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Two-photon annihilation of singlet cold dark matter due to noncommutative space-time

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Detecting cosmic rays, in particular, gamma rays, coming from dark matter annihilation or decay is an indirect way to survey the nature of dark matter. In commutative space-time, the annihilation of dark matter candidates (weakly interacting massive particles) to photons proceeds through loop corrections. However, it is possible for weakly interacting massive particles as well as other standard model singlet particles to couple with photons directly in noncommutative space-time. In this paper, we study two-photon annihilation of singlet weakly interacting massive particles in noncommutative space-time. If noncommutative interactions are relevant to the relic abundance, one can exclude some dark matter masses by using Fermi-Lat data.

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I. INTRODUCTION

Detection of the annihilation of dark matter into monochromatic gamma rays in upcoming telescopes is certainly an appropriate way to unambiguously determine their unknown nature. The most popular candidates of dark matter are weakly interacting massive particles (WIMPs), which are accommodated in some models beyond the standard model such as supersymmetry models with R parity [1,2], extradimensional models with conserved Kaluza-Klein parity [3], the T-parity conserved little Higgs model [4], and so on. Also singlet particles, either scalars [5] or fermions [6–8], can be served as cold dark matter. In all the above scenarios, the weak interactions of WIMPs are the main key to explain the thermal production of them in the early Universe (for a review, see Ref. [1,2]). Additionally, these weak interactions can provide an opportunity to search dark matter through their production in high energy accelerators [9], their direct detection [10], and their indirect detection, i.e., astrophysical observations of their annihilation or decay products in our Galaxy or beyond. In fact, through the WIMP scenario, the weak interaction of dark matter would produce observable standard model particles, such as charged antimatter particles, photons, and neutrinos. Among these, neutrino and photons have an advantage in comparison to the others, because they keep their source information during streaming. Moreover, the very small cross sections of neutrinos make their flux very difficult to detect. Therefore, the gamma-ray signatures of dark matter have been investigated extensively (for a review, see Ref. [11] and references therein). The continuum gamma-ray emission from dark matter annihilation could be confused with astrophysical backgrounds, e.g., emission from galactic cosmic rays or from millisecond pulsars. Hence, the study of the monochromatic gamma ray is important. Monochromatic gamma-ray signatures have been studied for some dark matter candidates in the literature [12,13].

On the other hand, noncommutative (NC) quantum field theories have been considered in the recent decade extensively because of some motivations coming from string theory [14] and measurement arguments based on quantum mechanics and classical gravity [15]. In the NC field theory, one encounters new properties such as the UV-IR mixing problem [16], the violation of Lorentz invariance.[1] From the phenomenological point of view, by comparing the results of the noncommutative version of usual physical models with the present data, lower bounds on the noncommutative scale have been estimated conservatively at about 1–10 TeV [18]. The NC field theories are constructed on space-time coordinates, which are operators and do not obey commutative algebra. In the case of the canonical version of NC space-time, the coordinates satisfy the following algebra:

$$[\hat{x}^\mu, \hat{x}^\nu] = i\theta^{\mu\nu}, \tag{1}$$

where a hat indicates a NC coordinate and $\theta^{\mu\nu}$ is a real, constant, and antisymmetric matrix. According to the Weyl-Moyal correspondence, to construct the NC field theory, an ordinary function can be used instead of the corresponding NC one by replacing the ordinary product with the star product as follows:

$$f \star g(x, \theta) = f(x, \theta) \exp\left(\frac{i}{2} \overleftarrow{\partial}_\mu \theta^{\mu\nu} \overrightarrow{\partial}_\nu\right) g(x, \theta). \tag{2}$$

Because of the above correspondence, a neutral particle (as well as a charged particle) can couple with the $U(1)$ gauge field in the adjoint representation. Some effects of

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[1] As one can see from (1), NC parameter $\theta^{\mu\nu}$ is a constant antisymmetric matrix which specifies a preferred direction in space-time. However, quantum field theory on NC space-time possesses symmetry under a twisted Poincaré algebra whose representation content is identical to the usual Poincaré symmetry [17].

this new coupling were studied in the literature [19]. In particular, singlet particles, which can be served as cold dark matter, can couple with the $U(1)$ electroweak gauge field in this manner [8]. For instance, this interaction can be relevant to the production of dark matter with masses about 100 GeV provided that the NC scale is about 1 TeV.

