Probing Kaluza-Klein Dark Matter with Neutrino Telescopes

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Abstract

In models in which all of the Standard Model fields live in extra “universal” dimensions, the lightest Kaluza-Klein (KK) particle can be stable. Calculations of the one-loop radiative corrections to the masses of the KK modes suggest that the identity of the lightest KK particle (LKP) is mostly the first KK excitation of the hypercharge gauge boson. This LKP is a viable dark matter candidate with an ideal present-day relic abundance if its mass is moderately large, between 600 to 1200 GeV. Such weakly interacting dark matter particles are expected to become gravitationally trapped in large bodies, such as the Sun, and annihilate into neutrinos or other particles that decay into neutrinos. We calculate the annihilation rate, neutrino flux and the resulting event rate in present and future neutrino telescopes. The relatively large mass implies that the neutrino energy spectrum is expected to be well above the energy threshold of AMANDA and IceCube. We find that the event rate in IceCube is between a few to tens of events per year.
1 Introduction

The premiere astrophysical conundrum is the nature and identity of dark matter. The accumulated body of evidence in favor of the existence of dark matter is by now overwhelming: Studies of the cosmic microwave background [1], high redshift supernovae [2], galactic clusters and galactic rotation curves [3] indicate that the matter density of the universe is \( \Omega_M \simeq 0.3 - 0.4 \). Constraints from big-bang nucleosynthesis, however, limit the baryonic matter density to a small fraction of this number [4]. Furthermore, the observed density of luminous matter is also very small, \( \Omega_L < 0.01 \) [5]. Therefore, the vast majority of the mass in the universe is dark. Additionally, cosmic microwave background studies and large scale structure formation requires that the majority of the dark matter be cold (non-relativistic) [6].

Dark matter could exist in several forms. Perhaps the most interesting possibility is a neutral, stable, weakly interacting particle arising from physics beyond the Standard Model. Candidates for such an animal abound, including the lightest supersymmetric particle, the axion, etc. The possibility of a stable Kaluza-Klein (KK) excitation as particle dark matter was raised many years ago [7] and more recently [8]. Models in which all of the Standard Model fields propagate in “universal” extra dimensions [9] (for earlier work, see [10]) provide the most natural home for KK dark matter [11, 12, 13]. This is because bulk interactions do not violate higher dimensional momentum conservation (KK number), and in these models all of the couplings among the Standard Model particles arise from bulk interactions. To generate chiral fermions at the zero mode level, the extra compact dimension(s) must be modded out by an orbifold. For five dimensions this is \( S^1/Z_2 \), while in six dimensions \( T^2/Z_2 \) is suitable and has other interesting properties [9] including motivation for three generations from anomaly cancellation [14] and the prevention of fast proton decay [15]. An orbifold does, however, lead to some of the less appealing aspects of the model. Brane-localized terms can be added to both orbifold fixed points that violate KK number. If these brane localized terms are symmetric under the exchange of the two orbifold fixed points, then a remnant of KK number conservation remains, called KK parity. All odd-level KK modes are charged under this discrete symmetry thereby ensuring that the lightest level-one KK particle (LKP) does not decay. This is entirely analogous to exactly conserved R-parity in supersymmetric models which ensures the lightest supersymmetric particle is stable. The stability of the LKP suggests it could well be an interesting dark matter candidate.

The identity of the lightest KK particle crucially depends on the mass spectrum of the first KK level. At tree-level, the mass of each excitation is simply

\[
\left( m_i^1 \right)^2 = \frac{1}{R^2} + \left( m_i^0 \right)^2
\]

