Secluded WIMPs, QED with massive photons, and the galactic center gamma-ray excess

E. C. F. S. Fortes1,2,3, V. Pleitez1, F. W. Stecker2,4
1 Instituto de Física Teórica, Universidade Estadual Paulista,
Rua Dr. Bento Teobaldo Ferraz 271, 01140-070 São Paulo, SP, Brazil
2 NASA Goddard Space Flight Center, Greenbelt, MD, 20770, USA
3 University of Maryland, College Park, MD, 20742, USA and
4 University of California, Los Angeles, CA 90095, USA

We propose a secluded WIMP dark matter model consisting of neutral fermions as the dark matter candidate and a Proca-Wentzel (PW) field as a mediator. In the model that we consider here, dark matter WIMPs interact with standard model (SM) particles only through the PW field of ~ MeV – multi-GeV mass particles. The interactions occur via a U(1)′ mediator, V ′ µν, which couples to the SM by kinetic mixing with U(1) hypercharge bosons, B, B′. One import difference between our model and other such models in the literature is the absence of an extra singlet scalar, so that the parameter with dimension of mass M2 V ′ is not related to a spontaneous symmetry breaking. This QED based model is also renormalizable. The mass scale of the mediator and the absence of the singlet scalar can lead to interesting astrophysical signatures. The dominant annihilation channels are different from those usually considered in previous work. We show that the GeV-energy γ-ray excess in the galactic center region, as derived from Fermi-LAT Gamma-ray Space Telescope data, can be attributed to such secluded dark matter WIMPs, given parameters of the model that are consistent with both the cosmological dark matter density and the upper limits on WIMP spin-independent elastic scattering.

PACS numbers: 12.60.Cn, 95.35.+d

I. INTRODUCTION

The clear astronomical [1] and cosmological evidence for large amounts of dark matter (DM) in the universe [2] has led to the construction of various theoretical models that go beyond the standard model (SM) weak-scale theories and which attempt to account for the DM abundance in the universe [3]. The observational evidence for DM has motivated various experimental searches to find dark matter [4].

Data from colliders are used to search for evidence of dark matter particles. Experiments with detectors like DAMA, CoGeNT, CDMS, XENON and LUX are used to search for evidence of the recoil energy of nuclei that would be produced by scattering with dark matter particles [5, 6]. High-energy colliders like LHC (Large Hadron Collider), have obtained significant upper limits on the annihilation of WIMPs to quarks [7]. They also offer very interesting possibilities to investigate interactions involving DM mediators.

Space-borne detectors have been used to search for evidence of the products of of dark matter annihilation, particularly γ-rays and cosmic-ray positrons. These searches have conservatively produced constraints on dark matter annihilation, both from cosmic γ-ray studies [8] and cosmic-ray positron studies [9]. However, analyses of the Fermi-LAT data have indicated the existence of an "excess" flux of γ-rays above that expected from cosmic rays interacting with interstellar gas. This flux appears to be extended around the region of the galactic center. It appears to peak in the 2 – 3 GeV energy range. This "excess" has been interpreted to be a possible indication of the annihilation of weakly interacting dark matter (WIMP) particles having a mass in the 20 – 45 GeV range, annihilating primarily into quarks or, less likely, in the 7 – 12 GeV mass range annihilating primarily into charged τ leptons [10, 11]. We note, however, the determination of various possible components of γ-ray emission from the galactic center is complicated and other possible interpretations of this excess have been suggested [12].

Neutralino supersymmetric WIMPs, viz., the lightest supersymmetric dark matter particles, have been a popular choice to be the DM WIMPs because they are stable and neutral and their cross section naturally leads to the correct cosmological DM density. However, as of now, the LHC has not found any evidence for such particles. Therefore, other candidate WIMP models have been explored and should be further explored.

In this work we consider a model of WIMP dark matter in which the dark sector is just quantum electrodynamics (QED) extended with a new massive photon field, usually dubbed a Proca-Wentzel (PW) field. It is well known that this is a renormalizable theory because it couples to a conserved vector current. Hence, the dark sector is made up of Dirac fermions, η, that only interact via a PW field, here denoted by V ′ µν, that serves to mediate between DM fermions and standard model particles.

Diagonalization of the kinetic terms in the Lagrangian give both the standard Z boson and an extra neutral gauge boson that we denote as Z′. We will show that this DM model can produce the observed cosmological DM abundance and that its annihilation into standard model (SM) particles can also lead to astronomically observable fluxes of γ-rays in the galactic center and in dwarf galaxies.

In our model secluded DM interactions occur only through Z′ mediators which subsequently decay to SM particles. Thus, they have an very small elastic scattering cross section with nuclei. This distinguishes such secluded WIMP models from other WIMP models. The decay of Z′’s into quark-antiquark and τ−τ+ channels produces pions, among which are π0's that decay to produce γ-rays. The π0-decay γ-ray spectrum has a characteristic peak at mπ0/2 [13] and is bounded by the rest-mass of the WIMP, mπ (e.g., [14]). The Z′ decay, particularly into charged leptons and light quarks, also yields γ-rays.
through internal bremsstrahlung. In this process, a γ-ray spectrum is produced that peaks near $m_\gamma$ [15]. Electrons resulting from this process can produce γ-rays via Compton scattering in the interstellar medium.

