Brightening the (130 GeV) Gamma-Ray Line

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Abstract

The gamma-ray line from dark matter (DM) annihilation is too weak to observe, but its observation will uncover much information, e.g., the DM mass and an anomalously large annihilation rate $\sim 0.1$ pb into di-photon. In this work, we construct a minimal effective theory (EFT) incorporating DM and heavier charged particles. A large annihilation rate is obtained from operator coefficients with resonance or strong coupling enhancement. The EFT is stringently constrained by the XENON100 and WMAP data. Without resonance, Dirac DM or colored charged particles are ruled out. It is pointed out that the di-gluon mode may correctly determine the DM relic density. Interestingly, this framework also provides an origin for the Higgs di-photon excess at the LHC. We apply the general analysis to the NMSSM, which can elegantly interpret the tentative 130 GeV gamma-ray line. A top-window model is also proposed to explain the gamma-ray line.

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I. INTRODUCTION AND MOTIVATIONS

The existence of dark matter (DM) has been confirmed by its gravitational effects, and its energy fraction $\sim 25\%$ today is also measured. However, the conclusive evidences that may reveal the DM particle properties are still absent. Among a variety of (indirect) detecting objects on DM, the gamma-ray from the DM dense region (such as the center of the Galaxy) is especially promising by virtue of weak astrophysical influence on its propagation [1]. Of particular interest is the monochromatic gamma-ray line, which has rather clear background. But it is highly suppressed because the DM $\chi$ can only annihilate to photons via the charged loop.

However, once such a spectral line is observed, it will uncover very important information of DM. In this article, we assume an extracted DM mass from $E_\gamma$ and an anomalously large annihilation rate $\langle \sigma v \rangle_{2\gamma} \sim 0.1$ pb into di-photon (it is taken as a referred value throughout the work, unless specified), then attempt to reconstruct the DM properties and dynamics to the most extent. Inspired by the recent discovery of a gamma-ray line at $E_\gamma \simeq 130$ GeV, which is claimed in the Ref. [2, 3] after re-analyzing the Fermi Large Area Telescope (FERMI-LAT) data published in 2009 [4], it is conjectured that the line may originate from DM annihilating into gamma. Best fit of the data shows a DM of mass around 130 GeV and annihilation rate at level 0.1 pb. Later independent analysis also confirms the line [5]. The line has a sharper peak which is hard to explain by FERMI-bubbles [6], while DM+DM $\rightarrow \gamma\gamma$ gives a better fit [5]. The Ref. [7] also shows a strong evidence of the gamma-ray from the inner galaxy and draws a similar conclusion. This line has received much attention from astrophysics [5, 6, 8, 9] and particle physics [10]. In spite of queries [8], the gamma-ray line from DM activity itself is of great theoretical interest, and deserving a deep study.

The topic can be studied in the effective theory (EFT) framework, by minimally including an operator $a_C\chi^\dagger \chi C^\dagger C$ where $C$ is the charged particle. The anomalously bright gamma-ray line is due to large $a_C$. We further demand the EFT be compatible with other constraints on the DM, i.e., the WMAP and XENON100 bound [12]. Independent of the mechanism generating large $a_C$, we can arrive:

- The charged particle $C$ in the loop should be heavier than the DM, otherwise it would render too large annihilation rate into $CC$, which leads to too small DM relic density. On top of that, the injection from such a large flux of charged particles into the cosmic-ray probably has been excluded by the PAMELA.

- The charged particle carrying both QED and QCD charges needs careful inspections. Along with the di-photon annihilating mode, there is an enhanced di-gluon mode with estimated rate $\langle \sigma v \rangle_{2G} \sim 0.1(\alpha_s^2/\alpha^2)\langle \sigma v \rangle_{2\gamma} \simeq 1$ pb, which makes an illustrative
coincidence.

The large $a_C$ can be generated through Breit-Weigner resonance mechanism, or the strong interaction between DM and the charged loop. For the former scenario, properties of the scalar/vector resonance can be further stringently restricted by symmetries, e.g. the CP, and the above consideration. The next-to-minimal supersymmetric standard model (NMSSM) [13] is a good realization of this scenario. We find it is capable of interpreting the tentative 130 GeV gamma-ray line. For the latter scenario, the XENON100 bound excludes the Dirac DM, as well as both Dirac and Majorana DM if the charged particle $C$ carries color. Interestingly, in any scenario, the possible SM-like Higgs $h$ to di-photon excess at the LHC [14] may share the same origin, if we incorporate the operator $a_C^h h C^i C$.

This paper is organized as following: In the section II, we perform a general analysis based on the minimal EFT. In the next two sections exploration on the enhancement mechanism is presented. The Section V includes the conclusion and discussions. And some necessary complementarity is casted in the Appendix.

II. GENERALITY AND GUIDANCE

As is well known, the DM can not directly annihilate into photons due to its QED charge neutrality, while transition at the loop level is generically highly suppressed. So, it is nontrivial to obtain an abnormally large annihilating rate, saying $\langle \sigma v \rangle_{2\gamma} \sim 0.1$ pb. Some more powerful model independent statements can be made, if it is further combined with other aspects of DM. To see that, we consider the minimal effective operators [42] relevant to the gamma-ray line anomaly

Fermion DM : $a_C \bar{\chi} \Gamma \chi \bar{C} C$, $a_{\bar{C}} \bar{\chi} \Gamma \chi \bar{C} C$, $a_W \bar{\chi} \Gamma \chi W^+ W^-$, \hspace{1cm} (1)

Scalar DM : $a_t \chi^\dagger \chi \bar{C} C$, $a_{\tilde{w}} \chi^\dagger \chi W^+ W^-$, \hspace{1cm} (2)

where the gamma matrix $\Gamma \subset \{1, \gamma^5, \gamma^\mu, \sigma^{\mu\nu}\}$. Lorentz and $SU(3)_C \times U(1)_{QED} \times CP$ symmetries are implied, and some operators will vanish due to these symmetries. $C$ is a charged fermion and $\bar{C}$ is a charged scalar, both of which are not confined to the SM. However, in the sense of inducing DM annihilating into gamma, the scalar loop is not as effective as the fermionic loop, unless there is a large enhancement from highly charged particles such as a double charged scalar. Note that operators containing $H^+ W^-$ where $H^+$ is the charged Higgs from 2HDM-like model are not included, since their contribution are always putative null, e.g., in the non-linear unitary gauge the vertex $H^+ W^- \gamma$ vanishes [15].