where \( R \) is the compactification radius that could be as large as \( 1/(300 \text{ GeV}) \) without conflict with experiment [9]. However, brane-localized terms can be added on the orbifold fixed points that significantly modify the masses and higher dimensional wavefunctions. The tree-level (matching) contributions at the cutoff scale of the higher dimensional theory are not calculable, but can be estimated using naive dimensional analysis. The result is that the size of these terms are
suppressed by a volume factor, of order $\Lambda R$ where $\Lambda$ is the cutoff. More importantly, these brane-localized terms are renormalized upon evolving from the matching scale to the mass scale of the light KK modes [16]. For universal extra dimensions, Cheng, Matchev, and Schmaltz showed that the tree-level mass formula receives significant one-loop radiative corrections from log-enhanced brane-localized terms on the orbifold fixed points [11]. These radiative corrections are, in many cases, larger than the shifts in the tree-level masses resulting from the masses of the zero modes. Indeed, here we will generally assume that these contributions dominate over the tree-level volume-suppressed matching contributions. Then, with the further assumption that the KK excitation of the Higgs does not receive a (significant) brane-localized negative contribution to its mass, the identity of the lightest KK state is identified as the first KK excitation of the photon. Like the ordinary photon, the KK photon is an admixture between the first KK hypercharge gauge boson and the first KK neutral SU(2) gauge boson. However, this identification is somewhat misleading, as [11] point out, since the mixing angle for the level-one KK gauge bosons is generally much smaller than the Weinberg angle. A leading order approximation to the mass of the lightest $B^1$ state is

$$m^2_{B^1} \simeq \frac{1}{R^2} \left[ 1 + \frac{g'^2}{16\pi^2} \left( -\frac{39\zeta(3)}{2\pi^2} - \frac{1}{3} \ln \frac{\Lambda R}{2\pi v^2} \right) \right]$$

(2)

neglecting higher order $\mathcal{O}(\sin \theta^1 \nu^2)$ corrections. This approximation is equivalent to identifying LKP $\equiv \gamma^1 \simeq B^1$, which we do for the remainder of the paper.

The relic density of the $B^1$ has been calculated in a recent paper by Servant and Tait [12]. Assuming the LKPs were once in thermal equilibrium, they found that the relic density is in the favorable region for providing the cold dark matter of the universe, $\Omega_{B^1} h^2 = 0.16 \pm 0.04$, when the mass is moderately heavy, between 600 to 1200 GeV. The range of mass depends on the relative importance of coannihilation with KK modes near in mass to the LKP. We shall see later that coannihilation is active throughout the parameter space when the KK mass spectrum is obtained using the radiative corrections arising from renormalized brane-localized terms. Direct detection of the LKP as dark matter has been considered by Cheng, Feng, and Matchev [13]. They emphasized that unlike the case of a neutralino LSP, the bosonic nature of the LKP means there is no chirality suppression of the annihilation signal into fermions. The annihilation rate of the LKP is therefore roughly proportional to the (hypercharge)$^4$ of the final state, leading to a large rate into leptons. The large mass of the LKP suggests that dark matter detection experiments sensitive to very heavy mass relics should be among the most promising methods of detection. In this paper, we explore the signal expected at present and future high energy neutrino telescopes that can probe precisely this type of heavy dark matter.1

2 Relic density and KK mass spectrum

The relic density of $B^1$'s depends on the $B^1$ mass, the annihilation cross section, and the coannihilation rate. The annihilation and coannihilation cross sections are determined by SM couplings

1Note that [12] [13] also remarked on the potential importance of indirect detection at neutrino telescopes.
and the mass spectrum of the first KK level. In [12], the relic density of $B_1^1$'s was calculated in the following approximation: the masses of all of the first level KK modes are equal, and electroweak breaking effects (i.e., fermion and gauge boson masses) are neglected. In the low-velocity limit their result is

$$\langle \sigma v \rangle = \frac{95g_1^4}{324\pi m_{B_1}^2},$$

into fermion final states. There are also annihilations into Higgs boson pairs, but this is only a few percent additive correction to the above and so can be neglected. A more general calculation for arbitrary masses of intermediate state KK fermions is straightforward, but as we will see, not necessary to obtain the quantitative results we present below.

