revived this noble idea by putting the monopole in the hidden sector and introducing the Higgs portal interaction to connect the hidden and the visible sectors.

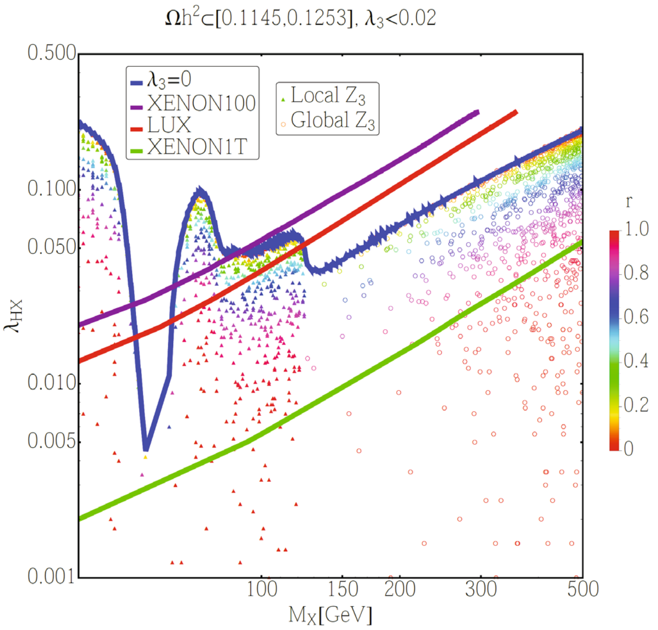
$\Omega h^2 \subset [0.1145, 0.1253]$, $\lambda_3 < 0.02$

Fig. 5. (Color online) Illustration of difference between global and local Z_3 symmetries. We have chosen M_{H_2} = 20 GeV, $M_{Z'}$ = 1TeV and $\lambda_3 < 0.02$ as an example. Colors in the scatterred triangles and circles indicate the relative contribution of semi-annihilation, r defined in Eq. (22). The curved blue band, together with the circles, gives correct relic density of X in the global Z_3 model. And the colored triangles appears only in the local Z_3 model.

We shall consider $SO(3)_X$-triplet real scalar field $\vec{\Phi}$ and add the following Lagrangian to the SM Lagrangian:

$$\mathcal{L}_{\rm new} = -\frac{1}{4} V^a_{\mu\nu} V^{a\mu\nu} + \frac{1}{2} D_\mu \vec{\Phi} \cdot D^\mu \vec{\Phi} - \frac{\lambda_\Phi}{4} \left(\vec{\Phi} \cdot \vec{\Phi} - v_\phi^2 \right)^2 - \frac{\lambda_{\Phi H}}{2} \left(\vec{\Phi} \cdot \vec{\Phi} - v_\phi^2 \right) \left(H^\dagger H - \frac{v_H^2}{2} \right). \quad (23)$$

We added $\lambda_{\Phi H}$ term describing the Higgs portal interaction, which is a new addition to the renowned 't Hooft-Polyakov monopole model.

For nonzero $\langle \vec{\Phi}(x) \rangle = (0, 0, v_\Phi)$, the original dark gauge symmetry $SO(3)_X$ is broken into its subgroup $SO(2)_X (\approx U(1)_X)$. Then, the hidden sector particles are composed of massive dark vector bosons $V_\mu^{\pm}$[2] with masses $m_V = g_X v_\Phi$, massless dark photon $\gamma_{h,\mu} \equiv V_\mu^3$, heavy (anti-)monopole with mass $m_M \sim m_V/\alpha_X$, and massive real scalar ϕ (dark Higgs boson). The massive hidden vector $V^\pm$ are stable due to the unbroken $SO(2)_X$ whereas the hidden monopole is stable due to topological reason. And the dark Higgs boson will mix with the SM

---

[2] Here ±1 in $V_\mu^\pm$ indicate the dark $U(1)_X$ charge, and not the usual electric charges.

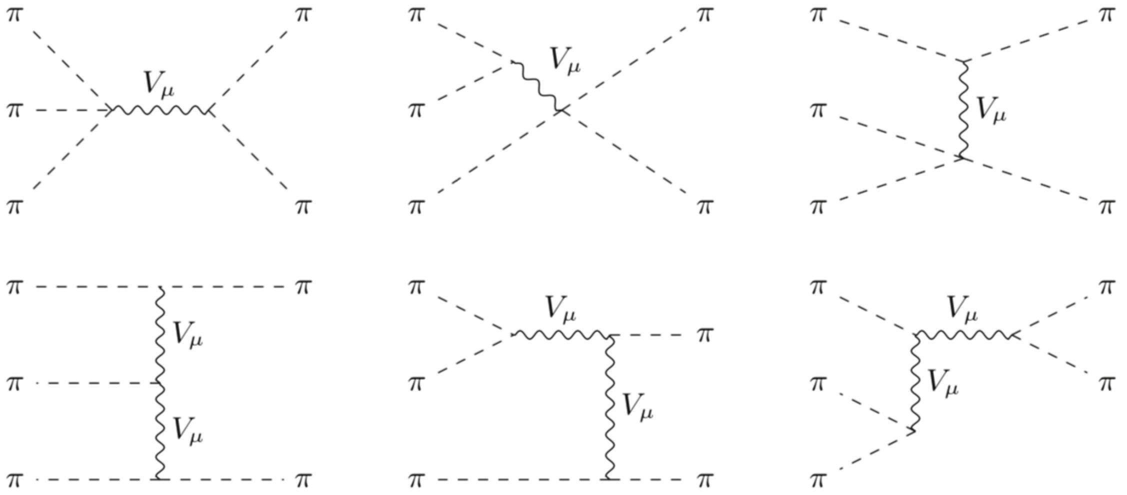

Fig. 6. Feynman diagrams contributing to 3 → 2 processes for the dark pions involving dark vector meson interactions. The 5-point contact interaction from the original WZW is not shown explicitly here.

Higgs boson through the Higgs portal term as usual.

Note that the kinetic mixing between γ_h and the SM $U(1)_Y$-gauge boson is forbidden at renormalizable level unlike the Abelian $U(1)_X$ case. This is because of the non-Abelian nature of the hidden gauge symmetry. Also the VDM is stable even if we consider higher dimensional operators because of the unbroken $SO(2)_X$. This would not have been the case, if the $SU(2)_X$ were completely broken by a complex $SU(2)_X$ doublet, where the VDM would decay in general in the presence of nonrenormalizable interactions [63]. Of course, it would be fine as long as the lifetime of the decaying VDM is long enough so that it can still be a good CDM candidate. In the VDM model with a hidden sector monopole, the unbroken $SO(2)_X$ subgroup not only guarantees the stability of VDM $V_\mu^\pm$, but also contributes to the dark radiation at the level of ~ 0.1. We refer the readers to the original paper on more details of phenomenology of this model [23].

