Indirect Detection of Little Higgs Dark Matter

Maxim Perelstein and Andrew Spray

Cornell Institute for High-Energy Phenomenology, Cornell University, Ithaca, NY 14853

Abstract

Little Higgs models with T parity contain an attractive dark matter candidate, the heavy photon. We compute the cross section of the heavy photon annihilation into Z-photon pairs, which turns out to be substantially higher than the previously computed cross section for the two photon final state. Unfortunately, even with this enhancement, the monochromatic photon flux from galactic heavy photon annihilation is unlikely to be detectable by GLAST or the currently operating atmospheric Cerenkov telescopes. We also compute the flux of high-energy neutrinos from the annihilation of the heavy photons captured by the Sun and the Earth. The maximum flux of upward-going muons due to such neutrinos is about 1 yr$^{-1}$km$^{-2}$. 
1 Introduction

The presence of dark matter in the universe has been firmly established by observations of galaxy rotation curves, large scale structure, and cosmic microwave background radiation. The microscopic nature of dark matter, however, remains unknown. One attractive scenario is that it consists of stable weakly-interacting massive particles (WIMPs), with masses around the weak scale, \(\sim\)100-1000 GeV. If the WIMPs were in thermal equilibrium with other species in the early universe, their relic abundance naturally matches the observed dark matter density. Further, many theoretical extensions of the Standard Model (SM) of particle physics include such particles. The most famous example is no doubt the neutralino of most SuperSymmetric (SUSY) models; another well known example is the lightest Kaluza-Klein excitation in models with universal extra dimensions.

In this paper, we consider the dark matter candidate emerging in another popular class of theories, the Little Higgs models [1] (for reviews, see [2, 3]). In these theories, the Higgs is a pseudo-Nambu-Goldstone boson associated with spontaneous global symmetry breaking in an extended electroweak sector, which occurs at a scale \(f\sim1\) TeV. The Higgs acquires a mass due to an explicit breaking of the global symmetries by gauge and Yukawa interactions, but a special “collective” manner in which this explicit breaking occurs ensures that the Higgs is relatively light, \(m_h\ll f\). Several models based on this idea have been constructed, most of which are tightly constrained by precision electroweak data. However, introducing an additional discrete symmetry, the T parity, weakens the precision electroweak constraints significantly [4]. The Littlest Higgs model with T parity (LHT) [5] provides an example of a fully realistic and natural model consistent with precision electroweak data [6], and is sufficiently simple to allow for detailed phenomenological analyses [7, 8, 9, 10].

An interesting consequence of T parity is the stability of the lightest T-odd particle (LTP), typically the heavy partner of the hypercharge gauge boson \(B_H\). (This particle is often referred to as the “heavy photon”, even though it does not couple to the electric charge.) The \(B_H\) is weakly interacting, has a mass in the range 80 – 500 GeV, and it was shown to have the correct relic abundance to account for all of the observed dark matter in certain regions of the parameter space [7, 11]. Potential non-gravitational signatures of the heavy photon dark matter have been discussed in [11] (direct detection and indirect detection via high-energy photons produced by \(B_H\) annihilations in the galactic center) and in [12] (indirect detection via high-energy positrons). In this paper, we continue the evaluation of indirect signatures for this dark matter candidate. First, we complete the previous analysis of the monochromatic photon flux [11] by computing the flux from the \(B_HB_H\rightarrow Z\gamma\) reaction, whose cross section turns out to be substantially larger than the \(B_HB_H\rightarrow 2\gamma\) reaction analyzed in [11]. Unfortunately, even with this enhancement the near-future prospects for the detection of the gamma line do not look promising, especially when the EGRET constraint on the continuous photon flux is taken into account. Second, we evaluate the neutrino flux from the annihilation of the heavy photons captured by the Sun and the Earth. While the predicted fluxes are too low to be observed by the near-future detectors such as ICECUBE, a further improvement of two-three orders of magnitude in
detector sensitivity would allow to probe interesting parts of the parameter space.