Other Z' models have recently been discussed in the context of astrophysical γ-ray production (see, e.g., Ref. [16]). However, our model, in which the Z' is the mediator of the DM interactions through the annihilation channel $\eta \rightarrow Z' Z'$, as shown in Figure 1, and with $m_{Z'} \ll m_\eta$, was not considered in Ref. [16]. (See also the discussion in Section [11].)

The outline of the paper is as follows: In section [11] we present the model. In section [111] we discuss the differences between our model and previously explored WIMP DM models. In section [1111] we calculate the relic density of our Dirac fermion DM candidate. In section [11111] we discuss the physics production of a fermion DM candidate. In section [111111] we present the model. In section [1111111] we discuss the differences between our model and previously explored WIMP DM models. In section [11111111] we calculate the relic density of our Dirac fermion DM candidate. In section [111111111] we discuss the physics production of a fermion DM candidate.
where $s_W$ is the usual weak mixing angle, $\xi = g_{Vb}/\sqrt{1 - g_{Vb}^2}$, $r = M_1^2/M_2^2$, and where the matrix $U$ is given by

$$
U = \begin{pmatrix}
    c_W & s_W & 0 \\
    -s_W & c_W & 0 \\
    0 & 0 & 1
\end{pmatrix}
$$

(8)

We can write the $3 \times 3$ mass eigenstates matrix as

$$
\begin{pmatrix}
    B \\
    W_3 \\
    V
\end{pmatrix} = \begin{pmatrix}
    c_W & -s_W c_{\alpha} & s_W s_{\alpha} \\
    s_W & c_W c_{\alpha} & -c_W s_{\alpha} \\
    0 & s_{\alpha} & c_{\alpha}
\end{pmatrix} \begin{pmatrix}
    A \\
    Z \\
    Z'
\end{pmatrix}
$$

(9)

with

$$
t_{2\alpha} = -\frac{2 s_W \xi}{1 - s_W^2 \xi^2 - r},
$$

(10)

where $t_{2\alpha}$ denotes $\tan 2\theta_{\alpha}$, $\alpha$ denotes $\theta_{\alpha}$, $c_{\alpha}$ denotes $\cos \theta_{\alpha}$ and $s_{\alpha}$ denotes the $\sin \theta_{\alpha}$. For small values of $\theta_{\alpha}$, we can expand $t_{\alpha} = t_{2\alpha}/2 - t_{4\alpha}/8$ and use the relation $c_{\alpha} = (1 + t_{\alpha}^2)^{-1/2}$ to perform the calculations.

The masses of the $Z$ and $Z'$ are found by diagonalizing the matrix [7]. They are given by:

$$
M_{Z,Z'} = \frac{M_Z'}{2}[(1 + s_W^2 \xi^2 + r) \pm \sqrt{(1 - s_W^2 \xi^2 - r)^2 + 4 s_W^2 \xi^2}],
$$

(11)

where $M_Z = g v_b/(2c_W)$, assuming that $\xi \ll 1$ and $r < (1 - s_W^2 \xi^2)$. (Again, we assume that the lighter neutral vector boson is $Z$.)

### III. MAIN DIFFERENCES COMPARED TO OTHER HEAVY PHOTON MODELS

Electroweak models with an extra $U(1)$ symmetry are among most well motivated extensions of the Standard Model. In a general context we use the notation $U(1)_1 \otimes U(1)_2$. All these models allow a kinetic mixing between the field strength tensors of both $U(1)$ gauge bosons, $F_{1\mu}, F_{2\mu}$.

We can separate this sort of models in several groups. For instance,

**A1** Both $U(1)$ groups are visible. It means that SM fermions carry both quantum numbers. This possibility usually arises in the context of grand unified theories (GUTs). Examples are the models of Babu et al. and that of Galison and Manohar [17]. Models in which $U(1)_1 = U(1)_2 = U(1)_{B-L}$ are of this type.

**A2** There are fermions which carry only one of the $U(1)$ charges, others carry both charges and some with no charges. As an example of these sort of models we have del Aguila et al. in [17].

**A3** One of the $U(1)$ factor is visible, the SM particles carry one of the $U(1)$ charge, for example $U(1)_1$, an the other, $U(1)_2$, is dark. Our model is of this type.

Another way to classify these sort of models is by considering the mass scale at which the kinetic mixing occurs. Although it is not usually explicitly say, this is an important point.

**B1** Models in which $U(1)_1 = U(1)_2$. The kinetic mixing occurs before the SSB. Our model and the one of Ref. [21] are of this type.

**B2** Models in which $U(1)_2 = U(1)_0$, i.e., the kinetic mixing occurs after the SSB and the mixing is with the massless photon. An example of this type of model is that of Holdom in [17].

**B3** Neither $U(1)_1$ nor $U(1)_2$ are related to the symmetry of the SM directly, but $U(1)_1 \otimes U(1)_2 \rightarrow U(1)_Y$ after SSB. The Galison and Manohar model [17] is of this type.

There are also models in which

**C1** The mixing among the two $U(1)$ factor occur in the kinetic and also in the mass terms. In this case the photon has a component on $Z'$. For small values of $\theta_{\alpha}$, we can expand $t_{\alpha} = t_{2\alpha}/2 - t_{4\alpha}/8$ and use the relation $c_{\alpha} = (1 + t_{\alpha}^2)^{-1/2}$ to perform the calculations.