VI. LONG-LIVED DM DUE TO ACCIDENTAL SYMMETRIES

1. EWSB and CDM from Strongly Interacting Hidden Sector

Another nicety of models with hidden sector is that one can construct a model where all the masses of the SM particles and DM are generated by dimensional transmutation in the strongly interacting hidden sector [15–17, 38]. Basically the light hadron masses such as proton or ρ meson come from confinement, which is derived from massless QCD through dimensional transmutation. One can ask if all the masses of observed particles can be generated by quantum mechanics, in a similar manner with the proton mass in the massless QCD. The most common way to address this question is to employ the Coleman-Weinberg mechanism for radiative symmetry breaking. Here I present a new model based on nonpertubative dynamics such as technicolor, chiral symmetry breaking in ordinary QCD or the Cooper pair and the energy gap in BCS superconductivity.

Let us consider a scale-invariant extension of the SM with a strongly interacting hidden sector [15–17,38]:

$$\begin{aligned}\mathcal{L} = \mathcal{L}_{\text{SM,kin}} + \mathcal{L}_{\text{SM,Yukawa}} - \frac{\lambda_H}{4}(H^\dagger H)^2 \\ - \frac{\lambda_{SH}}{2}\, S^2 H^\dagger H - \frac{\lambda_S}{4} S^4 - \frac{1}{4}\mathcal{G}^a_{\mu\nu}\mathcal{G}^{a\mu\nu} \\ + \sum_{k=1,\ldots,f} \overline{\mathcal{Q}}_k \left[iD\cdot\gamma - \lambda_k S \right] \mathcal{Q}_k. \quad (24)\end{aligned}$$

Here $\mathcal{Q}_k$ and $\mathcal{G}^a_{\mu\nu}$ are the hidden sector quarks and gluons, and the index k is the flavor index in the hidden sector QCD. We introduced a real singlet scalar S and replaces all the mass parameters by the field S in order to respect classical scale symmetry. In this model, we have assumed that the hidden sector strong interaction is vectorlike and confining like the ordinary QCD. Then we use the known aspects of QCD dynamics in order to study the hidden sector QCD.

In this model, dimensional transmutation will take place in the hidden sector and generate the hidden QCD scale and chiral symmetry breaking with nonzero $\langle \bar{\mathcal{Q}}_k \mathcal{Q}_k \rangle$. Once a nonzero $\langle \bar{\mathcal{Q}}_k \mathcal{Q}_k \rangle$ is developed, the $\lambda_k S$ term generate the linear potential for the real singlet S, which in turn results in the nonzero $\langle S \rangle$. Then the hidden sector current quark masses are induced through λ_k terms, and the EWSB can be triggered through λ_{SH} term if it has a correct sign. Then the Nambu-Goldstone boson in the hidden sector, hidden pion or dark pion π_h, will get nonzero masses, and becomes a good CDM candidate. Their dynamics at low energy can be described by chiral Lagrangian method. Thus one can calculate

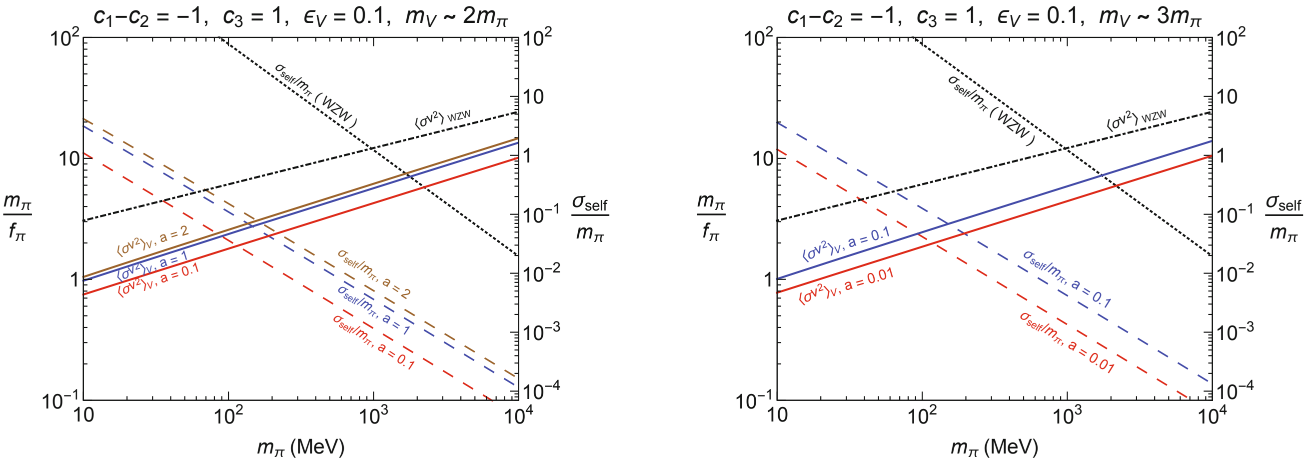

Fig. 7. (Color online) Contours of relic density ($\Omega h^2 \approx 0.119$) for m_π and m_π/f_π and self-scattering cross section per DM mass in cm^2/g as a function of m_π. The case without and with vector mesons are shown in black lines and colored lines respectively. We have imposed the relic density condition for obtaining the contours of self-scattering cross section. Vector meson masses are taken near the resonances with $m_V = 2(3)m_\pi\sqrt{1+\epsilon_V}$ on left(right) plots. In both plots, $c_1 - c_2 = -1$ and $\epsilon_V = 0.1$ are taken. See Ref. [48] for the definitions of these parameters.

thermal relic density of π_h and the DM direct detection cross section. See Ref. [15–17,38] for more details. Also hidden sector baryons $\mathcal{B}_h$ will be formed, the lightest of which would be long lived due to the accidental h-baryon number conservation. However, their dynamics is nonperturbative and I would not discuss this issue further here.