To achieve a large $\langle \sigma v \rangle_{2\gamma}$ and maintain the main merit of DM dynamics, the candidates running in the charged loop are more or less selected. Denoting by $S_i$ the set of DM $2 \rightarrow 2$
annihilation mode and without loss of generality, let $S_1$ be the one from which di-photons come after closing the charged states to form a loop. Some cases arise:

- If $S_1$ is the on-shell type with final states $X_1 \bar{X}_1$, then $\langle \sigma v \rangle_{S_1} \sim 10^4 \langle \sigma v \rangle_{S_1,2\gamma} \sim 10^3$ pb. Injection from such a large flux charged particles would have been observed by PAMELA [16] from the significant excess of the positron or anti-proton flux. On top of that, the DM relic density would be too small, unless we consider the subtle thermal Breit-Weigner enhancement effect which is active only today [17].

- If $S_1$ is properly off-shell, then the above problem is resolved since the $S_1$ is forbidden today. But we have to examine its annihilation rate at the early universe, i.e., comparing the rate of the forbidden annihilation mode $\langle \sigma v \rangle_{S_1,T_f}$ with 1 pb, where $T_f = m_\chi/x_f$ with $x_f \sim 25$ the typical decoupling temperature of DM. Generically one can expand the annihilation rate as [18]

$$\langle \sigma v \rangle_{S_1} = (a + b/x_f)v_2,$$

where the final two-body phase space gives the velocity of out-going particles in the CM frame, $v_2 = (1 - z^2 + z^2 v_{rel}^2/4)^{1/2}$ with $z = m_{X_1}/m_\chi$ and $v_{rel}$ the relative velocity of initial particles. For a properly large $z > 1$, the relic density of DM is [18]

$$\Omega h^2 = \frac{1.07 \times 10^9 x_f}{g_*^{1/2} M_{Pl} J}, \quad J \simeq a \frac{z}{\mu_f} e^{-\mu_+ x_f},$$

where $\mu_+ = (1 - 1/z^2)^{1/2}$. The relic density is very sensitive to $z$, e.g., from $z = 1.05$ to $z = 1.10$, it increases roughly one order. But in principle it is possible to obtain $\langle \sigma v \rangle_{2\gamma} \sim 0.1$ pb and $\langle \sigma v \rangle_{X_1 \bar{X}_1} \simeq 1$ pb simultaneously, if we accept large fine-tuning.

- If $z$ is large enough then the forbidden channel is completely ignorable. As a consequence, we need a new (dominant) channel $S_2$ to reduce the DM number density. In actual model building, it naturally happens. But a more interesting case arises as following.

- If the charged particle also carry color charge, then the $2\gamma$ mode is subdominant to the two gluon mode:

$$\frac{\langle \sigma v \rangle_{2G}}{\langle \sigma v \rangle_{2\gamma}} \sim 0.1 \frac{\alpha_s^2}{\alpha^2} \simeq \mathcal{O}(20),$$

which is estimated in light of fermions with unit charge, and the origin of numerical factor 0.1 can be traced back to the property of charged particles. This numerical coincidence means that if the two-gamma rate is $\sim 0.1$ pb, the right relic density is achieved via the di-gluon mode. For the proof of that the di-$Z$ mode is at most the same order of di-photon mode and thus irrelevant, see the Appendix A for details.
Generically, we will have several annihilation modes producing gamma line, $\chi\bar{\chi} \rightarrow \gamma\gamma, \gamma X$ with $X = Z, h$. The first mode creates a line at $E_\gamma = m_\chi$ while the second mode produces a line with lower energy $E'_\gamma \approx E_\gamma(1 - m_X^2/m_\chi^2)^{1/2}$. They have comparable cross sections except for a very significant phase space suppress. Therefore we focus on the two gamma final states, since the two modes quantitatively differ only by some constant, such as the difference between $e$ and $g_2$. However, how to distinguish the two lines is very interesting as discussed in Ref. [19].

To end up the general discussion in EFT, we would like to mention that there is a possible relation between the di-phonon excess for the Higgs search at the CMS/ATLAS [14] and for the DM search at the sky. The common point is the new charged loop. Through the same loop alone which the DM annihilates into two photons, the SM-like Higgs $h$ can decay into two photons with appreciable width, if the coupling to the Higgs is significantly. It is can be described simply by further including the effective operators $a_h^c h \bar{C} C$ or $a_h^h \bar{C} \tilde{C}^\dagger \tilde{C}$. However we are not going to discuss this in detail due to its triviality in the EFT.

### III. ENHANCEMENT FROM RESONANCE

In this section, we present the effective analysis by specifying the role of resonance. The simple top window model is constructed, and in particular we survey its implication on conventional supersymmetric model such as the NMSSM.

#### A. Scalar resonance

We consider the scalar resonance which appears almost everywhere in models with extended Higgs sector. To get a sufficiently large enhancement, $s-$channel resonant annihilation is the most conventional mechanism. In this case, the resonance $\phi$ takes mass $m_\phi \simeq 2m_{\text{DM}}$, and the cross section manifest of the enhancement can be parameterized as

$$\sigma v = \frac{T_I T_F}{32\pi} \frac{1}{m_\chi^2} |M|^2 \sim \frac{\alpha^2 m_\phi^2}{m_\chi^4} \frac{1}{(1 - r)^2 + \gamma},$$

(6)

where $r = 4m_X^2/m_\phi^2$ and $\gamma = (\Gamma_\phi/m_\phi)^2 \ll 1$. Here $\alpha$ stands for an effective coupling and will be specified in concrete examples, while $T_{I,F}$ takes 1/2 or 1/4 and so on, standing for the average of initial degree of freedoms or the symmetry factor of final states.

It is convenient to define $f_B = 1/((1 - r)^2 + \gamma)$, then the DM annihilating cross section into $X_i\bar{X}_i$ (thought $\phi$) can be rewritten as

$$(\sigma v)_{X_i} = \frac{f_B T_I}{m_\phi^4 m_\chi^2} |M(\chi\chi \rightarrow \phi)|^2 \Gamma(\phi \rightarrow X_i\bar{X}_i).$$

(7)
Without loss of generality, we take $X_2$ as the dominant mode, and immediately get the upper bound of the branching ratio of $\phi$ decay (to particles other than di-photon):

$$\frac{\text{Br}(\phi \to X_2 \bar{X}_2)}{\text{Br}(\phi \to \gamma \gamma)} = \frac{(\sigma v)_{S_2}}{(\sigma v)_{2\gamma}} \lesssim 10.$$  \hspace{1cm} (8)

To arrive it we have set $(\sigma v)_{S_2} \sim 1 \text{ pb}$ as the standard annihilation rate as well as the referred value $(\sigma v)_{2\gamma} \sim 0.1 \text{ pb}$. Therefore we get a model independent bound $\text{Br}(\phi \to \gamma \gamma) \gtrsim 10\%$, as negates the resonance from simple two-Higgs-doublet-model (2HDM) by virtue of their considerably coupling to fermions or light massive vector boson (it is absence for CP-odd Higgs). This bound has far-reaching implication on the collider. Provided that the production cross section of $\phi$ is sufficiently large (for example, when the charged loop meanwhile carries color as in the top-window model discussed later), the gamma-ray line observed at the sky predicts a clear di-photon excess at the peak around $2m_\chi (\approx 260 \text{ GeV})$ at the LHC.