The full KK mass spectrum with one-loop corrections is presented in [11]. The general expression for the correction to the level one KK fermion masses is given by

$$m_{f_1} = \frac{1}{R} \left[ 1 + \frac{9}{32\pi^2} \ln \Lambda R \sum_a C_a(f) g_a^2 \right].$$

(The Yukawa corrections to this formula can be safely ignored since the KK top plays a negligible role in the analysis to follow.) The sum is over all of the SM gauge groups, and $C_a(f)$ is the quadratic Casimir for the fermion transforming under the group $a$. The scale $\Lambda$ is the strong coupling cutoff scale of the extra dimensional theory. An important observation from this general formula is that the fractional shift in the mass to the KK quarks is generally an order of magnitude larger than for the KK leptons. This is simply due to the much larger QCD corrections over the electroweak corrections. In particular, consider the following fractional shifts

$$r_{\ell_R} \equiv \frac{m_{\ell_R} - m_{B_1}}{m_{B_1}} = \frac{g^2}{32\pi^2} \left[ \frac{28}{3} \ln \Lambda R + \frac{39\zeta(3)}{2\pi^2} - (2\pi R v)^2 \right],$$

$$r_{u_R} \equiv \frac{m_{u_R} - m_{B_1}}{m_{B_1}} = \frac{6g_2^2}{16\pi^2} \ln \Lambda R + \frac{g^2}{32\pi^2} \left[ \frac{13}{3} \ln \Lambda R + \frac{39\zeta(3)}{2\pi^2} - (2\pi R v)^2 \right],$$

$$r_{d_R} \equiv \frac{m_{d_R} - m_{B_1}}{m_{B_1}} = \frac{6g_2^2}{16\pi^2} \ln \Lambda R + \frac{g^2}{32\pi^2} \left[ \frac{4}{3} \ln \Lambda R + \frac{39\zeta(3)}{2\pi^2} - (2\pi R v)^2 \right].$$

The non-logarithmically enhanced terms contribute at most a percent (for $1/R = 300$ GeV), and are negligible for the radii of interest here ($1/R \gtrsim 600$ GeV). The ratio of the fractional shifts is

$$\frac{r_{q_R}}{r_{\ell_R}} \simeq \frac{18g_2^2}{14g^2} \simeq 12 \quad \text{where} \quad r_{q_R} \simeq r_{u_R} \simeq r_{d_R}.$$

Notice that the ratio is independent of the compactification radius and also independent of the cutoff scale. Hence, the relative KK particle mass hierarchy is fixed by the just the SM couplings.

What is the size of the radiative correction for a given KK mode? This is dependent on the cutoff scale in theory. Strong coupling in extra dimensional theories appears when $g^2 N$ is order one, where $g$ is the coupling and $N$ is the number of (KK) particles exchanged. For five universal dimensions this has been estimated to be $\Lambda \simeq 20/R$ [9]. This results in the fractional
shifts $r_{B^1_{LR}} \simeq 0.011$ and $r_{q^1_R} \simeq 0.14$. For six universal dimensions the fractional shifts are about half as big.

Several observations based on this discussion of the KK mass spectrum are in order: First, the right-handed KK lepton is within a percent or so of the $B^1$ mass, implying that coannihilation of the $B^1$ with $\ell^1_R$ is always important. This means that for the $B^1$ to have the appropriate abundance to be the dark matter of the Universe, it must have a mass between about 600 to 800 GeV using the results of [12]. Second, we see that the estimate of the relic density [4] in which all of the level one KK modes were taken to have the same mass is a reasonably good approximation since none of the KK modes is more than about 20% heavier than the $B^1$.

3 Capture and annihilation in the Sun

The calculation of the flux of neutrinos from WIMP annihilations in the Sun (and Earth) has been explored in some detail, especially for the case of neutralino dark matter [17]. The basic idea is to begin with the relatively well-known local dark matter density from the galactic rotation data, compute the interaction cross section of the WIMPs with nuclei in the Sun, compare the capture rate with the annihilation rate to determine if these processes are in equilibrium, and then compute the flux of neutrinos that result from this rate of WIMP capture and annihilation.