**C2** The mixing among the two $U(1)$ factors occurs only in the kinetic term. Examples of these sort of model are our model and that of Ref. [21]. In these models the photon has no component on $Z'$.

We can also classify the models according to

**D1** Extra scalars, for instance singlets, are added to break the additional $U(1)$ symmetry.

**D2** The only scalar in the model is that of the SM. Our model is of this type.

Models may also be classified according to the fermion, scalar or vector nature of both the DM and the mediator.

**E1** Our model has a Dirac fermion as DM and a real massive vector as mediator. The model of Ref. [23] postulates a real field scalar to be the dark force and a complex scalar to be the dark matter.

Our model is based on the following considerations. It is well known that massive quantum electrodynamic is a renormalizable and free of anomalies theory. It means that spontaneous symmetry breaking in order to give mass to the vector field can be implemented, but it is not mandatory.

The important point is that the photon (massless or not) couples to a conserved current. This occurs because the high energy behavior of the vector propagator $\sim k^2/V^2$, vanishes when contracted with the conserved current, say $J_\mu$. The $U(1)'$ symmetry in the kinetics term is broken by the mass term $(m_{Z'}^2/2)V_{\mu}V^\mu$, but its only consequence is the constraint $\partial_\mu V^\mu = 0$, which comes from the equation of motion. It is still possible to restore the gauge invariance if a massive scalar field called Stueckelberg field is invoked. If $S(x)$ denotes the gauge Stueckelberg scalar, satisfying the equation $(\Box - m_S^2)S(x) = 0$, the Proca-Wentzel Lagrangian is, in fact,
equivalent to the Stueckelberg Lagrangian when \( S(x) = 0 \) (For the main references, see [23]). Usually one of the Abelian group is related to \( U(1)_{B-L} \) [24].

In our model we have a Dirac fermion and a vector mediator and the mediator is lighter than the DM candidate. This is also the case in Ref. [25]. However, in the latter model, the leading interaction between the Standard Model and the dark sector is the kinetic mixing between the photon and the \( U(1) \) gauge boson: \( \mathcal{L} = (1/2)\epsilon F_{\mu\nu}^\gamma F^{\mu\nu}\). The leading kinetic mixing with the photon is also used in Ref. [26]. In the present model, the kinetic mixing is with the gauge boson of the \( U(1)_Y \) factor before the spontaneous symmetry breaking induced by the Higgs boson in the SM that we have denoted \( B' \). Is in this condition that the mixing in Eq. (5) occurs. Hence, the leading interaction between the SM particles and the dark sector is through the \( W^3 \) and \( B \) component on \( Z' \). See equation (9). It is through the mixing with the \( Z \) boson that the particles in the SM have the couplings with \( Z' \) given in the Appendix.

A model in which the photon has no dark component is that in Ref. [21], however, they introduce an scalar singlet \( \Phi_H \), which we denote by \( v \), in which the Higgs boson in the SM is denoted \( v_H \). The scalar singlet makes the difference of the model in Ref. [21] and ours. In former case the interactions the of Higgs bosons are:

\[
\begin{align*}
hff : & -i c_h \frac{m_f}{v_h} & \\
hWW : & 2 i c_h \frac{m_W^2}{v_h} & \\
hZZ : & 2 i c_h \frac{m_h^2 (c_\alpha - \xi s_W c_\alpha)^2}{v_h} - i s_h \frac{m_h^2}{v_h} & \\
hZ'Z' : & 2 i c_h \frac{m_h^2 (s_\alpha + \xi s_W c_\alpha)^2}{v_h} - i s_h \frac{m_h^2}{v_h} & \\
hZZ' : & 2 i c_h \frac{m_h^2}{v_h} [(c_\alpha - s_W \xi s_\alpha)(s_\alpha + s_W \xi c_\alpha)] - i s_h \frac{m_h^2}{v_h} & \end{align*}
\]

where \( v_h \) is the VEV of the scalar singlet (designated by \( \xi \) in Ref. [21]) that is needed to break the \( U(1)_Y \) symmetry and \( c_h = \cos \theta_h, s_h = \sin \theta_h \) are the cosine and sine of the mixing angle:

\[
\frac{\phi_{SM}}{\phi_H} = \left( \begin{array}{cc}
    c_h & s_h \\
    -s_h & c_h
  \end{array} \right)
\]

with \( h \) and \( H \) being the mass eigenstates. The scalar singlet carries \( U(1)_Y \) charge denoted \( e' \) in [18].