2. Strongly interacting massive particle (SIMP) scenario within the hidden QCD model

In the original models by Ko *et al.* [15–17, 38] discussed in the previous subsection, the Wess-Zumino-Witten (WZW) interaction was not considered because it is higher order in momentum expansion. If one includes the WZW term, then the DM number changing processes, 3 → 2, becomes possible and one may be able to achieve the correct relic density from this. Also 2 → 2 DM self-scattering can be large enough ($\sigma_{\rm self}/m_{\rm DM} \sim O(1)$ barn/GeV) to solve some of the vanilla ΛCDM paradigm, such as the core-cusp puzzle [65]. This new way to achieve both the relic density and the large self scattering cross section is often called Strongly Interacting Massive Particle (SIMP) scenario [66]. However, it turns out that the original proposal by Hochberg *et al.* for dark pion DM [67] is unlikely to be compatible with the validity of chiral perturbation theory, since one has to have $m_\pi/f_\pi \sim O(4\pi)$.

In Ref. [48], the present author showed that this problem can be significantly relieved if one includes the contributions from dark vector mesons (analogy of ρ and ω in the ordinary QCD) because of new 3 → 2 diagrams shown in Fig. 6 in addtion to the contact interaction from the original WZW term. Also light dark vector mesons make additional contirbutions to the dark pion self scattering through s, t and u-channel exchanges of dark vector mesons. Including these new contributions to the dark pion DM 3 → 2 and 2 → 2 scatterings from light dark vector mesons and assuming narrow width approximation for them, we find that the phenomenologically viable parameter space is about $m_\pi/f_\pi \sim$ a few (Fig. 7), which is well below 2π, the validity region of the chiral perturbation theory. It is also much smaller than the original proposal $\sim 4\pi f_\pi$ [67].

Summarizing this section, dark pion DM remains a good DM candidate, whose longevity is due to the accidental flavor symmetry of underlying dark gauge theory. Depending on the parameter space, one can achieve either WIMP or SIMP scenario. There are also dark baryons DM whose annihilation into dark pions is nonperturbative, which is not discussed here.

VII. LIGHT MEDIATORS AND SELF-INTERACTING DM

Another nice feature of the dark matter models with local dark gauge symmetry is that the model includes new degrees of freedom beyond DM particle: namely, dark gauge bosons and dark Higgs boson(s), that can play the role of force mediators from the beginning because of the rigid structure of the underlying gauge theories. In fact one can utilize the light mediators in order to explain the GeV scale γ-ray excess, or the self-interacting DM which would solve three puzzles in the CDM paradigm [65]: (i) core-cusp problem, (ii) missing satellite problem and (iii) too-big-to-fail problem. These would have been simply impossible if we adopted the EFT approach for DM physics, since there are no extra fields in the dark sector other than DM itself. (see Ref. [24] for more detail on this issue).

In the EFT approach for the DM, these new degrees of freedom are very heavy compared with the DM mass as well as the energy scale we are probing the dark sector (*e.g.*, the collider energy scale). However, we don't know anything about the mass scales of these mediators, and ignoring them as in the DM EFT approach would be too strong an assumption. Without these light mediators, we could not explain the GeV scale γ-ray excess as described in Sec. III.3, or have strong self-interacting DM. This illustrates one of the limitations of DM EFT approaches.

VIII. HIGGS INFLATION ASSISTED BY THE HIGGS PORTAL

Another interesting issue related with DM models with local dark gauge symmetris is the Higgs inflation [68,69] in the presence of the Higgs portal interaction to the dark sector [27]:

$$\mathcal{L} = -\frac{1}{2\kappa}\left(1 + \xi\frac{h^2}{M_{\rm Pl}^2}\right) R + \mathcal{L}_h + \lambda_{\phi H}\phi^2 h^2 \qquad (25)$$

in the unitary gauge, where $\kappa = 8\pi G = 1/M_{\rm Pl}^2$ with $M_{\rm Pl}$ being the reduced Planck mass, and $\mathcal{L}_h$ is the Lagrangian of the SM Higgs field only. Here ϕ denotes a generic dark Higgs field which mixes with the SM Higgs field after dark and EW gauge symmetry breaking.

In the presence of the Higgs portal interaction, we have recalculated the slow-roll parameters. Relegating the details to Ref. [27], I simply show the results: at a bench mark point for Fig. 2 of Ref. [27], we get the following results:

$$n_s = 0.9647 \ , \ r = 0.0840 \ , \qquad (26)$$

for $N_e = 56$, $h_*/M_{\rm Pl} = 0.72$, $\alpha = 0.07422199$ and $\xi = 12.8294$ for a pivot scale $k_* = 0.05{\rm Mpc}^{-1}$. There is a parameter space where the spectral running of n_s is small enough at the level of $|n_s^{'}| \lesssim 0.01$. It is amusing to notice that the r could be as large as $\sim O(0.1)$ in the presence of the Higgs portal interactions to a dark sector, in a much less sensitive way to the top quark and the Higgs boson masses in the standard Higgs inflation scenario [70,71].

IX. HIGGS PHENOMENOLOGY, EW VACUUM STABILITY AND DARK RADIATION (DR)

Now let us discuss Higgs phenomenology within this class of DM models. Due to the mixing effect between the dark Higgs and the SM Higgs bosons, the signal strengths of the observed 125 GeV Higgs boson will be universally reduced from "1" in a universal manner [18, 20]. Also the 125 GeV Higgs boson could decay into a pair of dark Higgs and/or a pair of dark gauge boson, which is still allowed by the current LHC data [22]. These predictions will be further constrained by the next round experiments.

Finally the dark Higgs can make the EW vacuum stable upto the Planck scale without any other new physics [19,20], and this was very important in the Higgs-portal assisted Higgs inflation discussed in the previous section.

In most cases, there is generically a singlet scalar which is nothing but a dark Higgs, which would give a new motivation to consider singlet extensions of the SM. Traditionally a singlet scalar was motivated mainly by why-not or $\Delta\rho$ constraint, or the strong first order EW phase transition which could be working for electroweak baryogenesis if there are new sources of CP violation. Being a singlet scalar, the dark Higgs will satisfy all these motivations, as well as stability of DM by local dark gauge symmetry. It would be important to seek for this singlet-like scalar at the LHC or the ILC, but the colliders cannot cover the entire mixing angle down to $\alpha \sim 10^{-8}$ (for MeV dark Higgs) relevant to DM phenomenology. One possible avenue would be to look for (a new resonance in) the di-Higgs channel, which would be an important topic at the LHC in the coming years.

Massless (or very light) dark gauge boson or light dark fermions in hidden sectors could also contribute to dark radiation (DR) of the universe. In a class of models we constructed, the amount of extra dark radiation is rather small by an amount consistent with the Planck data due to Higgs portal interactions [21,23,25].

X. INTERACTIONS BETWEEN DR AND DM AND SUPPRESSION OF MATTER POWER SPECTRUM

In certain models for DM and DR, there could be interaction between DR and DM through exchange of light mediator in the t-channel, which may affect the matter power spectrum at large k (or small scale) [72–74].