In the usual space-time, there does not exist any model which predicts direct coupling between WIMPs and photons, because they are electrically neutral. Hence the annihilation of WIMPs into photons proceeds through loop corrections. However, this process is possible at the tree level for the standard model singlet particles through adjoint representation of $U(1)$ gauge theory in NC space-time. In this paper, we calculate the annihilation cross section of singlet dark matter into two photons in NC space-time. Although this proceeds at the tree level, its contribution is suppressed with θ^4 or equivalently $\frac{1}{\Lambda^8}$, where Λ is the NC scale. However, this NC-induced interaction can be relevant to the thermal production of singlet dark matter in some parameter regions [8]. Therefore, the study of this process helps one to constrain the corresponding parameter regions by using gamma-ray experiments such as Fermi-Lat [20].

This paper is organized as follows: In Sec. II, we give a brief review of the singlet extended noncommutative standard model. In Secs. III and IV, we study the annihilation of a singlet fermion and scalar, respectively, into two photons. Finally, we discuss our conclusions in Sec. V.

II. A BRIEF REVIEW OF SINGLET EXTENDED NONCOMMUTATIVE STANDARD MODEL

The Weyl-Moyal correspondence, Eq. (2), leads to a few restrictions on a gauge theory in NC space-time [21]: (a) Only $U(n)$ gauge theories have a NC extension without any enlargements. Because of some existing terms proportional to the identity matrix due to the Weyl-Moyal correspondence, the usual $SU(n)$ gauge theories in particular are not permissible. (b) Only $n \times n$ matrix representations of $u(n)$ algebra respect the closeness condition. For instance, in the $U_\star(1)$ case for arbitrary fixed charge q, only the matter fields with charges $\pm q$ and zero are permissible. (c) In a gauge theory consisting of several simple gauge groups, the matter fields cannot carry more than two NC gauge group charges. Hence, the extension of the standard model based on $SU(3) \times SU(2) \times U(1)$ gauge theory to NC space-time is problematic. There exist, however, two approaches to construct the standard model gauge theory in NC space-time.

In the first approach, the gauge group is restricted to $U(n)$, and the symmetry group of the standard model is achieved by the reduction of $U(3) \times U(2) \times U(1)$ to $SU(3) \times SU(2) \times U(1)$ by an appropriate symmetry breaking [22]. Namely, two extra $U(1)$ factors are reduced through two extra Higgs particles (rather than the standard model) during appropriate Higgs mechanisms. The number of possible particles in each family is six—left-handed leptons, right-handed charged leptons, left-handed quarks, right-handed up quarks, right-handed down quarks, and Higgs which transform under the standard model gauge group as follows:

$$\Psi_L^l(x) \equiv \begin{pmatrix} \nu(x) \\ e(x) \end{pmatrix}_L \to V(x) \star \Psi_L^l(x) \star \upsilon^{-1}(x), \tag{3}$$

$$e_R(x) \to e_R(x) \star \upsilon^{-1}(x), \tag{4}$$

$$\Psi_L^q(x) \equiv \begin{pmatrix} u(x) \\ d(x) \end{pmatrix}_L \to V(x) \star \Psi_L^q(x) \star U^{-1}(x), \tag{5}$$

$$u_R(x) \to \upsilon(x) \star u_R(x) \star U^{-1}(x), \tag{6}$$

$$d_R(x) \to d_R(x) \star U^{-1}(x), \tag{7}$$

$$H(x) \equiv \begin{pmatrix} H^+(x) \\ H^0(x) \end{pmatrix} \to V(x) \star H(x), \tag{8}$$