Moreover, the scalar potential has several dimensionless parameters in the quartic interaction terms:

\[-\mathcal{V}^{(4)}(\Phi, S) = \lambda (\Phi_{SM}^4 + \Phi_{SM} S^3) + \rho (S^4 S^2) + \kappa S^3 \Phi_{SM} \Phi_{SM}, \]

which obeys the relation [21]:

\[
\tan(2\theta_h) = \frac{\kappa v_h v_S}{\sqrt{\rho v_h^2 - \lambda v_h^2}}.
\]

In a model with a scalar singlet, the SM-like Higgs scalar has weaker interactions with fermion and vector bosons. Such interactions are suppressed by a factor of \( c_h \) relative to SM interactions. The vertices \( hZZ, hZZ' \) and \( hZ'Z' \) have an extra term proportional to \( s_h \). Only when \( c_h = 1 \), do the interactions with the SM particles for both models agree. However, in the models of Refs. [18] and [21], the dark vector still interacts with the singlet \( h \). In general, the phenomenology of both models is a bit different in processes like \( h \to Z'Z', Z'Z' \). This may imply that the region allowed for the four parameters \( m_h, g_{y'V_B} \) and \( m_V \) is different in both models, although an overlap may also exist, as follows from the DM calculation above. For example, in the models of Refs. [18] and [21], the mass of the vector field \( m_V = g_{y'VS} \) and the mass of \( h \) also depends on \( v_S \). Hence, constraints on \( m_V \) are also constraint on \( m_h \) which could be in disagreement with accelerator data. This relation does not exist in our model in which there is no scalar singlet. In our model \( m_V \) is a free parameter. These interactions are equal to those in the appendix only when \( c_h = 1 \) in equation (12).

Notice that we are considering the \( Z' \) to be lighter than the \( Z \), then in our Eq. (11) the signal \( e' \) corresponds to the lighter boson. In fact, it may have zero mass if \( M_V = 0 \). In this case we have a dark QED indeed. We also note that the values that we take for \( g_{y'V_B} \) (called \( \epsilon \) in experimental papers) are much lower than the present experimental upper limits [27] but may be within the sensitivity range of future experiments [28].

**IV. COSMOLOGICAL DM DENSITY**

As noted before, in this model, in addition to SM particles, there is a dark sector composed of a Dirac fermion, \( \eta \), and a Proca-Wentzel vector field, \( V_\mu \approx Z' \), in the limit of small mixing angle. We consider \( \eta \) to be the only component of DM. In the regime where \( m_\eta \gg M_Z \), the primary annihilation channel will be \( \eta + \eta \to Z' + Z' \) [18], as shown in Figure 1. The differential annihilation cross section in this case is given by:

\[
\frac{d\sigma_{ann}}{d\Omega} = \frac{c_\eta^4 g_\eta^4}{64\pi s (1 - g_{y'V_B}^2)} \sqrt{s - 4m_\eta^2} \times [A_1 + \cos \theta \xi A_2 + A_2 \cos \theta (4m_\eta^2 + s + A_2 \sin \theta)],
\]

where \( d' \equiv d\Omega \sin \theta d\theta \), \( A_1 = s (32m_\eta^2 - 8m_\eta^2 s - s^2) \) and \( A_2 = \sqrt{s (s - 4m_\eta^2)} \), where the summation over final spins and
average over initial spins are taken into account.

After integrating equation (14) to obtain the total annihilation cross section, we used the approximation for the square of the center-of-momentum energy, \( s \approx 4m_\eta^2 + m_\chi^2 \), which is valid for non-relativistic particles. The annihilation cross-section can then be expressed in the form \( \langle \sigma_{\text{ann}}v \rangle = a + bv^2 \), where the \( a \) and \( b \) are given in equation (15). Note that relation between the velocity in the center-of-momentum \( (v_{\text{cms}}) \) and the relative velocity \( (v) \) is given by \( v_{\text{cms}} = v/2 \) [29]. We then find

\[
\langle \sigma_{\text{ann}}v \rangle \approx \frac{c^4 a^4}{32\pi(1 - g_{\text{w}}^2)m_\eta^2} + \frac{3}{8} \frac{c^4 b^4}{32\pi(1 - g_{\text{w}}^2)m_\eta^2}v^2.
\]

Using this expression, we calculated relic density as a numerical solution to the Boltzmann equation, discussed in [30, 31].

\[
(\Omega h^2)^2 \approx \frac{1.04 \times 10^9 \text{ GeV}^{-1}}{M_{\text{pl}}} \frac{X_f}{\sqrt{g_*(X_f)}} \frac{1}{a + 3b/X_f},
\]

where \( X_f \) is given by

\[
X_f = \ln \left[ c(c + 2) \frac{45 \theta_m M_{\text{pl}}(a + 6b/X_f)}{8 \pi^2 \sqrt{g_*(X_f)}} \right].
\]

and where \( g = 2 \) for fermionic DM, the Planck max, \( M_{\text{pl}} = 1.22 \times 10^{19} \text{ GeV} \), \( c \) is a parameter of order unity considered here as \( 5/4 \) and \( g_*(X_f) \) is the number of relativistic degrees of freedom at freeze-out. Given one neutral Dirac massive particle and one additional neutral gauge boson and all the SM content, \( g_*(X_f) \approx 113.25 \), \( X_f = m_\eta/T_f \) and \( T_f \) is the temperature at freeze-out. For relics with mass in the range of electroweak scale, \( X_f \) is in the range 20 – 30.

Another expression for \( \langle \sigma_{\text{ann}}v \rangle \), given in Ref. [18], is often used to give the annihilation cross section for fermionic dark matter. In the case where the mass of the mediator is negligible, it’s given by

\[
\langle \sigma_{\text{ann}}v \rangle = \frac{1}{2} \frac{\pi (\alpha')^2}{m_\eta^2}.
\]

The factor 1/2 in equation (18) comes from averaging over the \( U(1)' \) charges. Here, \( \alpha' \) is strength of the interaction in the dark sector. The value for \( \alpha' \) that gives the correct DM abundance is \( \sim 7.8 \times 10^{-4} \).