2 The Model

The LHT model has been described in detail elsewhere [5, 6, 7]; here, we briefly summarize the features of the model important for our analysis. The extended electroweak sector of the LHT model has a global \( SU(5) \) symmetry, which is spontaneously broken to \( SO(5) \) at a scale \( f \sim 1 \text{ TeV} \). At energies below the cutoff \( \Lambda \sim 4\pi f \), the dynamics of this sector is described by a non-linear sigma model. To incorporate gauge interactions, the subgroup \( \mathcal{G} = (SU(2) \times U(1))^2 \) of the \( SU(5) \) is weakly gauged. The \( T \) parity interchanges the two \( SU(2) \times U(1) \) factors. At the scale \( f \), the gauge symmetry \( \mathcal{G} \) is broken down to \( SU(2) \times U(1) \), identified with the SM electroweak group. The gauge bosons corresponding to the broken generators, \( W_H^\pm, W_H^0 \) and \( B_H \), are T-odd and acquire masses at the scale \( f \):

\[
M(W_H) \approx g f, \quad M \equiv M(B_H) \approx \frac{g' f}{\sqrt{5}} \approx 0.16 f. \tag{1}
\]

The uneaten \( SU(5)/SO(5) \) Nambu-Goldstone bosons decompose into a T-even \( SU(2) \) doublet \( H \), identified with the SM Higgs, and a T-odd \( SU(2) \) triplet \( \Phi \), which acquires a mass \( m_\phi = \sqrt{2} m_h f/v \) at one loop. In the fermion sector, the LHT model contains vector-like T-odd partners for the left-handed SM quarks and leptons. We will assume a common mass scale \( \tilde{M} \) for all of these particles. This choice is motivated both by simplicity and (in the case of the T-odd quarks) by flavor constraints [13]. Finally, the model contains an additional pair of weak-singlet, charge-2/3 quarks: \( T^+_+ \) (T-even) and \( T^-_- \) (T-odd). These are required to ensure the cancellation of the one-loop quadratic divergence in the Higgs mass parameter from the SM top. The mass \( M_{T^-} \) is a free parameter, and

\[
M_{T^+} = M_{T^-} \left(1 - \frac{m_t^2 f^2}{v^2 M_{T^-}^2}\right)^{-1/2}. \tag{2}
\]

With our assumptions, the spectrum and the (renormalizable) couplings of the LHT model are completely described in terms of the three parameters, \( f, \tilde{M}, \) and \( M_{T^-} \), in addition to the familiar parameters of the SM.

Due to the smallness of \( g' \) and the favorable group theory factor, the “heavy photon” \( B_H \) is substantially lighter than \( f \), and is always the lightest among the T-odd bosons of the theory. The fermion mass parameters \( \tilde{M} \) and \( M_{T^-} \) are at the scale \( f \), so it is natural to consider the part of the parameter space where the \( B_H \) is the lightest T-odd particle (LTP). In this region, the heavy photon is stable and can contribute to dark matter. In this paper we assume that all of the observed dark matter density is due to the \( B_H \). Calculations of the relic density have shown [7, 11] that there are three regions of the parameter space in which this is possible. In the first two, the \( B_H \) density in the early universe is controlled by

\[\text{In Eq. (1) and throughout this paper, we neglect corrections of order } v^2/f^2.\]
$B_H$ pair-annihilation through an $s$-channel Higgs. The correct relic density is obtained when the center of mass energy of a $B_H B_H$ collision in the nonrelativistic regime is near, but slightly displaced from, the resonance. There are typically two solutions, one on either side of the resonance. In the third region, the $B_H$ relic density is set by coannihilation processes with T-odd fermions, and hence $M$ must be close to $\tilde{M}$. The precise measurement of the present dark matter density implies tight correlations between the model parameters in these regions. It was found in [11] that they can be approximately described by

$$M = \begin{cases} 
(m_h/2.38) - 10 & \text{“Low” pair annihilation region} \\
(m_h/1.89) + 44 & \text{“High” pair annihilation region} \\
M - 20 & \text{coannihilation region}
\end{cases}$$

(All masses are in GeV.) With the reduced uncertainties from the 3-year WMAP data set [14], $\Omega_{\text{dm}} h^2 = 0.104 \pm 0.009$, the allowed variation of $M$ around the central values in (3) is at most of order ±5 GeV at the 2σ level. We will thus treat the allowed regions as lines defined by Eq. (3).

3 Monochromatic Photon Flux from $B_H B_H \rightarrow Z \gamma / h \gamma$ Annihilation

A promising avenue for indirect WIMP detection is through high energy gamma rays [15]. Since the galaxy is transparent to photons in the interesting energy range, any features in the photon spectrum are preserved. This can both aid in distinguishing the signal from the background, and provide information about the WIMP properties. In particular, a pair-annihilation of WIMPs into two-body final states containing photons leads to a flux of nearly monochromatic gamma rays, allowing for efficient background subtraction. In the context of the LHT model, the monochromatic gamma ray flux from the process $B_H B_H \rightarrow 2\gamma$ was analyzed in Ref. [11]. In this section, we will analyze the additional two processes which produce monochromatic photons, $B_H B_H \rightarrow Z \gamma$ or $h \gamma$.