where $\upsilon(x)$, $V(x)$, and $U(x)$ are $U(1)$, $U(2)$, and $U(3)$ gauge transformations, respectively. These transformations along with the following transformation for the gauge fields:

$$B_\mu \to \upsilon \star B_\mu \star \upsilon + \frac{i}{g} \upsilon \star \partial_\mu \star \upsilon, \tag{9}$$

$$W_\mu \to U \star W_\mu \star U + \frac{i}{g} U \star \partial_\mu \star U, \tag{10}$$

$$G_\mu \to V \star G_\mu \star V + \frac{i}{g} V \star \partial_\mu \star V, \tag{11}$$

where B_μ, W_μ, and G_μ are $U(1)$, $U(2)$, and $U(3)$ gauge fields, respectively, define the $U(3) \times U(2) \times U(1)$ gauge theory including gauge and Yukawa interactions. Of course, the conservation of the gauge symmetry in the Yukawa interactions for the up quarks leads to the following gauge transformation for the charge conjugate of the doublet Higgs field:

$$H^C \to V(x) \star H^C \star \upsilon^{-1}. \tag{12}$$

As was said, in addition to the standard model contents, there exist two new Higgs to reduce two additional $U(1)$ factors. The details of this model are beyond the scope of this paper (for the details of model building, see Ref. [22]). Moreover, a singlet particle either fermion or scaler with the following gauge transformation:

$$\Phi(x) \to \upsilon(x) \star \Phi(x) \star \upsilon^{-1}(x) \tag{13}$$

can be accommodated. The gauge interactions of these particles as well as Yukawa coupling between singlet fermions and scalars are permissible. However, the couplings between a standard model Higgs and singlet scalar, which is possible in the usual space-time, violate the gauge symmetry in this model [8].

In the second approach, one can construct the $SU(n)$ gauge group in noncommutative space-time by using Seiberg-Witten maps [23]. In this manner, a version of the NC standard model has been constructed which includes only the content of the usual one [24].[2] Explicitly, the Lagrangian of this theory is similar to the commutative standard model, but the fields and products are replaced by the NC fields and star products, respectively. For practical purposes, the NC fields have to be written with respect to the usual fields using Seiberg-Witten maps. Although the interactions of the standard model receive the NC corrections, one encounters some new interactions between the gauge fields themselves or between gauge fields and matter fields which proceed through loop corrections in commutative space-time [24]. In this approach, a singlet particle, either fermion or scalar, can also be transformed under gauge transformation according to (13). Therefore, it can be involved in the gauge interactions through the following minimal coupling:

$$\hat{D}_\mu \hat{\Phi} = \partial_\mu \hat{\Phi} - ig'(\hat{B}_\mu \star \hat{\Phi} - \hat{\Phi} \star \hat{B}_\mu), \tag{14}$$

where hats on the fields are used to emphasize that these fields are defined in NC space-time. It is clear that the Yukawa coupling between the singlet fermion and the singlet scalar is gauge invariant. In addition, the interaction terms between S and the standard model Higgs doublet H, such as $H^\dagger HS$ and $H^\dagger HS^2$, do not violate the gauge symmetry if H transforms under the following representation [8]:

$$H \to V \star H \star v^{-1}. \tag{15}$$

Therefore, a singlet particle beyond the standard model can be coupled with the $U(1)$ gauge field in both versions of the NC standard model. Using Seiberg-Witten maps and Weyl-Moyal correspondence, we expand the relevant action to the annihilation of singlet particles into two photons in terms of the NC parameter θ. To obtain the lowest order of NC corrections, we need the following Seiberg-Witten maps of the singlet particles either fermion or scalar [23]:

$$\begin{aligned}\hat{\Phi} = {}& \Phi + g'\theta^{\mu\nu}B_\nu\partial_\mu\Phi + g'^2\theta^{\mu\nu}\theta^{\kappa\lambda} \\ & \times \bigg[\frac{1}{2}B_\mu B_\kappa \partial_\nu\partial_\lambda\Phi - \partial_\mu B_\kappa B_\nu \partial_\lambda \Phi + \frac{1}{4}B_\nu\partial_\kappa B_\mu \partial_\lambda\Phi \\ & + \frac{1}{8}\partial_\mu B_\kappa \partial_\nu B_\lambda \Phi + \frac{1}{8}\partial_\kappa B_\nu \partial_\mu B_\lambda \Phi\bigg],\end{aligned} \tag{16}$$

and of the $U(1)$ gauge field:

$$\hat{B}_\mu = B_\mu + e\theta^{\nu\rho}B_\rho\bigg[\partial_\nu B_\mu - \frac{1}{2}\partial_\mu B_\rho\bigg]. \tag{17}$$

III. ANNIHILATION OF SINGLET FERMIONS INTO TWO PHOTONS

In commutative space-time, the singlet particles interact with the standard model particles through a renormalizable coupling between the singlet scalar and the standard model Higgs. As was said, the transcription of this coupling in NC space-time leads to the violation of $U(3) \times U(2) \times U(1)$ gauge symmetry. Therefore, singlet particle production is possible to be explained only by using the NC-induced interaction in the $U(3) \times U(2) \times U(1)$ model. Moreover, if one can consider the coupling between the singlet scalar and the standard model Higgs, the NC contribution in the annihilation of singlet particles to photons is comparable to or larger than the commutative one for some region of parameter space. That is why it proceeds at tree level in NC space-time (Fig. 1) contrary to the commutative one, which proceeds through loop quantum corrections. In particular, we are interested in the parameter region where the NC contribution is larger than the commutative one. Hence, we ignore the interference of commutative and NC terms.

The action describing a singlet fermion field in NC space-time is

$$S = \int d^4x (\bar{\hat{\psi}} \star i\gamma^\mu \hat{D}_\mu \hat{\psi} - m\bar{\hat{\psi}} \star \hat{\psi}). \tag{18}$$

After using the mentioned Seiberg-Witten maps and expanding the star product up to the second order of θ, we can write the action as follows:

$$\begin{aligned} S = {}& \int d^4x \bar{\psi}\bigg[(i\gamma^\mu\partial_\mu - m) - \frac{e}{2}\theta^{\nu\rho} \\ & \times (i\gamma^\mu(B_{\nu\rho}\partial_\mu + B_{\mu\nu}\partial_\rho + B_{\rho\mu}\partial_\nu) - mB_{\nu\rho})\bigg]\psi \\ & + ig'^2\theta^{\alpha\beta}\theta^{\kappa\lambda}[\partial_\kappa\bar{\psi}i\partial_\alpha \not{B}\partial_\beta\psi B_\lambda - \partial_\beta\bar{\psi}i\partial_\alpha\not{B}\partial_\kappa\psi B_\lambda \\ & - \partial_\alpha\bar{\psi}i\partial_\kappa\not{B}\partial_\beta\psi B_\lambda + \frac{1}{2}\partial_\alpha\bar{\psi}i\not{\partial}B_\kappa\partial_\beta\psi B_\lambda], \end{aligned} \tag{19}$$

where $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$. The second-order terms with respect to θ are relevant to the third diagram of Fig. 1, which includes only on-shell external lines. Hence, we have not written those terms which are in the order of θ^2 and vanish for the on-shell particles. After electroweak symmetry breaking, B_μ is written in terms of photon and Z_0 as follows:

$$B_\mu = -\sin\theta_W Z_{0\mu} + \cos\theta_W A_\mu, \tag{20}$$

in which θ_W is the electroweak mixing angle. Therefore, the relevant vertices to the fermionic singlet particle annihilation into two photons are given by

[2] We should mention that in the works [23,24], the charge quantization problem inherent in NC gauge field theories discussed and treated in Refs. [21,22] was dismissed by *mapping* three different noncommutative gauge field degrees of freedom to a single ordinary gauge field [25].