Given the values of our \( a \) and \( b \) coefficients, and using the standard approximation \( \langle \sigma_{\text{ann}}v \rangle = a + 6b/X_f \), we find from evaluating equations (15) and (18) that \( \alpha' \approx 8 \times 10^{-4} \), which is in good agreement with the result obtained in Ref. [18] and which yields the observed DM abundance.

Figures 2 – 5 give the results for \( \langle \sigma_{\text{ann}}v \rangle \) and \( \Omega h^2 \) as a function of \( m_\eta \) in the limit where \( M_{\chi'}/m_\eta \rightarrow 0 \), taking the values for the model parameters as indicated in the figure captions.

V. GAMMA-RAYS FROM \( \eta \) ANNIHILATION

We have pointed out that our particular secluded WIMP dark matter model has certain attractive theoretical features.
hilation models can fit the galactic distribution of the $\gamma$-ray excess, these authors assume a generalized Navarro-Frenk-White profile:

$$\rho(r) = \frac{\rho_0}{[(r/R_s)^\gamma + 1]^{\gamma+1}},$$  \hspace{1cm} (19)$$

where $\gamma$ is the profile's inner slope and $R_s$ is the scale radius and where they have normalized the DM profile by taking $\rho_{\text{localDM}} = 0.4$ GeV cm$^{-3}$.

In Ref. [11] it is concluded that the observed $\gamma$-ray spectrum is best fit by WIMPS with a mass $\sim 20 - 50$ GeV that annihilate to quarks with a cross section $\langle \sigma v \rangle = 0(10^{-26})$ cm$^3$s$^{-1}$. Motivated by that work, we consider here two specific WIMP mass models, viz., 20 GeV and 32 GeV. With the parameter choices for these models as given and discussed in Section IV, our results are compatible with the observed DM density as parametrized by the value $\Omega h^2 = 0.119$, and also with the value for $\langle \sigma v \rangle$ required to explain the proposed $\gamma$-ray excess from DM annihilation.

**Model 1:** In this first example, we take $m_\eta = 20$ GeV, with the parameters as shown in the caption of Figures 2 and 3 chosen to fit the DM value for $\Omega h^2 \approx 0.119$ and the value $\langle \sigma v \rangle \approx 1.8 \times 10^{-26}$ cm$^3$/s. The lifetime of the light mediator, $\mathcal{Z}$', must be less than one second in order to guarantee that the $\mathcal{Z}'$ decays before the epoch of big bang nucleosynthesis [20]. We have chosen $M_{\mathcal{Z}'} = 0.5$ GeV, which implies that $M_{\mathcal{Z}} = 353$ MeV. Given this $M_{\mathcal{Z}'}$ mass and all the couplings of $\mathcal{Z}'$ with fermions (see appendix), we then find the $\mathcal{Z}'$ decay width, $\Gamma_{\mathcal{Z}'} \approx 1.76 \times 10^{19}$ GeV. The branching ratios for the $\mathcal{Z}'$ decay channels are found to be 33.9% into $u\bar{u}$, 25.4% into $e^+e^-$, 24% in $\mu^+\mu^-$, 8.48% into $d\bar{d}$ and 8.18% into $s\bar{s}$.

In this case, the annihilation $\gamma$-rays primarily result from the decay of the $\mathcal{Z}'$ into light quarks $u\bar{u}$, $d\bar{d}$, $s\bar{s}$, which accounts for 50.6% of the decay width. It is shown in Ref. [11] that the annihilations of DM particles of mass between ~ 18 and ~ 26 GeV into light quarks may significantly account for the ~ 2 - 3 GeV excess.

Choosing the set of parameters given by $g_\eta = 0.088$, $m_\eta = 20$ GeV, $M_{\mathcal{Z}'}/m_\eta \to 0$, $g_{VB} = 8 \times 10^{-3}$, and taking into account the constraint for tan $\theta_9$, designated by $\Gamma_9$ in equation (10), we obtain $c_9 \approx 1$. In this case, the spin independent cross section of $\eta$ with the proton and neutron is $\sigma_9^p = 2.31 \times 10^{-46}$ cm$^2$ and $\sigma_9^n = 5.42 \times 10^{-47}$ cm$^2$. The value of $\sigma_9^f$ is of the same order of magnitude as that of the upper limits obtained by the both the LUX (Large Underground Xenon) experiment and the XENON100 experiment for the WIMP mass range that we consider here [9]. In our calculations, we obtained our parameters using micromegas package [33] [34]. We note that $\sigma_9^{f_{\mathcal{Z}'}}$ and $\sigma_9^{f_{\mathcal{Z}I}}$ are dependent on $g_{VB}$ parameter, which in practice connects the dark and visible sectors of the model (cf., Model 2 below).