There are two separate channels by which WIMPs can scatter off nuclei in the Sun: spin-dependent interactions and spin-independent interactions. For accretion from spin-dependent scattering, the capture rate is [18]

$$C_{SD} \cong 3.35 \times 10^{18} \text{s}^{-1} \left( \frac{\rho_{\text{local}}}{0.3 \text{ GeV/cm}^3} \right) \left( \frac{270 \text{ km/s}}{v_{\text{local}}} \right)^3 \left( \frac{\sigma_{H,SD}}{10^{-6} \text{ pb}} \right) \left( \frac{1000 \text{ GeV}}{m_{B^1}} \right)^2$$  (9)

where $\rho_{\text{local}}$ is the local dark matter density, $\sigma_{H,SD}$ is the spin-dependent, WIMP-on-proton (hydrogen) elastic scattering cross section, $v_{\text{local}}$ is the local rms velocity of halo dark matter particles and $m_{B^1}$ is our dark matter candidate. The analogous formula for the capture rate from spin-independent (scalar) scattering is [18]

$$C_{SI} \cong 1.24 \times 10^{18} \text{s}^{-1} \left( \frac{\rho_{\text{local}}}{0.3 \text{ GeV/cm}^3} \right) \left( \frac{270 \text{ km/s}}{v_{\text{local}}} \right)^3 \left( \frac{2.6 \sigma_{H,SI} + 0.175 \sigma_{He,SI}}{10^{-6} \text{ pb}} \right) \left( \frac{1000 \text{ GeV}}{m_{B^1}} \right)^2$$  (10)

Here, $\sigma_{H,SI}$ is the spin-independent, WIMP-on-proton elastic scattering cross section and $\sigma_{He,SI}$ is the spin-independent, WIMP-on-helium elastic scattering cross section. Typically, $\sigma_{He,SI} \simeq 16.0 \sigma_{H,SI}$. The factors of 2.6 and 0.175 include information on the solar abundances of elements, dynamical factors and form factor suppression.

Although these two rates appear to be comparable in magnitude, the spin-dependent cross section is typically three to four orders of magnitude larger than the spin-independent cross section [13]. Therefore, solar accretion by spin-dependent scattering dominates, and hereafter we can safely ignore spin-independent scattering. The microscopic elastic scattering spin-dependent
cross section of a $B^1$ off of a proton was calculated in [13]. Their result can be well-approximated by

$$
\sigma_{H,SD} = \frac{g'^4 m_p^2}{648 \pi m_{B^1}^2 r_{q_{ik}}^2} \left(4 \Delta^p_u + \Delta^p_d + \Delta^p_s\right)^2
$$

(11)

since the scattering cross section is dominated by exchanges of right-handed KK quarks due to their larger hypercharge over left-handed KK quarks. Here, $m_p$ is the mass of the proton and the $\Delta^p_q$'s parameterize the fraction of spin carried by a constituent quark $q$ [19].

$$
\Delta^p_u = 0.78 \pm 0.02 \quad , \quad \Delta^p_d = -0.48 \pm 0.02 \quad , \quad \Delta^p_s = -0.15 \pm 0.07
$$

(12)

Inserting the spin-fractions, we obtain

$$
\sigma_{H,SD} = 0.9 \times 10^{-6} \text{ pb} \left(\frac{1000 \text{ GeV}}{m_{B^1}}\right)^4 \left(\frac{0.14}{r_{q_{ik}}}\right)^2.
$$

(13)

If the capture rates and annihilation cross sections are sufficiently high, the Sun will reach equilibrium between these processes. For $N$ number of $B^1$'s in the Sun, the rate of change of this number is given by

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

(14)

where $C^\odot$ is the capture rate and $A^\odot$ is the annihilation cross section times the relative WIMP velocity per volume. $C^\odot$ was given in [9], while $A^\odot$ is

$$
A^\odot = \frac{\langle \sigma v \rangle}{V_{\text{eff}}}
$$

(15)

where $V_{\text{eff}}$ is the effective volume of the core of the Sun determined roughly by matching the core temperature with the gravitational potential energy of a single WIMP at the core radius. This was found in [20] to be

$$
V_{\text{eff}} = 1.8 \times 10^{26} \text{ cm}^3 \left(\frac{1000 \text{ GeV}}{m_{B^1}}\right)^{3/2}.
$$

(16)