The $B_H B_H$ annihilation in the LHT model is dominated by processes with an $s$-channel Higgs boson exchange. The cross section for producing a given final state $X$ in this channel can be related to the partial decay width of an (off-shell) Higgs [11]:

$$\sigma_X u \equiv \sigma \left( B_H B_H \rightarrow X \right) u = \frac{g^4 v^2}{72 M^4} \frac{s^2 - 4 s M^2 + 12 M^4 \hat{\Gamma} (h \rightarrow X)}{(s - m^2_h)^2 + m^2_h \hat{\Gamma}^2_h} \sqrt{s}. \quad (4)$$

Here, $u$ is the relative velocity of the annihilating WIMPs; $s \approx 4 M^2$ in the galactic centre; and the hat on $\Gamma$ indicates that the substitution $m_h \rightarrow \sqrt{s}$ should be made in the expression for the on-shell Higgs decay width.

Using the well-known results for the partial Higgs decay width in the $Z \gamma$ channel [16, 17],
we obtain
\[ \hat{\Gamma}(h \to Z\gamma) = \frac{\alpha g^2}{2048 \pi^4 m_W^2} |A_F + A_G|^2 \left(1 - \frac{m_Z^2}{s}\right)^3. \]  

(5)

Here, \( A_F \) and \( A_G \) are the contributions to the matrix element from fermions and gauge bosons, respectively:
\[
A_F = -4 \sum_f \frac{\sqrt{2} y_f \tau_f^{1/2}}{g \tau_f^{1/2}} N_{c_f} Q_f V_f [I_1(\tau_f, \lambda_f) - I_2(\tau_f, \lambda_f)],
\]
\[
A_G = -\sum_g \frac{c_g^{\tau_W}}{\tau_g} V_g Q_g \left\{4(3 - t_W^2)I_2(\tau_g, \lambda_g) + \left[1 + \frac{2}{\tau_g}\right] t_W^2 - \left[5 + \frac{2}{\tau_g}\right] I_1(\tau_g, \lambda_g)\right\},
\]

(6)

where \( t_W = 0.548 \) is the tangent of the Weinberg angle, and
\[
\tau_i \equiv \frac{4m_i^2}{s}, \quad \lambda_i \equiv \frac{4m_i^2}{m_Z^2}.
\]

(7)

The sums run over all particles of the relevant type (including both the SM and the additional particles of the LHT model), with the electric charge, the multiplicity and the coupling to the \( Z \) boson of each particle given by \( Q_i, N_{c_i}, \) and \( V_i \), respectively. For fermions, \( V_i \) denotes the vector part of the \( i\bar{i}Z \) coupling; there is no contribution from the axial part. In particular, for the extra vector-like fermions of the LHT model, we obtain
\[
V_i = \frac{g}{c_W} \left(T_3(i) - s_W^2 Q(i)\right),
\]

(8)

where \( T_3(\tilde{U}, \tilde{N}) = 1/2, T_3(\tilde{D}, \tilde{E}) = -1/2, \) and \( T_3(T_+, T_-) = 0. \) The fermion and vector boson trilinear couplings to the Higgs are given by \( y_f/\sqrt{2} \) and \( c_i g M_W \eta^{\mu\nu} \), respectively. (With this normalization, \( y_f \) are the standard Yukawa couplings for the SM fermions and \( c_W = 1 \) for the SM \( W^\pm \) boson.) The functions \( I_{1,2} \) are given by
\[
I_1(a, b) = \frac{ab}{2(a-b)} + \frac{a^2 b^2}{2(a-b)^2} \left[f(a) - f(b)\right] + \frac{a^2 b}{(a-b)^2} [g(a) - g(b)],
\]
\[
I_2(a, b) = -\frac{ab}{2(a-b)} \left[f(a) - f(b)\right],
\]