**Model 2:** In this model, we take $m_\eta = 32$ GeV. We choose the interaction strength $g_\eta = 1.12 \times 10^{-7}$, we take $g_{VB} = 1 \times 10^{-6}$, and we increase the mass of the $\mathcal{Z}'$ mediator to 3.5 GeV. The dominant interaction channel is still $\eta + \eta \to \mathcal{Z}' + \mathcal{Z}'$. In this case, we find $\langle \sigma v \rangle = 1.88 \times 10^{-26}$ cm$^3$s$^{-1}$ and we fit the DM $\Omega h^2 \approx 0.119$ from the results shown in Figures 4 and 5. Our values for the spin independent cross section of $\eta$ with the proton and neutron are $\sigma_9^p = 6.04 \times 10^{-46}$ and $\sigma_9^n = 1.42 \times 10^{-52}$ cm$^2$. For $M_{\mathcal{Z}'} = 3.5$ GeV, the additional $c\bar{c}$ decay channel opens up, being now kinematically allowed. The $\mathcal{Z}'$ decay width, $\Gamma_{\mathcal{Z}'}$, will therefore be larger than that for Model 1, in this case $\Gamma_{\mathcal{Z}'} \approx 3.74 \times 10^{14}$ GeV. The branching ratios for the $\mathcal{Z}'$ decay are found to be: 25% in $u\bar{u}$, 24.9% in $c\bar{c}$, 18.8% into $e^+e^-$, 18.8% into $\mu^+\mu^-$, 6.28% into $d\bar{d}$, 6.28% into $s\bar{s}$. Thus, in addition to the other channels of Model 1, the $c\bar{c}$ now also significantly contributes to $\eta^0 \to \gamma\gamma$ production. For a large, $c\bar{c}$ decay channel it is shown in Ref. [11] that the annihilations of DM particles of mass between ~ 28 GeV and ~ 36 GeV into $c\bar{c}$ channels may significantly account for the ~ 2 - 3 GeV excess.

The expression for $\mathcal{Z}'$ decay width into fermion-antifermion pairs, $f\bar{f}$, is given by

$$\Gamma_{\mathcal{Z}'} \to f\bar{f} = \sum_i \frac{N_i g_i^2}{c_{\mathcal{F}i}} \frac{1}{12\pi m_{\mathcal{Z}'}^2} \sqrt{m_{\mathcal{Z}'}^2 - 4m_i^2} \times \left[ f_A^i \frac{1}{\mathcal{M}_{\mathcal{Z}'}^2 - 4m_i^2} + f_B^i \right] \left( M_{\mathcal{Z}'}^2 + 2m_i^2 \right),$$  \hspace{1cm} (20)$$

where the sum runs over quarks and leptons species kinematically allowed. $N_i$ is the color number, $m_i$ denotes the mass of the fermion and $f_A^i, f_B^i$ denotes axial/vectorial couplings of $\mathcal{Z}'$ to fermions. Therefore, in the model considered here, the $\mathcal{Z}'$ will decay more into up-type quarks than down-type quarks. Even if we increase $M_{\mathcal{Z}'}$ so that $M_{\mathcal{Z}'} > 2m_{\eta}$, the $bb$ channel will still not dominate. We could open this channel kinematically, but this would require a larger mass for $\eta$ in order to keep the primary annihilation channel as to be $\eta + \eta \to \mathcal{Z}' + \mathcal{Z}'$ in the mass range $m_\eta \gg M_{\mathcal{Z}'}$.

If, for example, we take the same set of parameters for Model 2, only increasing $M_{\mathcal{Z}'} \sim 9.28$ GeV, so that $\mathcal{Z}' \to b\bar{b}$ is now opened, than the branching ratios of $\mathcal{Z}'$ would be 20.1% in $u\bar{u}$, 20.1% in $c\bar{c}$, 15% into $e^+e^-$, 15% into $\mu^+\mu^-$, 14.9% in $\tau^+\tau^-$, 5.15% into $d\bar{d}$, 5.15% into $s\bar{s}$ and only 4.57% into $b\bar{b}$. One of the main phenomenological differences between this model and the other ones presented in the literature is that, as
can be seen in this example, \( Z' \to b\bar{b} \) will not be the most probable decay channel. The detailed expression for \( f^I_{A,V} \) can be found in the appendix.

Considering the values for \( g_\eta \) chosen for model 2, for instance, if \( M_{Z'} > 2m_\eta \), the \( Z' \) decays 100\% into \( \eta\eta \). However, smaller values for \( g_\eta \) can suppress this decay and change this ratio. The expression for \( Z' \to \eta\eta \) is given by

\[
\Gamma_{Z'\to\eta\eta} = \frac{1}{12\pi} \frac{1}{M_{Z'}} \frac{g_\eta^2}{g_{\eta\eta}} \sqrt{M_{Z'}^2 - 4m_\eta^2(M_{Z'}^2 + 2m_\eta^2)}. \quad (21)
\]

VI. CONCLUSIONS

We have proposed a model in which a Proca-Wentzel (PW) field consisting of \~ MeV–multi-GeV mass particles are mediators of secluded fermionic DM interactions. From the low-energy phenomenology it is known that electroweak precision data such as the mass of \( W \) boson, the decay width of \( Z \) boson, and some asymmetries, can constrain such a model. For our secluded DM model, the constraints on \( \frac{\xi}{\sqrt{1-M_{Z'}^2/M_{Z}^2}} \) are well satisfied. Indeed, they are much smaller than \( O(10^{-2}) \) \[21\]. Also, the mediator of DM interactions will not induce flavor-changing neutral current processes.

One important difference between our model and other ones presented in the literature is the absence of an extra singlet scalar. Thus, there is no equivalent to a “Higgs scalar” in the PW Lagrangian. This distinguishing characteristic may become important if no \( h \to invisible \) width is observed at LHC; such an result could be a clear signature of the model.