1. **Effective analysis**

In light of previous arguments, we need some rather heavy charged particles, while in the SM top quark and $W$ boson are the only two charged particles of mass around the weak scale. Accordingly, a $m_\chi > m_t$ hints a new charged particle. When $m_t > m_\chi > m_W$, the top quark will open a unique window. Otherwise, $W$ may run in the charged loop provided a vertex $\phi W^+W^-$, that implies nontrivially the identity of $\phi$.

Since the effective operators listed in the previous section are ascribed to the integrating out $s-$channel scalar resonance, the set of possible operators can be reduced greatly:

**Fermionic DM**

$$a_C \bar{\chi}(\gamma^5)\chi \bar{C}(\gamma^5)C, \quad a_C \bar{\chi}(\gamma^5)\chi \bar{C}^\dagger \tilde{C}, \quad a_W \bar{\chi} \chi W^+W^-.$$  \hspace{1cm} (9)

**Scalar DM**

$$a_C \bar{\chi}^\dagger \chi \bar{C}(\gamma^5)C, \quad a_C \bar{\chi}^\dagger \chi \bar{C}^\dagger \tilde{C}, \quad a_W \bar{\chi}^\dagger \chi W^+W^-.$$  \hspace{1cm} (10)

The coefficients $a$’s are proportional to $f_B^{1/2}/m_\phi^2$. And $\gamma^5$ in the parenthesis may or may not appear, depending on the CP quantum number of $\phi$. But specified to fermionic DM, $\gamma^5$ must be inserted so as to make the present DM annihilating rate avoid acute velocity suppressing, $v^2 \sim 10^{-6}$. This means the $\phi$ must be CP-odd, denoted as $\phi_A$ hereafter. As an immediate consequence, the $W-$window is closed. Additionally, $\bar{C}^\dagger \tilde{C}$ should be understood as $\bar{C}_L^\dagger \tilde{C}_R$, under CP transformation $\bar{C}_L^\dagger \tilde{C}_R \to \bar{C}_R^\dagger \tilde{C}_L$. However, QED does not change the chirality, thus we need a further large LR mixing. As an example, the stop-system in the SUSY just satisfies those requirements. On the contrary, for scalar DM, no matter complex or real, the $\phi$ should be CP-even and denoted as $\phi_h$. Hence the $\gamma^5$ should be removed.
We would like to add some further remarks. Firstly, the Lorentz and SM-gauge invariance force $\phi$ either transforms non-trivially under the $SU(2)_L$ symmetry or mixes with such states. Secondly, the constraint on the $\phi$ interactions indicated by the Eq. (8) should be satisfied. Finally, the $\phi$ also mediates the tree-level DM-nucleon interaction in the presence of a top window. Although this contribution is suppressed by velocity for the fermionic DM, the scalar DM requires inspection [20]. The resulting DM-proton inelastic scattering cross section is

$$ \sigma_{SI} = \frac{4\mu_p^2}{\pi} f_p^2, \quad f_p \approx \frac{2}{24} f_T \frac{a_t}{m_p/s_{m_t},} $$  \hspace{1cm} (11)$$

where $\mu_p \approx m_p$ is the DM-proton reduced mass and $f_T^{1/2} \approx 0.83$ [20]. Note the enhancement factor $f_B^{1/2}$ in $a_t$ is removed when we are calculating the DM-nucleon recoil rate using the effective operator. The present exclusion on $f_p$ is $10^{-8}$ GeV$^{-2}$ for DM of mass 100 GeV, put by the XENON100 [12]. It implies the upper bound

$$ a_t \lesssim 3.2 f_p m_p m_t / m_p \approx 2.8 \times 10^{-3} \left( \frac{f_p}{10^{-8} \text{ GeV}^{-2}} \right) \left( \frac{m_p}{100 \text{ GeV}} \right) \text{ GeV}^{-1}, $$  \hspace{1cm} (12)$$

which places a rather strong constraint for the top-window model.

Having outlined the most essential profile of WIMP that potentially has bright gamma-ray lines, we continue to make some quantitative discussion. Effectively, through the charged loop, Lorentz and CP invariance leads to the following operators for the CP-even and CP-odd $\phi$ respectively

$$ \alpha \frac{h_{\phi CC}^1}{4\pi} \frac{1}{4\Lambda_1} \bar{\phi} h F_{\mu \nu} F^{\mu \nu}, \quad \alpha \frac{h_{\phi CC}^2}{4\pi} \frac{1}{8\Lambda_2} \bar{\phi} A F_{\mu \nu} F^{\mu \nu}, $$  \hspace{1cm} (13)$$

with $\alpha \approx 1/137$. $h_{\phi CC}$ is the coupling constant between $\phi$ and charged particles. Factoring out the loop factor and couplings, moreover multiplying $1/4(8)$ for later convenience, the $\Lambda_{1,2}$ can be much below the weak scale (it is even enhanced by color or electric charge). Concrete expressions for the effective scale are casted in the Appendix A.