(9)

where
\[
f(x) = \begin{cases} \sin^{-1} \left(\sqrt{\frac{x}{2}}\right)^2 & \text{if } x > 1 \\ -\frac{1}{4} \left[\log \left(\frac{1 + \sqrt{1-x}}{1 - \sqrt{1-x}}\right) - i\pi\right]^2 & \text{if } x < 1 \end{cases}
\]
\[
g(x) = \begin{cases} \sqrt{x-1} \sin^{-1} \left(\frac{1}{4}\sqrt{\frac{x}{2}}\right) & \text{if } x > 1 \\ \frac{1}{2} \sqrt{1-x} \left[\log \left(\frac{1 + \sqrt{1-x}}{1 - \sqrt{1-x}}\right) - i\pi\right] & \text{if } x < 1 \end{cases}
\]

(10)

(11)

\[2\text{We neglect the contribution to the Higgs width from the loops of the T-odd scalar } \Phi.\]
Figure 1: Photon fluxes from WIMP annihilation into $Z\gamma$ (solid line) and $\gamma\gamma$ (dashed line), in the “high” pair-annihilation region of the LHT model. The fluxes in the “low” pair-annihilation and coannihilation regions are similar. The fluxes are normalized to $\bar{J}\Delta\Omega = 1$, and scale linearly with this parameter.

For a telescope with line of sight parameterized by $\Psi = (\theta, \phi)$ and an angular acceptance $\Delta\Omega$, the anomalous photon flux due to $B_H B_H \to Z\gamma$ is given by [15]

$$\Phi = \left(5.5 \times 10^{-10} \text{ s}^{-1} \text{ cm}^{-2}\right) \left(\frac{\sigma_{Z\gamma\nu}}{1 \text{ pb}}\right) \left(\frac{100 \text{ GeV}}{M}\right)^2 \bar{J}(\Psi, \Delta\Omega) \Delta\Omega,$$

where the function $\bar{J}$ contains all information about the dark matter distribution in the halo. The photon energy is

$$E_\gamma = M \left(1 - \frac{M_Z^2}{4M^2}\right).$$

The thermal broadening of the line is much smaller than the energy resolution of any existing telescope, and can be neglected.

The flux expected in the “high” pair-annihilation region of the LHT model\(^3\) is plotted in Fig. 1. For comparison, the flux due to $B_H B_H \to 2\gamma$, computed in [11], is also shown. Throughout the parameter space, the $Z\gamma$ final state provides a stronger monochromatic photon flux, with the ratio of the $Z\gamma$ to $\gamma\gamma$ flux varying between about 1.5 and 100. (The largest values of this ratio are obtained for $M \sim 300$ GeV, where the $\gamma\gamma$ cross section is suppressed due to an accidental cancellation.) Similar results are obtained in the “low” pair-annihilation and the coannihilation regions of the parameter space.

Experimental searches for the anomalous high-energy gamma rays using Atmospheric Cerenkov telescopes (ACTs), such as HESS [18] and VERITAS [19], are currently under

\(^3\)In all calculations of photon and neutrino fluxes, we assume $\bar{M} = M_{T-} = f$ in the pair-annihilation regions and $M_{T-} = f$ in the coannihilation region. The impact of varying these parameters on the flux predictions is very small.
way, and the space-based GLAST telescope [20] is expected to perform such a search starting in 2007. Numerical simulations indicate that dark matter density may have a sharp peak at the galactic center, in which case the flux of the anomalous gamma rays would be maximized for a line of sight towards the center of the Milky Way. The principal source of background for a search focusing on the galactic center region is the recently discovered powerful point-like gamma ray source [21], whose spectrum strongly suggests that its nature is astrophysical. Taking into account this background, Zaharijas and Hooper [22] estimate that the minimal WIMP-related monochromatic photon flux required for 5σ discovery is about $10^{-11}$ cm$^{-2}$s$^{-1}$ for GLAST (which is sensitive to photon energies up to about 300 GeV) and $10^{-12}$ cm$^{-2}$s$^{-1}$ for the ACTs (sensitive to $E_\gamma \gtrsim 200$ GeV). The $Z\gamma$ line predicted in the LHT model would be observable at GLAST and the ACTs only if the fluxes are enhanced by a strong spike in the dark matter concentration around the galactic center: depending on the $B_H$ mass, values of $\bar{J}\Delta\Omega$ in the $10^{2}\ldots10^{4}$ range are required. Many (though not all) models of the galactic halo contain such spikes: for example, the profile of Moore et al. [23] predicts $\bar{J}\approx10^{5}$ for $\Delta\Omega=10^{-3}$ sr.