Here let me introduce a model based on non-Abelian dark gauge symmetry which is spontanesously borken itto its non-Abelian subgroup. For simplicity, let us consider the case $SU(3)_X \to SU(2)_X$ [41].

$$\mathcal{L} = -\frac{1}{4}F^a_{\mu\nu}F^{a\mu\nu} + (D_\mu\Phi)^\dagger\left(D^\mu\Phi\right) - \lambda_\phi\left(|\Phi|^2 - v_\phi^2/2\right)^2 , \qquad (27)$$

where $F^a_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + g f^{abc}A^b_\mu A^c_\nu$, covariant derivative D_μ is defined by $D_\mu\Phi = \left(\partial_\mu - igA^a_\mu t^a\right)\Phi$, and generators t^as are normalized as ${\rm Tr}[t^a t^b] = \delta^{ab}/2$. After $\Phi(x)$ gets a non-zero vacuum expectation value (vev), in unitary gauge we would have

$$\langle\Phi\rangle = \begin{pmatrix} 0 & 0 & \frac{v_\phi}{\sqrt{2}} \end{pmatrix}^T , \Phi = \begin{pmatrix} 0 & 0 & \frac{v_\phi + \phi\left(x\right)}{\sqrt{2}} \end{pmatrix}^T . \qquad (28)$$

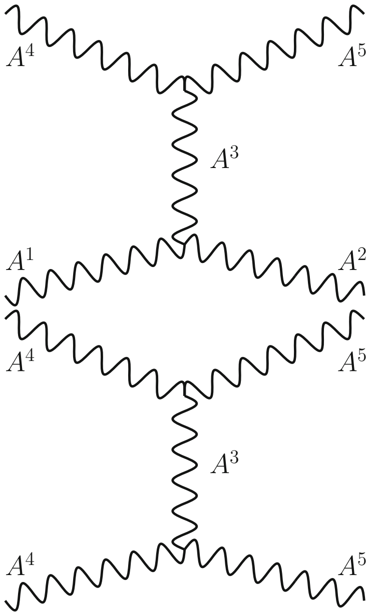

Fig. 8. Typical Feynmann diagrams for DM-DR (left) and DM-DM (right) scattering. $A^{m=4,\cdots,8}$ and $A^{a=1,2,3}$ denote the massive VDM and the massless DR, respectively.

Due to the spontaneous symmetry breaking by the above vacuum configuration, gauge bosons $A^{m=4,\cdots,8}$ obtain masses from the interaction term $g^2(A_\mu\Phi)^\dagger(A^\mu\Phi)$,

$$m_{A^{4,5,6,7}} = \frac{1}{2}gv_\phi,\ m_{A^8} = \frac{1}{\sqrt{3}}gv_\phi, \tag{29}$$

while gauge bosons $A^{a=1,2,3}$ associated with the unbroken gauge group $SU(2)$ are still massless. The stability of the massive VDM is guaranteed by the unbroken $SU(2)_X$ subgroup. Also there are three massless gauge bosons that play the role of non-Abelian DR. Interactions among DM-DM, DM-DR and DR-DR are all described by the original non-Abelian dark $SU(3)_X$ gauge theory with a single $SU(3)_X$ gauge coupling (see Fig. 8). In order to have a viable cosmological history, we need a very tiny dark gauge couplings and the Higgs portal couplings, so that the non-Abelian massive vector boson can not be the usual WIMP whose thermal density is determined by freeze-out mechanism [41]. Instead we have to call for the freeze-in mechanism to achieve the correct relic density of the massive VDM (Fig. 9).

In particular, the DM-DR scattering is generated by the t-channel exchange of DR (massless gauge boson), and its effect can modify the matter power spectrum (σ_8) in the large k region of the large scale structure. For example, we choose

$$\begin{aligned} &\Omega_b h^2 = 0.02227, && \Omega_c h^2 = 0.1184, \\ &100\theta_{\rm MC} = 1.04106, && \tau = 0.067, \\ &\ln\left(10^{10}A_s\right) = 3.064, && n_s = 0.9681, \end{aligned} \tag{30}$$

and treat neutrino mass the same way as PLANCK did with $\sum m_\nu = 0.06$ eV, which gives $\sigma_8 = 0.815$ in vanilla

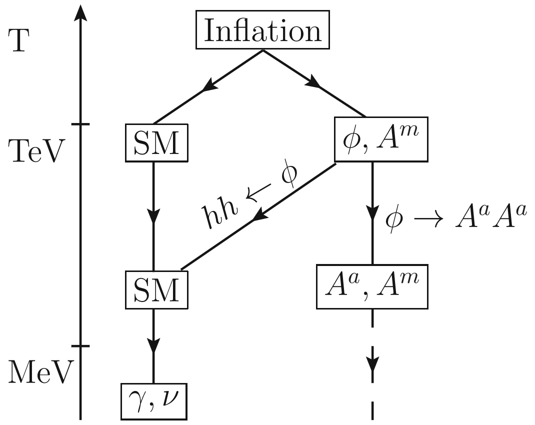

Fig. 9. Schematic picture for thermal history of DM A^m, DR A^a, dark Higgs boson ϕ and SM.

ΛCDM cosmology. Together with the same inputs as above, we take $\delta N_{\rm eff} \simeq 0.5$, $m_\chi \simeq 10$ TeV and $g_X^2 \simeq 10^{-7}$ in the interacting DM-DR case, we have $\sigma_8 \simeq 0.746$ which is much closer to the value $\sigma_8 \simeq 0.730$ given by weak lensing survey CFHTLenS (Fig. 10). I would like to emphasize that the most strongest and relevant constraint on this model comes from cosmological data, and not from collider or (in)direct DM detection experiments.

There are other models for DM-DR interactions, such as fermionic DR with light dark photon [39], or composite DM and DR from strongly interacting hidden sector [45]. The details can be found in the original papers.

XI. STRONG 1ST ORDER PHASE TRANSITION AND GRAVITATIONAL WAVE (GW) PRODUCTION

In the previous sections, I argued that the one of the generic features of the DM models with dark gauge symmetries is the existence of dark Higgs boson. I discussed their effects on the Higgs signal strength, the DM searches at colliders and Higgs inflation in Secs. III and VIII, respectively.

In this section, I discuss another role of dark Higgs in the early universe, namely the strong 1st order phase transition. Within the SM with the Higgs boson mass equal to 125 GeV, the electroweak phase transition can not be the strong 1st order. On the other hand, if there is a dark Higgs boson with Higgs portal couplings, the nature of EW phase transition can change into the strong 1st order, and/or the dark phase transition itself could be the strong 1st order one. As a result there could be gravitational wave productions from the bubble collision and turbulence, *etc.* (see Ref. [75] for example). Detection of GW produced during the strong 1st order phase transition in the early universe is going to be pursued in the future experiments such as eLISA and DECIGO.