Now we are at the position to evaluate the DM annihilation rate into gamma pair. Denoting by $\Gamma_{\phi \gamma \gamma}^{\mu \nu}$ the Feynman rules (see Fig. 5 for label) of Eq. (13), allowing for off-shell $\phi$, they are respectively given by [43]

$$ \phi \text{ CP - even : } \Gamma_{\phi \gamma \gamma}^{\mu \nu} (p_1, p_2, P) = \frac{\alpha h_{\phi CC}}{4\pi} \frac{1}{4\Lambda_1} [p_1 \cdot p_2 g^{\mu \nu} - (p_1^\mu p_2^\nu + p_2^\mu p_1^\nu)], $$  \hspace{1cm} (14)$$

$$ \phi \text{ CP - odd : } \Gamma_{\phi \gamma \gamma}^{\mu \nu} (q_1, q_2, P) = \frac{\alpha h_{\phi CC}}{4\pi} \frac{1}{8\Lambda_2} \epsilon^{\mu \nu \alpha \beta} (p_1)_\alpha (p_2)_\beta. $$  \hspace{1cm} (15)$$

After averaging initial states and summing over out-going photon polarization states, the resulting real scalar DM annihilation rate is

$$ (\sigma v)_{2\gamma} \approx \frac{f_B^2}{128\pi} \frac{g_{h \chi \chi}^2}{\Lambda_1^2} \frac{\alpha^2 h_{\phi CC}^2}{16\pi^2} \frac{1}{m_\chi^2} \approx 0.06 \left( \frac{f_B}{500} \right) \left( \frac{30 \text{ GeV}}{\Lambda_1} \right)^2 \left( \frac{g_{h \chi \chi} h_{\phi CC}}{0.3} \right)^2 \text{ pb}, $$  \hspace{1cm} (16)$$
where $g_{h\chi\chi}$ is the coupling constant between $\phi$ and DM. Owing to the scalar DM, we have parameterized the massive coupling as $g_{\phi\chi\chi} \rightarrow 2g'_{\phi\chi\chi}m_\chi$. For the Majorana DM,

$$
(\sigma v)_{2\gamma} \simeq \frac{f_B}{128\pi} \frac{g_{h\chi\chi}^2}{\Lambda^2} \frac{\alpha^2 h_{CC}^2}{16\pi^2} \approx 0.04 \left( \frac{f_B}{100} \right) \left( \frac{30 \text{ GeV}}{\Lambda} \right)^2 \left( \frac{g_{\phi\chi\chi} h_{CC}}{1.0} \right)^2 \text{pb.}
$$

(17)

FIG. 1: Dark matter annihilates into $2\gamma$ via top loop. The cut denotes for annihilates into top pair.

2. The top-window model

In the standard model (SM), lying on the top of the fermion mass ascending order, the top quark may hide some new dynamics. It is thus interesting to conjecture that the dark sector may have a close relation with the top quark, i.e., the dark sector only (strongly) interacting with top quark in the fermion sector [44]. Then it is reasonable to expect enhanced gamma-ray line via top-loop at levels close to the present experimental sensitivity. And interesting, the latest tentative gamma-ray line from dark matter annihilation [3] requires a DM of mass about 130 GeV, just lies within the top-window. Therefore, as a concrete example of the general effective operator analysis, we focus on a top-window model equipped with scalar resonance.

Due to the chiral structure of the SM, it is most likely that the $\phi$ dwells in an extra Higgs doublet $H'$ with hypercharge $+1/2$ so as to couple the top quarks at dimension-four level. A simple effective model, which might be a remnant of top dynamics reads

$$
\mathcal{L} = \frac{1}{2} \chi^2 \left( \mu^2 + \lambda H' H + \text{c.c.} + \kappa_1 |H'|^2 + \kappa_2 |H|^2 \right) + V(H, H') + y_3 \bar{Q}_3 H't_R.
$$

(18)

where $H$ is the ordinary Higgs doublet developing a vacuum expected value (VEV) $v = 174$ GeV. By contrast, $H'$ should have a VEV $\langle H' \rangle \ll v$ which makes the term $\phi W^+ W^-$ negligible [45]. Here we assume that the VEV hierarchy is realized by well organized Higgs potential $V(H, H')$ as discussed in [20]. In addition, since we have ascribed the unique
significant coupling $H'Q_3t_R$ to some unknown dynamics, we are free of the FCNC-problem induced by family changing Yukawa couplings $H'\bar{f}_L f'_R$. In summary, after the EW-breaking, we are left with the relevant terms

$$\mathcal{L} \supset \frac{1}{2} \left( \mu^2 + \kappa_2 v^2 \right) \chi^2 + \frac{1}{2\sqrt{2}} (\lambda v) \chi^2 \phi + \frac{1}{\sqrt{2}} (\kappa_2 v) \chi^2 h + \frac{\kappa_2}{4} \chi^2 h^2 + \left( \frac{y_3}{\sqrt{2}} \phi \bar{t}_L t_R + \text{c.c.} \right).$$  \tag{19}$$

In light of the general analysis, the scalar DM the $\phi$ is identified with the CP-even component of the neutral boson in $H'$. If we drop terms involving $H'$, the model just recovers the Higgs-porting model (see a recent discussion [21]). But since $\kappa_2$ is irrelevant on our purpose, so we can turn it off to reduce parameters.

Now we turn our attention to the phenomenological aspects of the model. First of all, the di-photon rate, in terms of the parameterization in the Eq. (17), is determined by

$$g'_{\phi\chi\chi} = \frac{\lambda}{2^{3/2}} \frac{v}{m_\chi}, \quad h_{\phi CC} = \frac{y_3}{\sqrt{2}}. \tag{20}$$

Taking $\lambda \simeq y_3 \simeq 1$ and $f_B = 500$ leads to a rate $\simeq 0.07$ pb. And the XENON100 constraint Eq. (12) has been arranged to be satisfied by means of rather large $f_B$ thus smaller couplings. Next, the mass difference between $m_t$ and $m_\chi$ is at a few ten percents level, so the forbidden annihilation mode (to top quark) is ignorable in terms of the Eq. (4). The di-gluon mode may properly account for the relic density. From Eq. (A10) and Eq. (A7) it is estimated that

$$\frac{\langle \sigma v \rangle_{2G}}{\langle \sigma v \rangle_{2\gamma}} = \frac{\alpha_s^2}{\alpha^2} \frac{1}{4N_c^2 Q_t^4} \approx 35,$$  \tag{21}

with $Q_t = 2/3$ the top quark charge.

An alternative model is to assume the dark sector consists of a SM vector-like Dirac fermion pair $(\psi, \psi^c)$, which is the the $SU(2)_L$ doublet and carries hypercharge $\pm 1/2$ respectively. Extra singlet fermion $S$ is introduced, then the model is

$$\mathcal{L} = \left( \lambda S \bar{\psi} H' + m_\psi \psi \psi^c + M_S S^2 \right) + V(H, H') + y_3 Q_3 H' t_R. \tag{22}$$

The fermonic DM brings important difference. The $\phi$ should be the CP-odd component of $H'$, which does not couple to $W^+W^-$ thus allows a large $\langle H' \rangle$ (but we have to ensure small $H' - H$ mixing, otherwise $\phi bb$ is too large). It leads to a phenomenologically viable singlet-doublet mixing Majorana DM. It can be regarded as an effective model of the supersymmetric model studied as the following (but top loop will be replaced by chargino loop).