However, the prospects of the future searches are further restricted by the constraints on the continuous component of the photon flux from the observations of the gamma rays from the galactic center in the energy range up to 30 GeV by EGRET [24, 22]. The continuous component of the flux in the LHT model was computed in Ref. [11]. The maximum value of the $Z\gamma$ photon flux compatible with the EGRET constraint (independently of the galactic halo profile) is shown in Fig. 2. It is clear that the maximum flux is substantially below the sensitivity of the current and near-future telescopes throughout the parameter space.

Given the presence of a point-like background source at the galactic center, other regions with possible dark matter overdensity were suggested as potential targets for a search for anomalous gamma rays. These include dwarf spheroidal companion galaxies to the Milky Way such as Saggitarius, Draco and Canis Major [25, 26], as well as the Large Magellanic Cloud and the M87 galaxy [27]. In addition, in models where the galaxy is built up from hierarchical dark matter clustering one should expect localized clumps of dark matter inside the Milky Way halo. The values of $\bar{J}\Delta\Omega$ expected for these objects are model-dependent. For example, the dark matter profiles in dwarf spheroidals surveyed in Ref. [25] give $\bar{J}\Delta\Omega\approx10^{-3}-1$ for $\Delta\Omega=10^{-3},10^{-5}$. For dark matter clumps in the halo, Baltz et. al. [29] estimate that a typical clump would have $\bar{J}\Delta\Omega\approx0.4$ at $\Delta\Omega=1.5\times10^{-4}$. (This estimate uses the data from a simulation by Taylor and Babul [28].) For a space-based telescope such as GLAST, the background flux for these targets can be estimated by a simple power-law extrapolation of the extragalactic gamma ray flux measured by EGRET [30]:

$$\frac{d\Phi}{dE d\Omega} = k \left( \frac{E}{100 \text{ GeV}} \right)^{-2.1},$$

where $k = 8.2 \times 10^{-11}$ cm$^{-2}$s$^{-1}$sr$^{-1}$GeV$^{-1}$. Assuming a telescope with energy resolution $\delta E/E = 0.1$, the signal/background ratio is close to 1 for sources with $\bar{J}\Delta\Omega\sim1$ in a model with $M \sim 100$ GeV. (The signal/background ratio decreases with increasing $M$ and/or decreasing $\bar{J}$.) However, the number of $Z\gamma$ events expected at GLAST ($A = 10^4$ cm$^2$) is well
Figure 2: Maximum $Z\gamma$ photon fluxes consistent with the EGRET bound on the continuous photon flux from WIMP annihilation. Solid line: “high” pair-annihilation region; dashed line: coannihilation region with $m_h = 300$ GeV. Also shown are the sensitivities of GLAST (dotted red line) and HESS (dotted blue line) [22].

below 1 event/year, so that no discovery is possible. Of course, this pessimistic prediction could be proven wrong if dark matter turned out to be significantly stronger clumped at short scales than presently thought, resulting in larger values of $\bar{J}$. Barring this possibility, a telescope with a larger effective area ($A \gtrsim 10^7$ cm$^2$) would be required to begin probing the heavy photon dark matter model in this channel.

The final annihilation process giving rise to monochromatic photon flux, $B_H B_H \to h\gamma$, is strongly suppressed. It cannot occur via an $s$-channel Higgs exchange, because the process $h^* \to h\gamma$ is forbidden by the Ward identity of Quantum Electrodynamics. The proof is as follows: Let the photon and on-shell Higgs have momenta $k$ and $p$, respectively. The matrix element has the form

$$M(h^* \to h\gamma) \equiv \varepsilon_{\mu}^{*}(k) \hat{\mathcal{M}}^{\mu},$$

where $\hat{\mathcal{M}}^{\mu}$ can be decomposed as

$$\hat{\mathcal{M}}^{\mu} = A(k, p) \, k^{\mu} + B(k, p) \, p^{\mu}. \quad (16)$$

The Ward identity requires that $k_{\mu} \hat{\mathcal{M}}^{\mu} = 0$. For the first term this is true since $k^2 = 0$. For the second term, $p \cdot k = 0$ only if the initial Higgs is on-shell (when the process is kinematically forbidden); thus, $B \equiv 0$. However, by the transverse nature of the polarization vector the first term provides no contribution to the matrix element; therefore, $M(h^* \to h\gamma) \equiv 0$. As a result, $B_H B_H \to h\gamma$ may only proceed via box diagrams with T-odd and T-even fermions in the loop, which receive no resonant enhancement. Moreover, the $B_H \bar{Q}q$ and $B_H \bar{L}l$ couplings are of order $g'/10 \approx 0.035$, further suppressing the cross section.
4 Neutrino Fluxes from $B_H$ Annihilations in the Sun and the Earth

Neutrinos produced in annihilations of heavy photons collected in the gravitational wells of the Sun and the Earth provide another potentially observable indirect signature of Little Higgs dark matter. The procedure for evaluating the neutrino fluxes in a given model is well established; a thorough review (in the context of SUSY) is given in [31]. Here we will follow this procedure to compute the neutrino fluxes expected in the LHT model.