The model that we have explored has characteristics in common with the model presented in \[18\]. Our results agree with theirs in the limit \( M_{Z'}/m_\eta \to 0 \). As those authors point out, there are interesting and testable astrophysical signatures for secluded DM models.

We have here explored one of these signatures in quantitative detail. We have considered the \( \gamma \)-rays that would be produced as a result of cosmic annihilations of secluded WIMPs of mass of \~ 20 and \~ 32 GeV. In particular, we have shown that the secluded DM model proposed here can potentially explain an apparent 2 – 3 GeV energy \( \gamma \)-ray excess in the galactic center region, this being an excess over that expected by taking account of other galactic \( \gamma \)-ray production processes. Such an excess has been inferred from an analysis of Fermi-LAT Gamma-ray Space Telescope data \[10\]. Our results are also consistent with a putative weak \( \sim(2-3)\sigma \) \( \gamma \)-ray signal claimed for the dwarf galaxy Reticulum 2 \[35\] and with the conservative constraints derived from Fermi Gamma-Ray Space Telescope observations in Refs. \[36\].

Most other DM interpretations of the galactic center excess stress the \( b\bar{b} \) channel. However, in our Model 1 light quarks and leptons dominate in producing the DM annihilation \( \gamma \)-rays while in our Model 2 the \( c\bar{c} \) channels dominate over the \( b\bar{b} \) channel. This is a predicted phenomenological difference between our model and other models. At present, owing to the

systematic uncertainties in determining the DM “signal” over the other processes contributing to the \( \gamma \)-ray “background” in the direction of the galactic center a definitive test of this difference is difficult \[11, 37\].

We note that for both of the secluded WIMP models that we consider here, our calculated values for the spin independent couplings with protons are consistent with the present experimental upper limits obtained by the XENON100 and LUX experiments \[6\]. The spin independent couplings with neutrons are suppressed. It is anticipated that by 2020 liquid xenon detectors will have the capability to measure spin independent cross sections as low as \( O(10^{-4}) \) cm\(^2\) \[38\]. Should future laboratory results yield a very small constraint on the WIMP elastic scattering cross section, the DM annihilation hypothesis for explaining the \( \gamma \)-ray excess from the galactic center region would then favor a secluded-WIMP model for the dark matter.

Note: While this work was being prepared other recent DM models that discuss further aspects of the DM interpretation of the galactic center \( \gamma \)-ray excess have been put forth \[22, 39\].

ACKNOWLEDGMENTS

We thank Julian Heeck and Matthew Wood for helpful discussions. E.C.F.S.F. thanks FAPESP for full support under contracts numbers 14/05505-6 and 11/21945-8. VP thanks to CNPq for partial support.

Appendix A: Interactions and vertices of the model

Couplings of \( Z \) and \( Z' \) to SM fermions:

\[
\bar{\psi} \psi Z : \frac{ig}{c_W} \frac{c_{\alpha_\lambda}}{(1 - s_W \xi_{\alpha_\lambda})} \left[ T^3 \frac{s_W^2 (1 - \xi_{\alpha_\lambda}/s_W) Q}{(1 - s_W \xi_{\alpha_\lambda})} \right],
\]

\[
\bar{\psi} \psi Z' : \frac{ig}{c_W} \frac{c_{\alpha_\lambda}(t_{\alpha_\lambda} + s_W \xi)}{(t_{\alpha_\lambda} + s_W \xi)} \left[ T^3 \frac{s_W^2 (1 - \xi_{\alpha_\lambda}/s_W) Q}{(t_{\alpha_\lambda} + s_W \xi)} \right]. \quad (A1)
\]

Triple gauge bosons couplings: Comparing to the SM couplings denoted by \( R \), they will be:

\[
R_{AW^*W^*-1} = 1,
R_{ZWW*-c_\alpha},
R_{ZWW*-s_\alpha} . \quad (A2)
\]
Higgs Couplings:

\[ h f f : -i m_f / v_h, \]
\[ h W W : 2 m_W^2 / v_h, \]
\[ h Z Z : 2 i m_W^2 (c_a - \xi s_w s_o), \]
\[ h Z Z' : 2 i m_Z^2 (s_a + \xi s_w c_a), \]
\[ h Z Z' : 2 i m_W^2 [ (c_a - s_w \xi s_a) (s_a + s_w \xi c_a) ]. \]

Coupling of Dirac Fermion and Z and Z' gauge bosons:

\[ \bar{\eta} \gamma^\mu Z : i g_\eta \frac{1}{\sqrt{1 - g_{\nu B}^2}} s_a, \]
\[ \bar{\eta} \gamma^\mu Z' : i g_\eta \frac{1}{\sqrt{1 - g_{\nu B}^2}} c_a. \]