**B. The NMSSM**

In the supersymmetric standard models (SSM), in addition to the (Majorana) neutralino LSP dark matter of mass around $m_Z$, a wealth of new heavy charged particles, CP-odd
Higgs and hence a viable resonance $\phi_A$, are furnished. It is thus of particular interest to investigate whether the above set up can be realized in the SUSY. Furthermore, as stressed at the beginning of this section, the coupling of the resonance to light states are stringently constrained. In spite of difficulties in the minimal-SSM, such an invisible $\phi_A$ can be readily accommodated in the NMSSM following the singlet limit. While the examination on light dark matter limit has been done [22], our scenario has not been currently examined yet as far as we are aware.

FIG. 2: The LSP (with acceptable relic density 0.09-0.12) at the plane of $(\sigma v)_{2\gamma} - \sigma_{SI}^p$. As one can see, the FERMI gamma-ray line can be accommodated even under stringent XENON100 exclusion. The colored notation is of the SM-like Higgs mass. We obtain the results by using the programme NMSSMtools [34].

Owing to the singlet sector, the NMSSM presents a clear realization of our scenario. For our purpose, the model is none other than the $Z_3$-NMSSM:

$$W \supset \lambda S H_u H_d + \frac{\kappa}{3} S^3,$$

$$-L_{soft} \supset m_S^2 |S|^2 + \left( \lambda A \lambda S H_u H_d + A_\kappa \frac{\kappa}{3} S^3 + c.c. \right),$$

Consider a slice of the parameter space: (A) $\lambda \sim 1$ moreover $\kappa \lesssim \lambda$ and $\tan \beta \sim 3$ (favored by naturalness to enhance the SM-like Higgs mass [23]), further the $v_s \equiv \langle S \rangle$ gives rise to the Higgsino-like charginos (with mass roughly $\mu = \lambda v_s \sim 200$ GeV) which replaces the top quark in the gamma loop; (B) The 130 GeV LSP dark matter has significant (even if not dominant) singlet component, and $\kappa S^3/3$ provides the vertex $\kappa A_s \tilde{S}^2$ with unsuppressed coupling. (C) The highly singlet-like CP-odd Higgs $A_1 \simeq A_s$ of mass around 260 GeV offers a proper resonant enhancement. Consequently $A\bar{b}b$ and $AhZ$ can be sufficiently suppressed,
and its only significant coupling is to the chargino $\sim \lambda A_1 \tilde{H}_u^+ \tilde{H}_d^-$. This hiding $A_1$ is a key to reconcile the WMAP and FERMI, and the situation of this parameter space to interpret FERMI is shown in the Fig. 2. We have restrict $m_{A_1}$ falls into 255-265 GeV. As one can see, only a very small portion of the points pass all constraints, labeled as the FERMI-region.

We close this section by making some further comments. The above parameter space is a portion of the natural NMSSM [23]. However, as observed there, the Higgsino generically occupies a large proportion of LSP, and its relic density is too small while $\sigma_{\text{SI}}^p$ is too large. To circumvent those problems, we may have to tune the parameters to obtain a viable LSP, as is reflected in the Fig. 3. From model building, we may simply go to the singlet-port dark sector by simply adding $\eta S\Phi^2/2$, where the lighter $\Phi$ component is the dark matter candidate.

**FIG. 3:** The distribution of the di-photon rate on input coordinates, left: $M_A - \kappa A_\kappa$; Right: $\mu - M_2$. The solid purple circles stand for parameter configurations satisfying WMAP, XENON100 and FERMI. They are well-tuned, since they are rather discrete even in the preferred window. Other parameter settings: $\lambda = 0.65, \kappa : 0.13 - 0.16, \tan \beta : 1.3 - 1.7, m_{Q_3} = m_{\tau_R} = 1000$ GeV, $A_t = 0$ while $A_\kappa$ varies between -400 to -300 GeV.

### C. Vector resonance

In this subsection we turn our attention to the vector resonance. There is only one neutral massive vector boson in the SM as well as its simple extension. Such a resonance means the DM mass should be around $m_Z/2$. However, in models with extended $U(1)_X$ local symmetries, new resonances with free masses are expected. Such models have been proposed in [24, 25], where spectral lines from $\gamma Z$ and/or $\gamma h$ final states are predicted. It is believed that the Landau-Yang theorem [26] excludes the di-photon mode. Nevertheless, this
theorem does not exclude the vertex $Z\gamma\gamma$ with off-shell $m_Z$, please see the Ref. [27] discussing anomalous three vector boson coupling within SM. And a concrete evidence for such coupling can be found in the earlier calculation of neutralino annihilation to di-photon [15], where a pole from $m_Z$ indeed exists.

This pole has deep relation with the anomaly of the theory, and it is non-vanishing if and only if the axial-coupling between $Z'$ and the charged massive fermions is present (mass splitting is needed to spoil thorough anomaly cancellation). As a consequence, in Eq. (9) we are left with the only one operator built from fermionic DM and charged particles:

$$a_C \bar{\chi} \gamma^5 \gamma_\mu \bar{C} \gamma^5 \gamma^\mu C,$$

which indicates $Z'$ can not come from vector-like theories. While for the complex scalar DM, the relevant effective vertex is $\chi^\dagger \partial_\mu \chi Z'_\mu$, which renders the DM annihilation suffering from velocity suppress. So we do not need to consider it here.

Whatever the charged particles are, the $Z'\gamma\gamma$ effective Lagrangian can be built by Lorentz and QED invariance. This leads to the following dimension-six operators

$$\left( \frac{\alpha}{4\pi} \frac{g'_{ZCC}}{4\Lambda_1^2} Z'_S F_\mu F_\alpha + \frac{\alpha}{4\pi} \frac{g'_{ZCC}}{4\Lambda_1^2} Z'_\mu \bar{F}_\mu F_\alpha \right),$$

$$+ \left( \frac{\alpha}{4\pi} \frac{g'_{ZCC}}{4\Lambda_2^2} Z'_A F_\mu F_\alpha + \frac{\alpha}{4\pi} \frac{g'_{ZCC}}{4\Lambda_2^2} Z'_\mu \bar{F}_\mu F_\alpha \right),$$

where the symmetric and antisymmetric 2-rank tensors are defined as $(Z_{S/A})_{\mu\nu} = \partial_\mu Z'_\nu \pm \partial_\nu Z'_\mu$. Under $C$ and $P$ symmetries, the vector field transforms as