The number of WIMPs $N$ collected in the Sun or the Earth obeys

$$\dot{N} = C - AN^2,$$

where $C$ is the capture rate and $A$ is the annihilation rate per WIMP. Schematically, the capture rate is given by

$$C \sim c \langle \sigma_{\text{el}} \rangle \frac{1}{M_{\text{GeV}}^2},$$

where $M_{\text{GeV}}$ is the WIMP mass in units of GeV, the quantity $\langle \sigma_{\text{el}} \rangle$ is essentially the weighted average of the elastic WIMP-nucleus scattering cross sections (in pb) over the atomic composition of the Sun or the Earth, and $c$ is a coefficient determined by the properties of the astronomical body in question: $c \approx 10^{30}$ s$^{-1}$ for the Sun and $c \approx 10^{20}$ s$^{-1}$ for the Earth. (See Ref. [31] for a more detailed discussion.) The only input from particle physics required to compute $C$ is the elastic scattering cross sections, which were computed in Ref. [11] for the LHT model. The annihilation rate per WIMP $A$ is schematically given by

$$A = \langle \sigma_{\text{an}} u \rangle \frac{1}{V_{\text{eff}}},$$

where $\sigma_{\text{an}}$ is the total WIMP annihilation cross section, the average is over the thermal distribution of the WIMPs captured in the Sun or the Earth, and $V_{\text{eff}}$ is the effective volume of the “WIMP-sphere” inside the astronomical body. (For details, see Ref. [31].)

Having computed $C$ and $A$, we solve Eq. (17) to obtain the total WIMP annihilation rate:

$$\Gamma_A = \frac{1}{2} C \tanh^2 \left( t \sqrt{AC} \right),$$

where $t \approx 4.5 \cdot 10^9$ years is the age of the Solar System. The experimental technique best suited to searching for high-energy neutrinos from WIMP annihilation relies on observing an upward-going muon created by a charged-current interaction of a muon neutrino in the rock below the detector. The rate of such muons per unit detector area is given by [31]

$$\Gamma_{\text{detect}} = \left( 2.54 \times 10^{-29} \text{ m}^{-2}\text{yr}^{-1} \right) \frac{\Gamma_A}{s^{-1}} M_{\text{GeV}}^2 \sum_i a_i b_i \sum_F B_F \left( \langle N z^2 \rangle \right)_{F,i} (M),$$

where $i$ are the possible neutrino types, $a_i$ and $b_i$ are (known) coefficients describing the neutrino scattering and muon propagation in the rock, and $F$ are the possible final states.
Figure 3: The rate of neutrino-induced upward-going muon events expected from the heavy photon annihilation in the Sun. Solid line: “low” pair-annihilation region (the flux in the “high” pair-annihilation region is similar); dashed line: coannihilation region with $m_h = 300$ GeV. Also shown is the expected sensitivity of the IceCube detector (red/dotted line).

of WIMP annihilation with branching fractions $B_F$. In the LHT model, the dominant annihilation channels are $W^+W^-$ and $ZZ$. The quantity

$$\langle N_{z^2} \rangle_{F,i}(M) \equiv \int_{E_{th}/M}^{1} \frac{dN_{F,i}(E_i, z)}{dz} z^2 dz$$

is the second moment of the spectrum of neutrino type $i$ from final state $F$. Here $E_{th}$ is the threshold energy of the detector, and $dN_{F,i}/dz$ is the neutrino spectrum, normalized per single WIMP annihilation into the final state $F$. This spectrum is a convolution of the initial neutrino spectrum at the production point with the propagation effects (including neutrino oscillations and absorption) on the way to the detector. In this analysis we used the neutrino spectra computed by Cirelli et al. [32], and assumed a detector with a threshold energy of 50 GeV, representative of the IceCube experiment [33]. Note however that lowering this threshold would not have a substantial effect on the rates, since the sub-threshold contribution to the rate scales as $(E_{th}/M)^3$ and is at most of order 10% throughout the interesting parameter range in the LHT model.