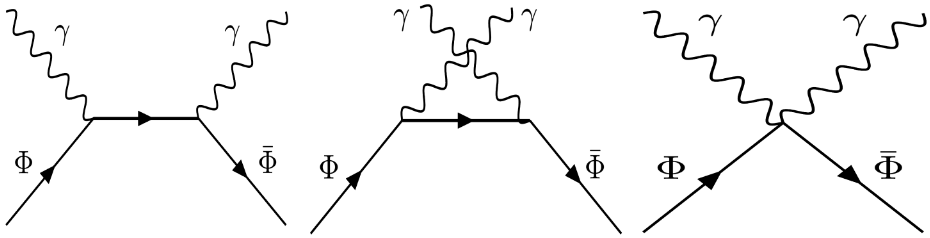

FIG. 1. The tree-level Feynman diagrams for the annihilation of two singlet particles into monochromatic gamma-ray lines in NC space-time. Here, Φ stands for a singlet either fermion or scalar.

$$\Gamma^{\mu}(\psi(k')\bar{\psi}(k)\gamma(q)) = -e[q\theta k\gamma^{\mu} + (\not{k} - m_{\nu})\tilde{q}^{\mu} - \not{q}\tilde{k}^{\mu}] \tag{21}$$

and

$$T^{\mu\nu} = ie^2[k_2\theta p_2\tilde{p}_1^{\mu}\gamma^{\nu} + p_2\theta p_1\tilde{k}_2^{\mu}\gamma^{\nu} - k_2\theta p_1\tilde{p}_2^{\mu}\gamma^{\nu} - k_1\theta p_1\tilde{p}_2^{\nu}\gamma^{\mu} + k_1\theta p_2\tilde{p}_1^{\nu}\gamma^{\mu} + p_2\theta p_1\tilde{k}_1^{\nu}\gamma^{\mu} + \frac{1}{2}p_2\theta p_1(\not{k}_1 - \not{k}_2)\theta^{\mu\nu}], \tag{22}$$

where we have used $e = g'\cos\theta_W$ and the notation $q\theta k = \theta^{\mu\nu}q_{\mu}k_{\nu}$ and $\tilde{q}^{\mu} = \theta^{\mu\nu}q_{\nu}$.

The cross section of the self-annihilation of the singlet fermion into two photons is obtained as follows:

$$\sigma v = 0.15 \times 10^4 \times \frac{s\alpha^2(2.89sm^6 - 7.07m^8 - 0.340s^2m^4 + 0.067m^2s^3 + 0.075s^4)}{m^4\Lambda_{\rm NC}^8}, \tag{23}$$

where s and v are the square center of mass energy and relativistic velocity, respectively. The thermal average cross section times the relativistic velocity can be expanded in terms of $\epsilon = \frac{s-4m^2}{4m^2}$ in the case of nonrelativistic singlet fermions. The first nonzero term of this expansion corresponds to $\epsilon = 0$. Therefore, we have

$$\langle\sigma v\rangle = 1.3 \times 10^5\alpha^2\frac{m^6}{\Lambda_{\rm NC}^8}. \tag{24}$$