$$CV C^{-1} = -V, \quad PV(\bar{x},t) P^{-1} = (-1)^{\mu} V(-\bar{x},t),$$

and $P\partial^\mu P^{-1} = (-1)^{\mu} \partial^\mu$. As a result only $\tilde{\Lambda}_{1,2}$ term conserves the CP, so in this work we only keep them (which is consistent with the presence of $\gamma^5$ in Eq. (26)). From the effective Lagrangian, the annihilation rate is calculated to be (for illustrative purpose we only show the $\tilde{\Lambda}_1$-related part)

$$\langle \sigma v \rangle_{2\gamma} \simeq \frac{\alpha^2 g^2_{Z'\chi\chi} g^2_{Z'CC}}{16\pi^2} \frac{f_B}{64\pi} \left( \frac{m_\chi}{\tilde{\Lambda}_1} \right)^2 \frac{1}{\tilde{\Lambda}_1^2}$$

$$\approx 0.03 \left( \frac{g^2_{Z'\chi\chi} g^2_{Z'CC}}{1.0} \right)^2 \left( \frac{f_B}{500} \right) \left( \frac{m_\chi}{100 \text{ GeV}} \right) \left( \frac{100 \text{ GeV}}{\tilde{\Lambda}_1} \right)^4 \text{ pb},$$

(28)

On the other hand, from direct calculation similar to the Ref. [15], we get

$$\langle \sigma v \rangle_{2\gamma} \simeq N_c^2 Q_c^4 (g^2_{Z'\chi\chi} g^2_{Z'CC}) \frac{\alpha^2 f_B}{64\pi^3 m_\chi^2} |A^{1/2}(\tau)|^2.$$
This leads to the effective scale

\[ \tilde{\Lambda}_1 = \frac{m_\chi}{2Q_C N_C^{1/2}} |\mathcal{A}^{1/2}(\tau)|^{1/2}, \quad (30) \]

where \( \tau = m_C/m_\chi \) and \( |\mathcal{A}^{1/2}(\tau)| = \tau (\arctan 1/\sqrt{\tau - 1})^2 \gtrsim 1 \) for \( \tau > 1 \).

IV. CHARGED-LOOP PORTING DARK MATTER

In this section we consider the scenario in the absence of resonance. Viewing from the UV-completion level (illustratively showed in Fig. 4, where a \( Z_2 \) symmetry can be consistently assigned on DM and charged particle), the operators list in the Eq. (9) are generated in the \( t- \) channel. To get a large annihilation cross section without resonance enhancement, it is expected that there is rather strong coupling between DM and the charged particles. That maybe consistent with the composite dark matter. However, in the case of decaying DM scenario [28], a large coupling constant can be avoided.

FIG. 4: Scalar/Fermionic dark matter annihilates into \( 2\gamma \) via a charged loop. \( Z_2 \) symmetry can be assigned to the dark matter and charged particles.

A. Annihilating scenario

In all cases, the heavy charged loop is the major port between DM and visible sector. This is a reminiscence of the dipole dark matter theory, where the DM-photon interactions are the leading order of DM-visible interactions. Considering the fermionic DM case, the effective operators up to dimension-seven should be incorporated:

\[ -\frac{\lambda_\chi}{2} \langle \bar{\chi} \sigma_{\mu\nu} \chi \rangle F_{\mu\nu}, \quad -i \frac{\lambda}{2} \langle \bar{\chi} \sigma_{\mu\nu} \gamma^5 \chi \rangle F_{\mu\nu}, \quad (31) \]

\[ \frac{e}{\Lambda_3^3} \bar{\chi} \chi F_{\mu\nu} F_{\mu\nu}, \quad i \frac{e}{\Lambda_3^3} \bar{\chi} \gamma^5 \chi F_{\mu\nu} \bar{F}_{\mu\nu}, \quad (32) \]
Eq. (31) are dubbed magnetic momentum DM (MDM) and electronic momentum DM (EDM) respectively, with $\lambda_\chi / d_\chi$ the magnetic/electric dipole momentum. If the DM is a Majorana fermion, Eq. (31) vanishes and we only have to consider Eq. (32). Otherwise, we expect all operators listed in Eq. (31) and Eq. (32) are comparably important, since they are generated at the same loop-level. The naive dimension counting gives

$$1/\lambda_\chi \sim \frac{e}{16\pi^2} \frac{g_{hcc}^2}{m_C}, \quad \Lambda_3 \sim \Lambda_4 \sim (m_C^2/\lambda_\chi)^{1/3},$$

(33)

up to an overall operator coefficients loop factors $f(m_C^2/m_C^2)$, whose exact expression is very involved due to multi propagators, and we leave it for further work.

Before dealing with the annihilating rate into gamma pair, we consider the possible bound on operators. In spite of loop suppression, the M(E)DM has long-distance interactions, mediated by the photon, which lead to a great enhancement on $\sigma_{SI}$. Especially, the EDM has a further $1/v^2 \sim 10^6$ enhancement [35, 36], consequently the XENON100 tightly bounds on them. The DM-proton (only for proton by virtue of the QED mediator) cross section is $\sigma_{SI}^p = \alpha \lambda_\chi^2$. In the light of Ref. [36], for a 100 GeV DM, the upper bound is roughly

$$\lambda_\chi \lesssim 10^{-19} \text{ e cm} \approx 1.5 \times 10^{-6} \text{ GeV}^{-1},$$

(34)

which indicates $m_C/g_{hcc}^2 \gtrsim 1$ TeV. In a way similar to the one given in Section. III A 1, the correlation between operator coefficients leads to an estimation rate of the $\bar{\chi}\chi \rightarrow \gamma\gamma$ process

$$(\sigma v)_{2\gamma} \sim \frac{e^2}{64\pi} \lambda_\chi^2 \lesssim 10^{-6} \text{ pb},$$

(35)

which is far below the sensitive bound. Thus the DM can not be a Dirac particle.