The rate of neutrino-induced upward-going muon events expected from the heavy photon annihilation in the Sun is shown in Fig. 3. The expected sensitivity of the IceCube detector is shown for comparison. The maximum possible rate (achieved at the low end of the allowed LTP mass range, $M \approx 100$ GeV) is about 1 event/yr/km$^2$ in the coannihilation region and about 0.5 events/yr/km$^2$ in pair-annihilation regions. Unfortunately these rates are well below the sensitivity of the IceCube. The sensitivity would need to be improved by a factor of a few hundred to a thousand before the fluxes predicted in the LHT models can be probed.

\footnote{We thank Marco Cirelli for providing us with the updated version of the spectra.}
The rates of events due to heavy photon annihilation in the Earth are even smaller, of the order $10^{-5}$ events/yr/km$^2$ or below throughout the parameter space.

5 Conclusions

In this paper, we discussed two signatures of the heavy photon dark matter predicted by the LHT theory. We computed the cross section of the process $B_H B_H \rightarrow Z\gamma$, leading to a monochromatic gamma ray signature. We also computed the flux of high-energy neutrinos from the annihilation of the heavy photons trapped in the Sun and the Earth. Unfortunately, the near-term prospects for observing both signatures are rather poor. The gamma ray signature could in principle be observed by GLAST, but only if dark matter is very strongly clumped at short distance scales. In the neutrino case, the predicted flux is too small to be observed at the IceCube.

This study complements two previous analyses of the discovery prospects for the LHT dark matter. While all predictions are subject to significant astrophysical uncertainties, it appears that the most promising search channels are the “secondary” gamma rays produced in hadronization and fragmentation of the primary WIMP annihilation products [11], and anomalous high-energy positrons [12]. In these channels, the signal may be observed by the near-future instruments, GLAST in the case of gamma rays and PAMELA and AMS-02 in the case of positrons. On the other hand, the LHT model predicts that no signal will be observed in the near-future direct detection and high-energy neutrino searches, while the monochromatic gamma ray signal is very unlikely. By testing the predicted pattern of signals, astroparticle experiments will provide an important test of the LHT dark matter hypothesis, complementary to the more direct searches for the new particles predicted by the LHT model at the Large Hadron Collider.

Acknowledgments — We are grateful to Dan Hooper for helpful correspondence, and to Andrew Noble for useful discussions. We thank Marco Cirelli for clarifying the results of Ref. [32] and providing us with the updated version of the neutrino flux tables. This research is supported by the NSF grant PHY-0355005.

References

[1] N. Arkani-Hamed, A. G. Cohen and H. Georgi, Phys. Lett. B 513, 232 (2001) [arXiv:hep-ph/0105239]; N. Arkani-Hamed, A. G. Cohen, E. Katz and A. E. Nelson, JHEP 0207, 034 (2002) [arXiv:hep-ph/0206021].

[2] M. Schmaltz and D. Tucker-Smith, Ann. Rev. Nucl. Part. Sci. 55, 229 (2005) [arXiv:hep-ph/0502182].

[3] M. Perelstein, arXiv:hep-ph/0512128.
[4] H. C. Cheng and I. Low, JHEP 0309, 051 (2003) [arXiv:hep-ph/0308199]; JHEP 0408, 061 (2004) [arXiv:hep-ph/0405243].

[5] I. Low, JHEP 0410, 067 (2004) [arXiv:hep-ph/0409025].

[6] J. Hubisz, P. Meade, A. Noble and M. Perelstein, JHEP 0601, 135 (2006) [arXiv:hep-ph/0506042].

[7] J. Hubisz and P. Meade, Phys. Rev. D 71, 035016 (2005) [arXiv:hep-ph/0411264].

[8] A. Freitas and D. Wyler, arXiv:hep-ph/0609103.

[9] A. Belyaev, C. R. Chen, K. Tobe and C. P. Yuan, arXiv:hep-ph/0609179.

[10] M. Carena, J. Hubisz, M. Perelstein and P. Verdier, arXiv:hep-ph/0610156.

[11] A. Birkedal, A. Noble, M. Perelstein and A. Spray, Phys. Rev. D 74, 035002 (2006) [arXiv:hep-ph/0603077].

[12] M. Asano, S. Matsumoto, N. Okada and Y. Okada, arXiv:hep-ph/0602157.