As a specific example, let me discuss the SM extension

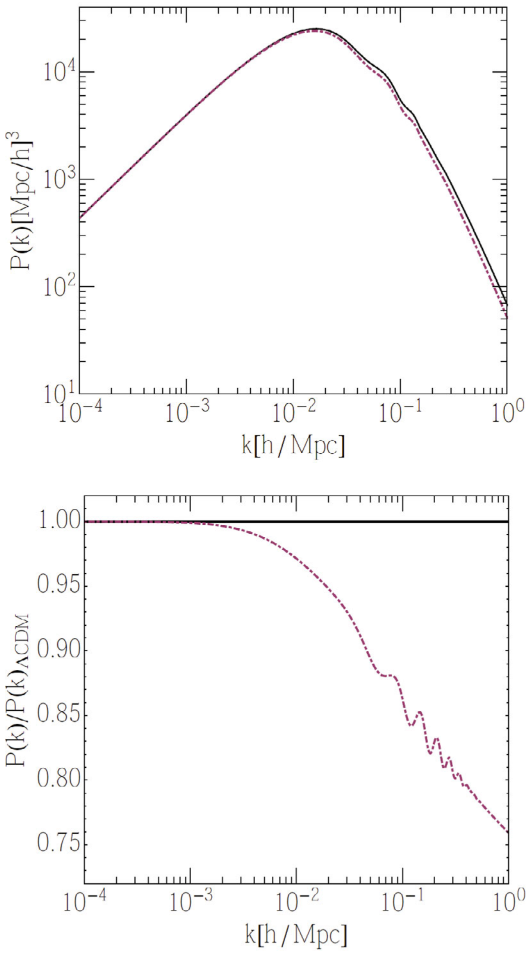

Fig. 10. (Color online) Matter power spectrum $P(k)$ (upper) and ratio (lower) with $m_\chi \simeq 10$TeV and $g_X^2 \simeq 10^{-7}$, in comparison with ΛCDM. The black solid lines are for ΛCDM and the purple dot-dashed lines for interacting DM-DR case, with input parameters in Eq. (30). We can easily see that $P(k)$ is suppressed for modes that enter horizon at radiation-dominant era. Those little wiggles are due to the well-known baryon acoustic oscillation.

with a singlet scalar [40], which would encompass the singlet fermion DM case discussed earlier in this review. For certain parameter space, one can anticipate a large deviation in the Higgs triple coupling up to $\sim 100\%$, and the strong enough GW productions observable at eLISA or DECIGO (Fig. 11).

If one considers a model with massive dark photon with dark Higgs boson which is the main theme of this review, the phase transition within this model has rich structure: one-step, two-step or three-step phase transitions, some of which could be strongly 1st order (Fig. 12). It turns out that one can observe the gravitational wave for dark photon mass heavier than $\sim$ 25 GeV, and this is complementary with the DM detection experiments in case dark photon is VDM by ignoring the kinetic mixing term [49]. This is another good example of complementarity between the GW detection and direct DM detection.

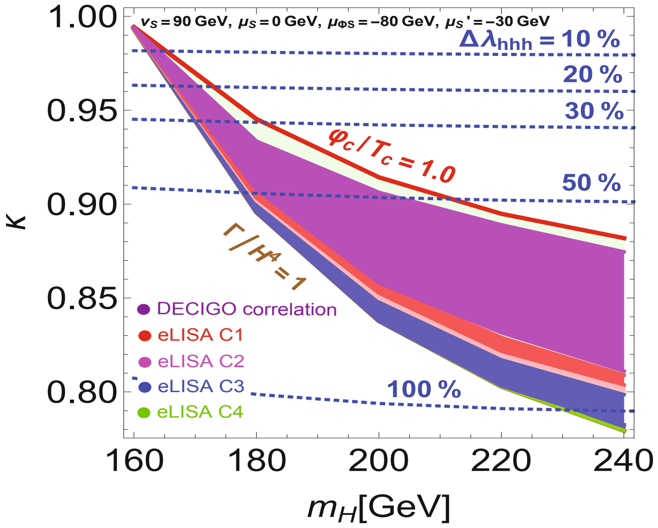

Fig. 11. (Color online) The detectability of GWs and the contours of the deviations in the hhh coupling, $\Delta\lambda_{hhh}$ in the $m_H - \kappa$ plane with $\kappa \equiv \cos\alpha$ (the mixing angle between the SM Higgs boson and the singlet scalar). The projected region of a higher sensitive detector design is overlaid with that of weaker one. The region which satisfies both $\phi_C/T_C > 1$ and $T_C > 0$ is also shown for a reference. See Ref. [40] for the numerical values of the input parameters and legends.

In another model with Z_3-symmetric scalar extension [46], it was found that this model could be probed only by the gravitational wave detection, and not by any other terrestrial experiments including colliders, unlike the SM extension with a singlet scalar. Therefore the eLISA and DECIGO can explore some types of BSM, which otherwise can never be probed (the so-called nightmare scenarios for the terrestrial particle physics experiments). This shows that the future GW detection experiments are indispensable tools for probing the many BSM models for the strong 1st order phase transitions in the early universe.

XII. CONCLUSION AND OUTLOOK

In this review article, I discussed a class of dark matter models where dark gauge symmetry plays an important role in stabilizing electroweak scalar DM or making them long lived enough compared with the age of the universe. I first discussed the limitation of the DM EFT or simplified DM models in the context of fermion and vector DM models with Higgs portal. DM EFT and naive simplified models are either nonrenormalizable or violate the full SM gauge invariance and unitarity in general, which could lead to unphysical results especially for DM phenomenology at high energy colliders. Then I discussed three explicit examples: (i) DM is stable due to unbroken

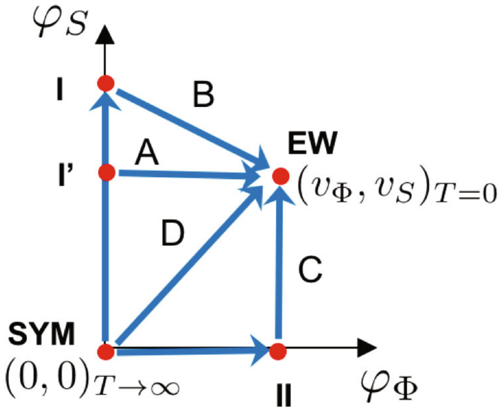

Fig. 12. (Color online) Possible phases: SYM (symmetric phase at $T = \infty$), I, I', II, EW (EW phase at $T = 0$ GeV) and transition types: A, B, C, D. Here Φ denotes the SM Higgs field.