Direct detection possibly gives a second constraint, no matter Dirac or Majorana DM. If the charged particles also carry color charge, then after replacing the QED field-strength with the gluon field-strength, we get the dimension-seven operators

$$f_G \frac{\alpha_s}{4\pi} \bar{\chi}\chi G^a_{\mu\nu} G^a_{\mu\nu}, \quad f_{G,5} \frac{\alpha_s}{4\pi} \bar{\chi}\gamma^5\chi G^a_{\mu\nu} \tilde{G}^a_{\mu\nu},$$

(36)

Again from naive dimension estimation, $f_{G,5} \sim f_G \sim 16\pi^2/e\Lambda_3^3$. Unlike the MDM or EDM operators, they always lead to direct detection signals. Concretely, the first operator gives contribution to $\sigma_{SI}^p$ in the form of Eq. (11), with [29]

$$\frac{f_p}{m_p} \sim -\frac{2}{9} f_G f_{TG} \lesssim 10^{-9} \text{ GeV}^{-2}/m_n,$$

(37)

which implies the lower bound $\Lambda_3 \lesssim (16\pi^2 f_G/e)^{1/3} \simeq 3.7$ TeV renders again a very small $(\sigma v)_{2\gamma} \sim 10^{-6}$ pb. Now we can draw the conclusion: in the charged loop ported DM scenario,
if we expect a large annihilation rate to gamma, XENON100 excludes both a Dirac DM and colored charged loops.

We would like to comment on the scalar DM case, the \((S^\dagger \partial_\mu S) \partial_\nu F^{\mu\nu}\) leads to contact interaction between DM and nucleon and thus safe. In fact, such model [11] has been recently proposed to explain the anomaly using the scalar DM.

**B. Decaying scenario**

The scenarios discussed previously involve either some tuning or strong couplings, while decaying DM gives an alternative more natural choice. We close the paper by giving a short comment on this scenario to explain the FERMI gamma-ray line. Here, the scalar DM is more or less favored. At the two-body decay level (three-body or more leads to too wide spectrum to account for the peak), the scalar DM can decay into the gamma pair while the fermionic DM can not. The unique effective operator is written as

\[
\frac{\alpha}{4\pi} \frac{1}{4\Lambda^2} \chi^\dagger \chi F^{\mu\nu} F_{\mu\nu},
\]

but this time the \(\chi\) obtains a TeV scale VEV \(v_\chi\) which breaks the \(Z_2\) symmetry and leads to the scalar DM (the real part of \(\chi\)) decaying to gamma pair. To fit the data, besides a mass of DM should be around 260 GeV, we further need its extremely narrow branching decay width to gamma pair \(\sim 10^{-29} \text{s}^{-1}\) [28]. The small decay width can be achieved by lifting the mass of charged particles running in the loop, typically \(\Lambda_c \rightarrow M_{\text{GUT}}\) (see an example in [33]):

\[
\begin{align*}
\Gamma_{\chi \rightarrow 2\gamma} &= \frac{\alpha^2}{1024\pi^3} \left( \frac{\sqrt{2} v_\chi}{\Lambda} \right)^2 \frac{m_\chi^4}{\Lambda^2} \\
&= 1.1 \times 10^{-29} \left( \frac{v_\chi}{10^3 \text{GeV}} \right)^2 \left( \frac{m_\chi}{260 \text{GeV}} \right)^4 \left( \frac{3 \times 10^{14} \text{GeV}}{\Lambda} \right)^4 \text{s}^{-1}. \quad (39)
\end{align*}
\]

Some comments are in orders. Firstly, in our notation the \(\Lambda\) lies below \(M_{\text{GUT}}\) about two orders, but actually it may be compensated by loop factors and large Yuawa couplings (set to unit in the above estimation). Secondly, the relic density is a generic problem for decaying DM, but non-thermal production such as Ref. [30] using freeze-in mechanism [31] may offer a solution. Last but never the least, the 130 GeV gamma-ray line in decaying DM scenario is easily compatible with the sharp excess in PAMELA positron fraction [16], which can be interpreted by a leptonic decaying DM (saying to \(e^+ e^-\)) with lifetime \(\tau \sim 10^{26} \text{s}\) and mass around 200 GeV. In model building, one has to introduce relevant dimension-six operators for DM decay to leptons and adjust the branching ratio to fit both datas.
For the fermionic DM, the Lorentz invariance force the presence of second fermion (in the final state) in the operator of decay. As a case in point, in the (lepton number violating) $R-$parity violating SUSY, the gravitino has the following two body decay modes [32]:

$$\tilde{G} \rightarrow \nu + \gamma, \quad W^+ + \ell^-, \quad Z^0 + \nu,$$

where gravitino mass is $m_{\tilde{G}} = 250$ GeV and the branching ratios are respectively given by 0.03, 0.69 and 0.28. Although it fails to explain both PAMELA and FERMI gamma-ray line, the first mode can explain the latter given proper decay width. Of particular interest, the neutrino (from the third mode) and gamma signal appear simultaneously, and the accompanied neutrino signal may be detected and thus provides a complementary detect method for this scenario.

V. CONCLUSION AND DISCUSSION

The gamma-ray line from DM annihilating in the galaxy center generically is well below the present detectable level. However, once the observation, from it we are able to extract very important information of the DM properties/dynamics. In this work, we present a minimal effective theory framework to understand the anomalously bright gamma-ray line from dark matter activity:

- In the EFT large annihilation rate is ascribed to operator coefficients with resonant or strong coupling enhancement.

- Due to the XENON100 bound, Dirac DM or colorful charged particles are ruled out in models with only strong couplings.

- If the charged particle in the loop carry $SU(3)_C$ charge, the di-gluon annihilation mode is about one order larger than the di-photon mode, that may properly account for the relic density.

- The SM-like Higgs may share the same charged loop, and therefore provide a source of Higgs di-photon excess at the LHC.

Applying the general analysis to the NMSSM, that is proved to accommodate neutralino LSP with large annihilation rate into di-photon and interpret the tentative 130 GeV gamma-ray line. Top-window model is also proposed to explain it.

Although not the central points of this work, we would like to end up by commenting its very promising collider detection prospect, if the 130 GeV gamma-ray line from DM activity will be confirmed. In light our general analysis in the text, at the LHC or Tavertron (but
beyond LEP) one can expect new light color-singlet charged particle $C$ can be produced: $q\bar{q} \rightarrow C^\dagger C$. While beyond the 130 GeV line and consider more wide scope, the LHC could put very strong exclusion on the model with colored loop.

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**Note added**

In the completion of this work, we note the appearance of work [40], they also note the importance of the di-gluon mode in the determining of DM relic density and more relevant phenomenologies are discussed there. The Ref. [41] specifically studies the 130 GeV gamma-ray line from LSP annihilation in the NMSSM, taking a quite similar idea to ours. We greatly thank the authors for sending us the updated version of NMSSMtools in which the CP-odd mass correctly adopts the running mass. Using it we re-scan our region, and find different results than [41].