One can look at this result from two points of view. First, since the annihilation of WIMPs into photons proceeds through loop corrections in the usual space-time, the cross section of the annihilation into photons is 4 or 5 orders of magnitude weaker than the annihilation cross section requested by the correct relic abundance. Therefore, the comparison of the NC result with the usual expectation in commutative space leads to $\Lambda_{\rm NC}$ being more than 1 TeV. Second, the NC-induced singlet fermion-photon interaction may be relevant to the thermal production of singlet fermion in the early Universe provided that [8]

$$\langle\sigma_{\rm ann}v\rangle_f \sim 3.9 \times 10^{-3}N_f\frac{m^2}{\Lambda_{\rm NC}^4}, \tag{25}$$

where N_f denotes the number of allowed pair charged fermions. Here, $\langle\sigma_{\rm ann}v\rangle_f$ is the thermal average of the annihilation cross section of the singlet fermionic dark matter into standard model massive particles. The correct relic abundance requires $\langle\sigma_{\rm ann}v\rangle_f \sim 1.4\times10^{-26}\,{\rm cm^3\,s^{-1}} \simeq 1.2\times10^{-9}\,{\rm GeV^{-2}}$. Combination of (24) and (25) leads to

$$\langle\sigma v\rangle \sim 1.9 \times 10^{-14}\ {\rm m^2}. \tag{26}$$

Obviously, $\langle\sigma v\rangle$ for dark matter with mass about 100 GeV is about Fermi-Lat bounds. Consequently, singlet fermionic dark matter in NC space-time may be excluded by Fermi-Lat [20] for masses larger than 100 GeV.

IV. ANNIHILATION OF A SINGLET SCALAR INTO TWO PHOTONS

Alternatively, a real singlet scalar can also be served as cold dark matter [5]. The annihilation of this candidate of dark matter into photons has been studied in the usual commutative theory [12]. Therefore, it is also interesting to study the gamma ray coming from the annihilation of the recent candidate of dark matter in NC space-time. The action describing the singlet scalar particles in NC space-time is written as follows:

$$S = \frac{1}{2}\int d^4x((\hat{D}_{\mu}\hat{\phi})^{\dagger}\star\hat{D}^{\mu}\hat{\phi} - m^2\hat{\phi}^{\dagger}\star\hat{\phi}). \tag{27}$$

After replacing the corresponding Seiberg-Witten maps from (16) and (17) and the star products up to the first order of θ, we have

$$S=\int d^4x\Big[\frac{1}{2}\partial_\mu\phi\partial^\mu\phi-\frac{1}{2}m^2\phi\phi+g'\theta^{\alpha\beta}\{\partial_\beta\phi\partial^\mu\phi\partial_\alpha B_\mu-\partial_\alpha\phi(\Box-m^2)\phi B_\beta\}-2g'^2\theta^{\alpha\beta}\theta^{\kappa\lambda}\Big[\partial_\alpha\phi\partial_\mu\partial_\kappa\phi B_\beta\partial^\mu B_\lambda$$
$$+\partial_\alpha\phi\partial_\mu\partial_\lambda\phi B_\beta\partial_\kappa B^\mu+\frac{1}{2}\partial_\beta\phi\partial_\lambda\phi\partial_\alpha B_\mu\partial_\kappa B^\mu+\partial_\alpha\partial_\mu\phi\partial_\beta\phi B_\lambda\Big(\partial_\kappa B^\mu-\frac{1}{2}\partial^\mu B_\kappa\Big)+\partial_\beta\partial_\mu\phi\partial_\kappa\phi\partial_\alpha B^\mu B_\lambda\Big]\Big] \tag{28}$$

for the real scalar field. Therefore, the vertex of the coupling between the singlet scalar and photon at the first order of θ is given by

$$\Gamma^\mu(\phi(k_1)\phi(k_2)\gamma(q))=-e\theta^{\alpha\beta}[k_{1\beta}k_{2\alpha}k_2^\mu+k_{2\beta}k_{1\alpha}k_1^\mu-k_{1\alpha}g_\beta^\mu(k_2^2-m^2)-k_{2\alpha}g_\beta^\mu(k_1^2-m^2)], \tag{29}$$