dark gauge symmetry Z_2 and Z_3 originating from $U(1)_X$ gauge symmetry, (ii) DM is stable due to topological reason, the famous 't Hooft-Polyakov monopole in the hidden sector, and the unbroken $U(1)$ subgroup gaurantees the stability of the vector DM in the monopole sector, and (iii) DM is long lived due to global flavor symmetry which is an accidental symmetry of underlying new strong interaction in the dark sector and DM could be either WIMP or SIMP. In the models where DM is stable or long-lived due to some underlying dark gauge symmetries, there appear generically new fields, namely dark Higgs and dark gauge bosons which can play important role in DM self-interaction or galactic center γ-ray excess, which are not possible in the Higgs portal EFT for vector DM. Depending on the mass scales of dark Higgs and dark gauge bosons and their couplings to the DM and the SM fields, DM phenomenology could be much richer than the simpler DM models without them. And dark Higgs can modify the inflation or the phase transition during the evolution of the early universe. If there is interaction between DM and DR, they can affect the large scale structure of the universe in terms of H_0 and σ_8. I discussed some of those aspects for illustration.

One of the generic predictions of the DM models with local dark gauge symmetry is the existence of a new neutral scalar boson which is mostly the SM singlet if the DM particles are either fermion or vector. It affects the DM signatures at high energy colliders because of the form factors with two scalar propagators with negative sign, Eq. (15). This feature is a consequence of the full SM gauge invariance and renormalizability, and can not be seen in the usual EFT approach or simplified DM models. It is crucial to include the interference between the SM Higgs boson and the dark Higgs boson in the DM search at high energy colliders. Also the Higgs signal strength would be reduced from "one" independent of production and decay channels. It would be an important target to find them at the current/future colliders.

Before closing this review, let me mention on possible future research direction. There are still a number of issues to be further explored within this approach, such as dark matter showering at high energy colliders [77–80] and bound state formations (for example, see Refs. [81, 82] for recent discussions on this issue), to name a few. Before comparison with various observations, one has to make firm theoretical predictions.

ACKNOWLEDGMENTS

This review paper is an extension of my previous review [76], with update of my works thereafter. The author is grateful to Professor Youngjoon Kwon for giving me this opportunity to write for the 50th Anniversary of Journal of Korean Physical Society (JKPS). He also thanks Seungwon Baek, Soo-Min Choi, Katsuya Hashino, Taeil Hur, Dong Won Jung, Mitsuru Kakizaki, Zhaofeng Kang, Shinya Kanemura, Jinmian Li, Toshinori Matsui, Natsumi Nagata, Alexander Natale, Myeonghun Park, Wan Il Park, Eibun Senaha, Yong Tang, Hiroshi Yokoya and Chaehyun Yu for enjoyable collaborations on the subjects presented in this review article. This work is supported by the Korea government (MSIP) (No. 2009-0083526) through Korea Neutrino Research Center at Seoul National University.

REFERENCES

[1] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716**, 1 (2012) doi:10.1016/j.physletb.2012.08.020 [arXiv:1207.7214 [hep-ex]].

[2] S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716**, 30 (2012) doi:10.1016/j.physletb.2012.08.021 [arXiv:1207.7235 [hep-ex]].

[3] J. P. Miller, E. de Rafael, B. L. Roberts and D. Stckinger, Ann. Rev. Nucl. Part. Sci. **62**, 237 (2012) doi:10.1146/annurev-nucl-031312-120340.

[4] A. Crivellin, G. D'Ambrosio and J. Heeck, Phys. Rev. D **91**, 075006 (2015) doi:10.1103/PhysRevD.91.075006 [arXiv:1503.03477 [hep-ph]].

[5] W. Altmannshofer and D. M. Straub, arXiv:1503.06199 [hep-ph].

[6] L. Calibbi, A. Crivellin and T. Ota, Phys. Rev. Lett. **115**, 181801 (2015) doi:10.1103/PhysRevLett.115.181801 [arXiv:1506.02661 [hep-ph]].

[7] T. Hurth, F. Mahmoudi and S. Neshatpour, Nucl. Phys. B **909**, 737 (2016) doi:10.1016/j.nuclphysb.2016.05.022 [arXiv:1603.00865 [hep-ph]].

[8] W. Altmannshofer, P. Stangl and D. M. Straub, Phys. Rev. D **96**, 055008 (2017) doi:10.1103/PhysRevD.96.055008 [arXiv:1704.05435 [hep-ph]].

[9] A. Antognini *et al.*, Science **339**, 417 (2013) doi:10.1126/science.1230016.

[10] R. Pohl, R. Gilman, G. A. Miller and K. Pachucki, Ann. Rev. Nucl. Part. Sci. **63**, 175 (2013) doi:

10.1146/annurev-nucl-102212-170627 [arXiv:1301.0905 [physics.atom-ph]].

[11] R. J. Hill and G. Paz, Phys. Rev. D **82**, 113005 (2010) doi:10.1103/PhysRevD.82.113005 [arXiv:1008.4619 [hep-ph]].

[12] D. Tucker-Smith and I. Yavin, Phys. Rev. D **83**, 101702 (2011) doi:10.1103/PhysRevD.83.101702 [arXiv:1011.4922 [hep-ph]].

[13] B. Batell, D. McKeen and M. Pospelov, Phys. Rev. Lett. **107**, 011803 (2011) doi:10.1103/PhysRevLett.107.011803 [arXiv:1103.0721 [hep-ph]].

[14] C. E. Carlson, Prog. Part. Nucl. Phys. **82**, 59 (2015) doi:10.1016/j.ppnp.2015.01.002 [arXiv:1502.05314 [hep-ph]].

[15] T. Hur, D. W. Jung, P. Ko and J. Y. Lee, Phys. Lett. B **696**, 262 (2011) doi:10.1016/j.physletb.2010.12.047 [arXiv:0709.1218 [hep-ph]].

[16] P. Ko, Int. J. Mod. Phys. A **23**, 3348 (2008) doi:10.1142/S0217751X08042109 [arXiv:0801.4284 [hep-ph]].

[17] T. Hur and P. Ko, Phys. Rev. Lett. **106**, 141802 (2011) doi:10.1103/PhysRevLett.106.141802 [arXiv:1103.2571 [hep-ph]].