The present $B^1$ annihilation rate is

$$
\Gamma = \frac{1}{2} A^\odot N^2 = \frac{1}{2} C^\odot \tanh^2 \left(\sqrt{C^\odot A^\odot} t_\odot\right)
$$

(17)

where $t_\odot \simeq 4.5$ billion years is the age of the solar system. The annihilation rate is maximized when it reaches equilibrium with the capture rate. This occurs when

$$
\sqrt{C^\odot A^\odot} t_\odot \gg 1.
$$

(18)

Combining our expression for the capture and annihilation rate [using $\langle \sigma v \rangle$ from [3]], we find

$$
\sqrt{C^\odot A^\odot} t_\odot = 2.4 \left(\frac{1000 \text{ GeV}}{m_{B^1}}\right)^{13/4} \frac{0.14}{r_{q_{ik}}}.
$$

(19)
process & annihilation fraction \\
$B^1 B^1 \rightarrow \nu_e \nu_e, \nu_\mu \nu_\mu, \nu_\tau \nu_\tau$ & $r_{f^1} = 0$ & $r_{q^1} = 0.14$ \\
$\rightarrow e^+ e^-, \mu^+ \mu^-, \tau^+ \tau^-$ & 0.012 & 0.014 \\
$\rightarrow \bar{u}u, \bar{c}c, t\bar{t}$ & 0.20 & 0.23 \\
$\rightarrow \bar{d}d, \bar{s}s, \bar{b}b$ & 0.11 & 0.077 \\
$\rightarrow \phi\phi^*$ & 0.007 & 0.005 \\

Table 1: The relative annihilation fraction into various final states. The numbers shown are not summed over generations, and the Higgs mass was assumed to be lighter than $m_{B^1}/2$.

Hence, throughout the mass range that leads to the ideal relic abundance of $B^1$'s considered by [12], we find that the Sun either reaches or nearly reaches equilibrium between $B^1$ capture and annihilation.

The Earth, being less massive and accreting dark matter only by scalar interactions, captures particles much more slowly than the Sun. For the optimistic case of $m_{B^1} = 500$ GeV with $\sigma_H \sim 10^{-6}$ pb (scalar only), the ratio of the age of the solar system to the equilibrium time is on the order of $10^{-5}$, which corresponds to a $10^{-10}$ suppression in the annihilation rate. We find no scenario in which KK dark matter annihilation in the Earth provides an observable signal.

## 4 Event rates and prospects for detection

Now that we have determined the annihilation rate in the Sun, we need to determine the outgoing flux of detectable particles. This corresponds to determining the annihilation fraction directly into muon neutrinos, as well as indirectly through decays. We are interested exclusively in muon neutrinos since at the energies relevant to $B^1$ annihilation, neutrino telescopes only observe muon tracks generated in charged-current interactions. In the approximation that all heavier level one KK modes have the same mass, the relative annihilation fraction can be determined from simply the hypercharge of the final state fermions. This is shown in Table 1 for the column $r_{f^1} = 0$. However, the annihilation into KK quarks is slightly further suppressed since the KK quarks are slightly heavier than the KK leptons. Using $r_{q^1} = 0.14$, our estimates of the modified branching fractions are shown in Table 1. Clearly the relative annihilation fractions are not particularly sensitive to the details of the spectrum so long as KK leptons are the same mass or lighter than KK quarks as the one-loop radiative corrections suggest.

Neutrinos are generated in annihilations directly, but can also be produced in the decays of tau leptons, quarks and Higgs bosons generated in annihilations. Only very short lived particles contribute to secondary neutrino flux, as longer lived particles lose the majority of their energy from scattering in the Sun before decaying. Bottom and charm quarks lose energy from
hadronization before decaying. Although top quarks do not hadronize, they generate neutrinos only through gauge bosons and bottom quarks generated in their decay. Neutrinos can also lose energy traveling through the Sun [21, 22, 23]. Muon neutrinos that escape the Sun then travel to the Earth where a small fraction of them are converted to muons through charged current interactions. Neutrino telescopes observe high energy muon neutrinos by identifying a muon track in the detector medium generated in a charged-current interaction (for a review, see [24]).