[13] J. Hubisz, S. J. Lee and G. Paz, JHEP 0606, 041 (2006) [arXiv:hep-ph/0512169]; M. Blanke, A. J. Buras, A. Poschenrieder, C. Tarantino, S. Uhlig and A. Weiler, arXiv:hep-ph/0605214; M. Blanke, A. J. Buras, A. Poschenrieder, S. Recksiegel, C. Tarantino, S. Uhlig and A. Weiler, arXiv:hep-ph/0609284.

[14] D. N. Spergel et al., arXiv:astro-ph/0603449.

[15] L. Bergstrom, P. Ullio and J. H. Buckley, Astropart. Phys. 9, 137 (1998) [arXiv:astro-ph/9712318].

[16] M. A. Shifman, A. I. Vainshtein, M. B. Voloshin and V. I. Zakharov, Sov. J. Nucl. Phys. 30, 711 (1979) [Yad. Fiz. 30, 1368 (1979)].

[17] J. F. Gunion, H. E. Haber, G. L. Kane and S. Dawson, The Higgs Hunter’s Guide, Persues, Cambridge, MA, 1990; see also arXiv:hep-ph/9302272.

[18] J. A. Hinton [The HESS Collaboration], New Astron. Rev. 48, 331 (2004) [arXiv:astro-ph/0403052].

[19] T. C. Weekes et al., Astropart. Phys. 17, 221 (2002) [arXiv:astro-ph/0108478].

[20] A. Morselli, A. Lionetto, A. Cesarini, F. Fucio and P. Ullio [GLAST Collaboration], Nucl. Phys. Proc. Suppl. 113, 213 (2002) [arXiv:astro-ph/0211327].

[21] K. Kosack et al. [The VERITAS Collaboration], Astrophys. J. 608, L97 (2004) [arXiv:astro-ph/0403422]; K. Tsuchiya et al. [CANGAROO-II Collaboration], Astrophys. J. 606, L115 (2004) [arXiv:astro-ph/0403592]; F. Aharonian et al. [The HESS Collaboration], Astron. Astrophys. 425, L13 (2004) [arXiv:astro-ph/0408145].
[22] G. Zaharijas and D. Hooper, Phys. Rev. D 73, 103501 (2006) [arXiv:astro-ph/0603540].

[23] B. Moore, F. Governato, T. Quinn, J. Stadel and G. Lake, Astrophys. J. 499, L5 (1998) [arXiv:astro-ph/9709051]; B. Moore, T. Quinn, F. Governato, J. Stadel and G. Lake, Mon. Not. Roy. Astron. Soc. 310, 1147 (1999) [arXiv:astro-ph/9903164].

[24] D. Hooper and B. L. Dingus, Phys. Rev. D 70, 113007 (2004) [arXiv:astro-ph/0210617]; arXiv:astro-ph/0212509.

[25] N. W. Evans, F. Ferrer and S. Sarkar, Phys. Rev. D 69, 123501 (2004) [arXiv:astro-ph/0311145].

[26] L. Bergstrom and D. Hooper, Phys. Rev. D 73, 063510 (2006) [arXiv:hep-ph/0512317]; S. Profumo and M. Kamionkowski, JCAP 0603, 003 (2006) [arXiv:astro-ph/0601249].

[27] E. A. Baltz, C. Briot, P. Salati, R. Taillet and J. Silk, Phys. Rev. D 61, 023514 (2000) [arXiv:astro-ph/9909112]; N. Fornengo, L. Pieri and S. Scopel, Phys. Rev. D 70, 103529 (2004) [arXiv:hep-ph/0407342].

[28] J. E. Taylor and A. Babul, Mon. Not. Roy. Astron. Soc. 348, 811 (2004) [arXiv:astro-ph/0301612].

[29] E. A. Baltz, M. Battaglia, M. E. Peskin and T. Wizansky, arXiv:hep-ph/0602187.

[30] P. Sreekumar et al. [EGRET Collaboration], Astrophys. J. 494, 523 (1998) [arXiv:astro-ph/9709257].

[31] G. Jungman, M. Kamionkowski and K. Griest, Phys. Rept. 267, 195 (1996) [arXiv:hep-ph/9506380].

[32] M. Cirelli, N. Fornengo, T. Montaruli, I. Sokalski, A. Strumia and F. Vissani, Nucl. Phys. B 727, 99 (2005) [arXiv:hep-ph/0506298].

[33] P. A. Toale [The ICECUBE Collaboration], arXiv:astro-ph/0607003.