[18] S. Baek, P. Ko and W. I. Park, JHEP **1202**, 047 (2012) doi:10.1007/JHEP02(2012)047 [arXiv:1112.1847 [hep-ph]].

[19] S. Baek, P. Ko, W. I. Park and E. Senaha, JHEP **1211**, 116 (2012) doi:10.1007/JHEP11(2012)116 [arXiv:1209.4163 [hep-ph]].

[20] S. Baek, P. Ko, W. I. Park and E. Senaha, JHEP **1305**, 036 (2013) doi:10.1007/JHEP05(2013)036 [arXiv:1212.2131 [hep-ph]].

[21] S. Baek, P. Ko and W. I. Park, JHEP **1307**, 013 (2013) doi:10.1007/JHEP07(2013)013 [arXiv:1303.4280 [hep-ph]].

[22] S. Choi, S. Jung and P. Ko, JHEP **1310**, 225 (2013) doi:10.1007/JHEP10(2013)225 [arXiv:1307.3948 [hep-ph]].

[23] S. Baek, P. Ko and W. I. Park, JCAP **1410**, 067 (2014) doi:10.1088/1475-7516/2014/10/067 [arXiv:1311.1035 [hep-ph]].

[24] P. Ko and Y. Tang, JCAP **1405**, 047 (2014) doi:10.1088/1475-7516/2014/05/047 [arXiv:1402.6449 [hep-ph]].

[25] P. Ko and Y. Tang, Phys. Lett. B **739**, 62 (2014) doi:10.1016/j.physletb.2014.10.035 [arXiv:1404.0236 [hep-ph]].

[26] P. Ko, W. I. Park and Y. Tang, JCAP **1409**, 013 (2014) doi:10.1088/1475-7516/2014/09/013 [arXiv:1404.5257 [hep-ph]].

[27] J. Kim, P. Ko and W. I. Park, JCAP **1702**, 003 (2017) doi:10.1088/1475-7516/2017/02/003 [arXiv:1405.1635 [hep-ph]].

[28] S. Baek, P. Ko and W. I. Park, Phys. Rev. D **90**, 055014 (2014) doi:10.1103/PhysRevD.90.055014 [arXiv:1405.3530 [hep-ph]].

[29] P. Ko and Y. Tang, JCAP **1501**, 023 (2015) doi:10.1088/1475-7516/2015/01/023 [arXiv:1407.5492 [hep-ph]].

[30] S. Baek, P. Ko and W. I. Park, Phys. Lett. B **747**, 255 (2015) doi:10.1016/j.physletb.2015.06.002 [arXiv:1407.6588 [hep-ph]].

[31] P. Ko, W. I. Park and Y. Tang, arXiv:1504.06944 [hep-ph].

[32] P. Ko and Y. Tang, Phys. Lett. B **741**, 284 (2015) doi:10.1016/j.physletb.2015.01.001 [arXiv:1410.7657 [hep-ph]].

[33] P. Ko and Y. Tang, JCAP **1602**, 011 (2016) doi:10.1088/1475-7516/2016/02/011 [arXiv:1504.03908 [hep-ph]].

[34] S. Baek, P. Ko, M. Park, W. I. Park and C. Yu, Phys. Lett. B **756**, 289 (2016) doi:10.1016/j.physletb.2016.03.026 [arXiv:1506.06556 [hep-ph]].

[35] P. Ko and Y. Tang, Phys. Lett. B **751**, 81 (2015) doi:10.1016/j.physletb.2015.10.021 [arXiv:1508.02500 [hep-ph]].

[36] P. Ko and H. Yokoya, JHEP **1608**, 109 (2016) doi:10.1007/JHEP08(2016)109 [arXiv:1603.04737 [hep-ph]].

[37] P. Ko, A. Natale, M. Park and H. Yokoya, JHEP **1701**, 086 (2017) doi:10.1007/JHEP01(2017)086 [arXiv:1605.07058 [hep-ph]].

[38] H. Hatanaka, D. W. Jung and P. Ko, JHEP **1608**, 094 (2016) doi:10.1007/JHEP08(2016)094 [arXiv:1606.02969 [hep-ph]].

[39] P. Ko and Y. Tang, Phys. Lett. B **762**, 462 (2016) doi:10.1016/j.physletb.2016.10.001 [arXiv:1608.01083 [hep-ph]].

[40] K. Hashino, M. Kakizaki, S. Kanemura, P. Ko and T. Matsui, Phys. Lett. B **766**, 49 (2017) doi:10.1016/j.physletb.2016.12.052 [arXiv:1609.00297 [hep-ph]].

[41] P. Ko and Y. Tang, Phys. Lett. B **768**, 12 (2017) doi:10.1016/j.physletb.2017.02.033 [arXiv:1609.02307 [hep-ph]].

[42] P. Ko and J. Li, Phys. Lett. B **765**, 53 (2017) doi:10.1016/j.physletb.2016.11.056 [arXiv:1610.03997 [hep-ph]].

[43] S. Baek, P. Ko and J. Li, Phys. Rev. D **95**, 075011 (2017) doi:10.1103/PhysRevD.95.075011 [arXiv:1701.04131 [hep-ph]].

[44] T. Kamon, P. Ko and J. Li, Eur. Phys. J. C **77**, 652 (2017) doi:10.1140/epjc/s10052-017-5240-8 [arXiv:1705.02149 [hep-ph]].

[45] P. Ko, N. Nagata and Y. Tang, Phys. Lett. B **773**, 513 (2017) doi:10.1016/j.physletb.2017.08.065 [arXiv:1706.05605 [hep-ph]].

[46] Z. Kang, P. Ko and T. Matsui, JHEP **1802**, 115 (2018) doi:10.1007/JHEP02(2018)115 [arXiv:1706.09721 [hep-ph]].

[47] B. Dutta, T. Kamon, P. Ko and J. Li, Eur. Phys. J. **C 78**, 595 (2018) doi:10.1140/epjc/s10052-018-6071-y [arXiv:1712.05123 [hep-ph]].

[48] S. M. Choi, H. M. Lee, P. Ko and A. Natale, Phys. Rev. **D 98**, 015034 doi:10.1103/PhysRevD.98.015034 [arXiv:1801.07726 [hep-ph]].

[49] K. Hashino, M. Kakizaki, S. Kanemura, P. Ko and T. Matsui, JHEP **1806**, 088 (2018) doi: 10.1007/JHEP06(2018)088 [arXiv:1802.02947 [hep-ph]].

[50] A. Djouadi, O. Lebedev, Y. Mambrini and J. Quevillon, Phys. Lett. B **709**, 65 (2012) doi: 10.1016/j.physletb.2012.01.062 [arXiv:1112.3299 [hep-ph]].