where we use $e=g'\cos\theta_W$. At the second order of θ we have

$$T^{\mu\nu}=2ie^2\Big[2k_1\theta p_1\tilde{p}_2^\nu p_1^\mu+2k_2\theta p_2\tilde{p}_1^\mu p_2^\nu+2k_1\theta p_2\tilde{p}_1^\nu p_2^\mu+2k_2\theta p_1\tilde{p}_2^\mu p_1^\nu-k_1\theta p_1k_2\theta p_2g^{\mu\nu}-k_2\theta p_1k_1\theta p_2g^{\mu\nu}$$
$$-\tilde{p}_1^\mu\tilde{p}_2^\nu(p_1.k_1+p_2.k_2)-\tilde{p}_1^\nu\tilde{p}_2^\mu(p_2.k_1+p_1.k_2)+p_1\theta p_2\tilde{k}_2^\mu(p_1^\nu-p_2^\nu)+p_1\theta p_2\tilde{k}_1^\nu(p_1^\mu-p_2^\mu)$$
$$+\frac{1}{2}p_1\theta p_2(k_1.p_1-k_2.p_1)\theta^{\mu\nu}+\frac{1}{2}p_1\theta p_2(k_2.p_2-k_1.p_2)\theta^{\mu\nu}\Big], \tag{30}$$

where the scalars are considered to be on the mass shell. Therefore, after a straightforward calculation, we obtain the cross section of the annihilation of singlet scalars into two photons as follows:

$$\sigma v=2\alpha^2\frac{71m^4s^3-159m^6s^2+150m^8s+291m^{10}-7.49m^2s^4+2.22s^5}{m^4\Lambda_{\rm NC}^8}. \tag{31}$$

Similar to the fermionic case, we consider nonrelativistic singlet scalars and expand the thermal average of the annihilation cross section times velocity in terms of $\epsilon=\frac{s-4m^2}{4m^2}$. Putting $s=0$, we obtain the first nonzero in the nonrelativistic limit as follows:

$$\langle\sigma v\rangle=6.5\times10^3\alpha^2\frac{m^6}{\Lambda_{\rm NC}^8}. \tag{32}$$

We also look at this result from two points of view. First, taking $m=100$ GeV and $\Lambda_{\rm NC}=1$ TeV leads to $\langle\sigma v\rangle\sim3.5\times10^{-13}$ (GeV)$^{-2}$, which is comparable with the expectations for the corresponding value in the non-resonance region in commutative space-time [12]. Second, the NC-induced interactions can be relevant to the thermal explanation of relic abundance provided that

$$\langle\sigma_{\rm ann}v\rangle_s\sim1.9\times10^{-2}N_f\frac{m^2}{\Lambda_{\rm NC}^4}, \tag{33}$$

where N_f denotes the number of allowed pair charged fermions. Here, $\langle\sigma_{\rm ann}v\rangle_s$ is the thermal average of the annihilation cross section of singlet scalar dark matter into standard model massive particles. After replacing the value of $\langle\sigma_{\rm ann}v\rangle_s$ and combining with (32), we obtain

$$\langle\sigma v\rangle\sim3.7\times10^{-17}\ {\rm m}^2. \tag{34}$$

Consequently, in this case, $\langle\sigma v\rangle$ is comparable to or larger than Fermi-Lat bounds [20] for dark matter with mass larger than about 3 TeV.

V. SUMMARY AND DISCUSSION

In this paper, we have considered an extension of the standard model which includes a singlet either fermion or scalar as cold dark matter in NC space-time. NC space-time induces a new coupling between singlet particles with photons through the adjoint representation. This coupling may be relevant to the thermal explanation of dark matter production in the early Universe provided that the NC scale is about 1 TeV, which is consistent with the phenomenologically obtained bounds [8]. In this paper, we have calculated the cross section of the annihilation of a singlet, either fermion or scalar, into two photons in NC space-time. We have found that if the NC scale is such that the NC-induced interactions are relevant to the production of dark matter, the recent Fermi-Lat results [20] will exclude masses larger than 100 GeV for fermionic dark matter and 3 TeV for scalar dark matter.

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