The interactions of Z' in Eq. (A1) are written in a simplified form. In fact, in order to calculate the Z' decay width we have used couplings in the form of

\[ \mathcal{L}_{NC} = \frac{g}{2 c_w} \sum_i [\bar{\psi}_i \gamma^\mu (g'_{V} - g'_{A} \gamma^5) \psi_i Z_0 + \bar{\psi}_i \gamma^\mu (f'_{V} - f'_{A} \gamma^5) \psi_i Z'_0], \]

where \( g'_{V}, g'_{A} \) denote respectively the vectorial and vector-axial coupling of Z boson with fermions and \( f'_{V}, f'_{A} \) denote these couplings but now for Z' gauge boson. The expressions which relate vectorial and vector axial couplings to left and right-handed couplings are given in Eq. [A6]

\[ g'_{A} = \frac{1}{2} (g'_L - g'_R), \]
\[ g'_{V} = \frac{1}{2} (g'_L + g'_R), \]
\[ f'_{A} = \frac{1}{2} (f'_L - f'_R), \]
\[ f'_{V} = \frac{1}{2} (f'_L + f'_R). \]

The detailed left and right-handed couplings of Z' to fermions are given below.

\[ f'^L_L = c_a (t_a + \xi s_w) \left[ T^3_3 \left( \frac{(t_a + \xi s_w)}{(t_a + \xi s_w)} \right) s_w Q_a \right], \]
\[ f'^L_R = c_a (t_a + \xi s_w) \left[ \left( \frac{(t_a + \xi s_w)}{(t_a + \xi s_w)} \right) s_w Q_a \right], \]
\[ f'^L_L = c_a (t_a + \xi s_w) \left[ T^3_3 \left( \frac{(t_a + \xi s_w)}{(t_a + \xi s_w)} \right) s_w Q_d \right], \]
\[ f'^L_R = c_a (t_a + \xi s_w) \left[ \left( \frac{(t_a + \xi s_w)}{(t_a + \xi s_w)} \right) s_w Q_d \right], \]
\[ f'^R_L = c_a (t_a + \xi s_w) \left[ T^3_3 \left( \frac{(t_a + \xi s_w)}{(t_a + \xi s_w)} \right) s_w Q_e \right], \]
\[ f'^R_R = c_a (t_a + \xi s_w) \left[ \left( \frac{(t_a + \xi s_w)}{(t_a + \xi s_w)} \right) s_w Q_e \right], \]
\[ f'^R_L = c_a (t_a + \xi s_w) T^3_3, \]
\[ f'^R_R = 0. \]

where \( Q_i \) denotes the charge of fermion, \( T^3_3 = -1/2, T^3_3 = 1/2. \)

For completeness we write here the right and left-handed couplings of Z' to fermions.

\[ g'^R_L = [c_a (1 - s_w t_o \xi)] T^3_3 \left( \frac{(1 - t_o \xi)}{s_w t_o \xi} \right) s_w Q_o, \]
\[ g'^R_R = [c_a (1 - s_w t_o \xi)] \left( \frac{(1 - t_o \xi)}{s_w t_o \xi} \right) s_w Q_o, \]
\[ g'^L_L = [c_a (1 - s_w t_o \xi)] T^3_3 \left( \frac{(1 - t_o \xi)}{s_w t_o \xi} \right) s_w Q_d, \]
\[ g'^L_R = [c_a (1 - s_w t_o \xi)] \left( \frac{(1 - t_o \xi)}{s_w t_o \xi} \right) s_w Q_d, \]
\[ g'^L_L = [c_a (1 - s_w t_o \xi)] T^3_3 \left( \frac{(1 - t_o \xi)}{s_w t_o \xi} \right) s_w Q_e, \]
\[ g'^L_R = [c_a (1 - s_w t_o \xi)] \left( \frac{(1 - t_o \xi)}{s_w t_o \xi} \right) s_w Q_e, \]
\[ g'^L_L = [c_a (1 - s_w t_o \xi)] \left( \frac{(1 - t_o \xi)}{s_w t_o \xi} \right) s_w Q_e, \]
\[ g'^L_R = 0. \]

[1] D. Burstein and V. Rubin, Astrophys. J. 297, 423 (1985).
[2] G. Hinshaw, et al., Astrophys. J. Suppl. 208:19 (2013).
[3] J. L. Feng, Ann. Rev. Astron. Astrophys. 48, 495 (2010).
[4] A. Pierce, In Kane, Gordon (ed.) et al.: Perspectives on LHC physics pp. 13-23. D. Cline, arXiv:1406.5200.
[5] R. Bernabei et al. [DAMA and LIBRA Collaborations], Eur. Phys. J. C 67, 39 (2010) arXiv:1002.1028 [astro-ph.GA]. C. E. Aalseth, et al. [CoGeNT Collaboration], Phys. Rev. Lett. 106, 131301 (2011). Z. Ahmed et al. [CDMS-II Collaboration], Phys. Rev. Lett. 106, 131302 (2011).
[6] E. Aprile et al. [XENON100 Collaboration], Phys. Rev. Lett. 109, 181301 (2012). D. S. Akerib et al. [LUX Collaboration], Phys. Rev. Lett. 112, 091303 (2014).
[7] Fox et al. Phys.Rev. D 85, 056011 (2012).
[8] M. Ackermann et al. [Fermi-LAT Collaboration], Phys. Rev. Lett. 107, 241302 (2011). M. Ackermann et al. [Fermi-LAT Collaboration], Phys. Rev. D 89, 042001 (2014). A. Abramowski et al., [H.E.S.S. Collaboration] Phys. Rev. Lett. 114, 081301 (2015).
[9] M. Aguilar et al. Phys. Rev. Lett. 113, 121102 (2014). L. Feng,