[51] A. Djouadi, A. Falkowski, Y. Mambrini and J. Quevillon, Eur. Phys. J. C **73**, 2455 (2013)

doi:10.1140/epjc/s10052-013-2455-1 [arXiv:1205.3169 [hep-ph]].

[52] O. Lebedev, H. M. Lee and Y. Mambrini, Phys. Lett. B **707**, 570 (2012) doi:10.1016/j.physletb.2012.01.029 [arXiv:1111.4482 [hep-ph]].

[53] A. De Simone, G. F. Giudice and A. Strumia, JHEP **1406**, 081 (2014) doi:10.1007/JHEP06(2014)081 [arXiv:1402.6287 [hep-ph]].

[54] G. Aad *et al.* [ATLAS Collaboration], Phys. Rev. Lett. **112**, 201802 (2014) doi:10.1103/PhysRevLett.112.201802 [arXiv:1402.3244 [hep-ex]].

[55] S. Chatrchyan *et al.* [CMS Collaboration], Eur. Phys. J. C **74**, 2980 (2014) doi:10.1140/epjc/s10052-014-2980-6 [arXiv:1404.1344 [hep-ex]].

[56] G. Aad *et al.* [ATLAS Collaboration], Eur. Phys. J. C **75**, 299 (2015) Erratum: [Eur. Phys. J. C **75**, 408 (2015)] doi:10.1140/epjc/s10052-015-3517-3, 10.1140/epjc/s10052-015-3639-7 [arXiv:1502.01518 [hep-ex]].

[57] V. Khachatryan *et al.* [CMS Collaboration], JHEP **1506**, 121 (2015) doi:10.1007/JHEP06(2015)121 [arXiv:1504.03198 [hep-ex]].

[58] P. Athron, J. M. Cornell, F. Kahlhoefer, J. Mckay, P. Scott and S. Wild, arXiv:1806.11281 [hep-ph].

[59] L. M. Krauss and F. Wilczek, Phys. Rev. Lett. **62**, 1221 (1989) doi:10.1103/PhysRevLett.62.1221.

[60] G. Belanger, K. Kannike, A. Pukhov and M. Raidal, JCAP **1301**, 022 (2013) doi:10.1088/1475-7516/2013/01/022 [arXiv:1211.1014 [hep-ph]].

[61] G. 't Hooft, Nucl. Phys. B **79**, 276 (1974). doi:10.1016/0550-3213(74)90486-6.

[62] A. M. Polyakov, JETP Lett. **20**, 194 (1974) [Pisma Zh. Eksp. Teor. Fiz. **20**, 430 (1974)].

[63] T. Hambye, JHEP **0901**, 028 (2009) doi:10.1088/1126-6708/2009/01/028 [arXiv:0811.0172 [hep-ph]].

[64] A. Falkowski, J. T. Ruderman and T. Volansky, JHEP **1105**, 106 (2011) doi:10.1007/JHEP05(2011)106 [arXiv:1101.4936 [hep-ph]].

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

[66] Y. Hochberg, E. Kuflik, T. Volansky and J. G. Wacker, Phys. Rev. Lett. **113**, 171301 (2014) doi:10.1103/PhysRevLett.113.171301 [arXiv:1402.5143 [hep-ph]].

[67] Y. Hochberg, E. Kuflik, H. Murayama, T. Volansky and J. G. Wacker, Phys. Rev. Lett. **115**, 021301 (2015) doi:10.1103/PhysRevLett.115.021301 [arXiv:1411.3727 [hep-ph]].

[68] F. L. Bezrukov and M. Shaposhnikov, Phys. Lett. B **659**, 703 (2008) doi:10.1016/j.physletb.2007.11.072 [arXiv:0710.3755 [hep-th]].

[69] F. Bezrukov and M. Shaposhnikov, JHEP **0907**, 089 (2009) doi:10.1088/1126-6708/2009/07/089 [arXiv:0904.1537 [hep-ph]].

[70] Y. Hamada, H. Kawai, K. Y. Oda and S. C. Park, Phys. Rev. Lett. **112**, 241301 (2014) doi:10.1103/PhysRevLett.112.241301 [arXiv:1403.5043 [hep-ph]].

[71] F. Bezrukov and M. Shaposhnikov, Phys. Lett. B **734**, 249 (2014) doi:10.1016/j.physletb.2014.05.074 [arXiv:1403.6078 [hep-ph]].

[72] J. Lesgourgues, G. Marques-Tavares and M. Schmaltz, JCAP **1602**, 037 (2016) doi:10.1088/1475-7516/2016/02/037 [arXiv:1507.04351 [astro-ph.CO]].

[73] Y. Tang, Phys. Lett. B **757**, 387 (2016) doi:10.1016/j.physletb.2016.04.026 [arXiv:1603.00165 [astro-ph.CO]].

[74] M. A. Buen-Abad, M. Schmaltz, J. Lesgourgues and T. Brinckmann, JCAP **1801**, 008 (2018) doi:10.1088/1475-7516/2018/01/008 [arXiv:1708.09406 [astro-ph.CO]].

[75] C. Caprini, R. Durrer and G. Servant, Phys. Rev. D **77**, 124015 (2008) doi:10.1103/PhysRevD.77.124015 [arXiv:0711.2593 [astro-ph]].

[76] P. Ko, New Physics: Sae Mulli **66**, 966 (2016). doi:10.3938/NPSM.66.966

[77] M. Buschmann, J. Kopp, J. Liu and P. A. N. Machado, JHEP **1507**, 045 (2015) doi:10.1007/JHEP07(2015)045 [arXiv:1505.07459 [hep-ph]].

[78] M. Zhang, M. Kim, H. S. Lee and M. Park, arXiv:1612.02850 [hep-ph].

[79] T. Cohen, M. Lisanti, H. K. Lou and S. Mishra-Sharma, JHEP **1711**, 196 (2017) doi:10.1007/JHEP11(2017)196 [arXiv:1707.05326 [hep-ph]].

[80] J. Chen, P. Ko, H. N. Li, J. Li and H. Yokoya, arXiv:1807.00530 [hep-ph].

[81] M. Cirelli, P. Panci, K. Petraki, F. Sala and M. Taoso, JCAP **1705**, 036 (2017) doi:10.1088/1475-7516/2017/05/036 [arXiv:1612.07295 [hep-ph]].

[82] A. Mitridate, M. Redi, J. Smirnov and A. Strumia, JCAP **1705**, 006 (2017) doi:10.1088/1475-7516/2017/05/006 [arXiv:1702.01141 [hep-ph]].