**Appendix A: Effective vertex from the charged loop**

![Diagram](image.png)

FIG. 5: fermion loop induced neutral scalar decay
1. $\phi \to 2\gamma$

In this Appendix we present the procedure of calculating the coefficients of the effective vertex used in the Section III A, quoted for convenience:

$$\alpha \frac{\phi_{CC}}{4\pi} \frac{1}{4\Lambda_1} \phi_h F_{\mu\nu} F^{\mu\nu}, \quad \alpha \frac{\phi_{CC}}{4\pi} \frac{1}{8\Lambda_2} \phi_A F_{\mu\nu} \tilde{F}^{\mu\nu}. \tag{A1}$$

The calculation is similar to the case of the Higgs with general couplings. For definiteness, here we focus on the process depicted in the Fig. 5, where CP-even resonance $\phi_h$ decays to gamma pair with single charged fermion running in the loop. At the one hand, the direct calculation of decay width gives

$$\Gamma (h \to \gamma \gamma) = \frac{\alpha^2 m_h^2}{256 \pi^3} \left| 2 N_e Q_C^2 h_{\phi CC} A_{1/2}^h (\tau) \right|^2, \tag{A2}$$

with $N_e = 3$ the color factor and $Q_C$ the electronic charge of $C$. The loop function $A(\tau)$ is

$$A_{1/2}^h (\tau) = \frac{2}{\tau^{3/2}} \left[ \tau + (\tau - 1) f(\tau) \right], \tag{A3}$$

$$f(\tau) = \begin{cases} \arcsin^2 \sqrt{\tau} & \tau \leq 1 \\ -\frac{1}{\tau} \left[ \log \frac{1+\sqrt{1-\tau}}{1-\sqrt{1-\tau}} \right]^2 & \tau > 1 \end{cases} \tag{A4}$$

only depending of the ratio $\tau = m_{\phi_h}^2 / 4m_C^2$. In general, the $\phi_h$ can be off shell and therefore the replacement $m_{\phi_h}^2 \to P^2$, with $P$ the four-momentum of $\phi_h$. For the CP-odd resonance $\phi_A$, the decay width takes the same form as the Eq. (A2) but the loop function is different:

$$A_{1/2}^A = 2\tau^{-1/2} f(\tau). \tag{A5}$$

On the other hand, from the effective operators in Eq. (A1) one can calculate

$$\Gamma (\phi_h \to \gamma \gamma) = \frac{\alpha^2 h_{\phi CC}^2}{1024 \pi^3 \Lambda_1^3} m_{\phi_h}^3, \quad \Gamma (\phi_A \to \gamma \gamma) = \frac{\alpha^2 h_{\phi CC}^2}{1024 \pi^3 \Lambda_2^3} m_{\phi_A}^3. \tag{A6}$$

So, comparing Eq. (A2) and Eq. (A6), we eventually get the effective suppressing scales

$$\Lambda_1 = m_\chi / |2N_e Q_C^2 A_{1/2}^h|, \quad \Lambda_2 = m_\chi / |2N_e Q_C^2 A_{1/2}^A|. \tag{A7}$$

where $m_\phi = 2m_\chi$ has been used. Those expressions justify the statement that the $\Lambda_{1,2}$ can be much below the weak scale (or $m_\chi$).

2. $\phi \to 2G, \quad \phi \to 2Z$

There are two other effective vertex require attention. First, if the charged particle also carries color charge, then the resonance can decay into two gluons. Repeating the procedure
FIG. 6: The loop function $|\mathcal{A}|^2$ varies as variable $\tau = m_\phi^2/4m_C^2$. Dashed line: for the CP-odd $\phi$; Solid line: for CP-even $\phi$. The former is always several times of the latter.

dealing with the vertex $\phi_h FF$ and $\phi_A F \bar{F}$, the corresponding coefficients parameterized exactly as the Eq. (A1) except for the replacement $\alpha \rightarrow \alpha_s$ and $\Lambda_i \rightarrow \Lambda_{i,s}$, we are able to get the precise effective operators. It is straightforward to get the gluonic partial decay width from direct and effective calculations:

$$\Gamma(\phi_h \rightarrow gg) = \frac{\alpha_s^2 m_h}{32\pi^3} |h_{\phi CC} A_{1/2}^h|^2 = \frac{\alpha_s^2 h_{\phi CC}^2}{128\pi^3 \Lambda_1^3} m_h^3,$$

(A8)

$$\Gamma(\phi_A \rightarrow gg) = \frac{\alpha_s^2 m_A}{32\pi^3} |h_{\phi CC} A_{1/2}^A|^2 = \frac{\alpha_s^2 h_{\phi CC}^2}{128\pi^3 \Lambda_2^3} m_A^3.$$

(A9)

Then we obtain the effective scales

$$\Lambda_{1,s} = m_\chi/|A_{1/2}^h|, \quad \Lambda_{2,s} = m_\chi/|A_{1/2}^A|.$$

(A10)

a few times of the values given in the Eq. (A7).

Now we turn attention to the di-$Z$ mode. Compared to the di-phonon mode, the difference lies in the additional terms in the polarization vector substitution: $\epsilon_\mu \epsilon'^*_\nu \rightarrow g_{\mu\nu} - p_1^\mu p_2^\nu/m_Z^2$. Presumably the resulting change is suppressed by the small parameter $m_Z^2/m_C^2$ (confirmed by the $h \rightarrow Z\gamma$ result [39]), then we are justified to ignore this effect at the leading order and approximately have

$$\Gamma(h \rightarrow ZZ) / \Gamma(h \rightarrow \gamma\gamma) \sim g^4 \hat{v}_f^4 \hat{e}^4 Q_f^2 \left( 1 - 4 M_Z^2/m_\phi^2 \right)^{3/2},$$

(A11)

the phase space suppressing factor is about 0.3 for $m_\phi = 260$ GeV, largely it cancels the enhancement from coupling ratio. Generally the vector coupling is of the form $\hat{v}_f \bar{\psi}^\mu \psi Z_\mu$.
with $\hat{v}_f = T^3_C - 2Q_C \sin^2 \theta_W$, here $T^3_C$ is the isospin quantum number of $C$.

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[42] Obviously, our effective discussion is still more fundamental than the Ref. [19] that only give operators consists of DM field and gauge/Higgs fields. And it is can be a generation of [37] that only includes light charged particles.

[43] In the Feynman rule of $\phi FF$ vertex, we have make it explicitly conserves the Ward-Takahashi identity by symmetric the $p_1^\mu p_2^\nu$ term.

[44] We note that alone this line, a pioneer attempt has been made in [24], in which a vector resonance is investigated.

[45] Such a spectator Higgs doublet is motivated in the Ref. [20] to explain the origin of iso-spin violation between the DM-nucleon interaction.