To calculate the event rate in neutrino telescopes, we first must calculate the spectrum of neutrinos from annihilation directly to neutrinos as well as from other annihilation modes that lead indirectly to neutrinos from decays. Calculating the flux of indirectly produced neutrinos requires modeling the physics of the final state, including hadronization and decay, and must be done numerically. Fortunately, the indirect sources have been previously calculated in supersymmetric models for neutralino annihilation in the Sun. The flux of muon neutrinos from the indirect final states of leptons, quarks, gauge bosons, and Higgs particles has been simulated and parameterized by Edsjo [21]. We use this parameterization, replacing the neutralino with the LKP, and using the relative annihilation rates into the various final states as shown in Table 1.

In this parameterization the flux of muons expected at an Earth-based detector is

$$\frac{d\Phi_\mu}{dz}(m_{B_1}, z) = \frac{p_1 m_{B_1}^2}{1 + \exp \left( z - \frac{m_{B_1}}{p_4} \right)} \left[ 1 - \exp \left( -\frac{p_5 m_{B_1}}{z \rho_8} \right) \right] \exp \left( -\frac{p_7 m_{B_1} z}{p_8} \right)$$

where $z \equiv E_\mu/m_{B_1}$ and $p_1 \ldots p_8$ are the parameters fitted to the simulation results that can be found in [21]. The accuracy of this parameterization was estimated to be within about 15% so long as the threshold is not too close to the LKP mass, $E_{\mu}^{th} \lesssim 0.2 m_{B_1}$. We have used a muon energy threshold of 50 GeV, although our results are not very sensitive to this choice. This energy threshold is well below the masses of the LKPs considered here, and thus well within the range where this parameterization is expected to be valid.

The calculation of the direct neutrino signal was done separately. Annihilation of LKPs into neutrinos results in a mono-energetic spectrum of neutrinos. Of course not all of the neutrinos escape the Sun, due to interactions in the solar medium. The probability of a neutrino escaping the Sun without interacting is given by [23]

$$P = e^{-E_\nu/E_k}$$

where $E_k \simeq (130, 160, 200, 230)$ GeV for $(\nu_\mu, \nu_\tau, \bar{\nu}_\mu, \bar{\nu}_\tau)$. We make the conservative approximation that all neutrinos that interact in the Sun are absorbed rather than regenerated. Note that this exponential suppression of very high energy neutrinos emitted from the Sun is implicitly included in the parameterization, Eq. (20).

Neutrinos exiting the Sun travel to the Earth and then occasionally interact through charged current interactions. The mono-energetic spectrum of muon neutrinos becomes a distribution of muon energies [25] due to the conversion interactions. The resulting muon travels a distance $R_\mu$ before its energy falls below the threshold energy $E_{th}$. This distance, called the muon range, is
Figure 1: The spectrum of muons at the Earth generated in charged-current interactions of muon neutrinos generated in the annihilation of 600 GeV (left side) and 1000 GeV (right side) dark matter particles. The elastic scattering cross section used for capture in the Sun was fixed at $10^{-6}$ pb for both graphs. The rates are proportional to that cross section. We used the branching ratios given in Table I.

given by $[26]$

$$R_\mu \simeq \frac{1}{\rho \beta} \ln \left[ \frac{\alpha + \beta E_\mu}{\alpha + \beta E_{th}} \right]$$  \hspace{1cm} (22)$$

where $\rho$ is the density of the detector medium, $\alpha \simeq 2.0$ MeV cm$^2$/g and $\beta \simeq 4.2 \times 10^{-6}$ cm$^2$/g.

The effective volume in which a muon producing interaction can occur and be observed is simply the muon range times the effective area of the detector.

Another important effect for all sources of neutrinos is flavor oscillations. In the Sun, for the energies of interest to us here, an originally outgoing muon or tau neutrino oscillates sufficiently so as to randomize its final flavor $[23]$. (Oscillation into electron neutrinos occurs only for a significantly lower energy source of neutrinos.) This gives a significant increase to the number of muon neutrinos that result from the annihilation process $B^1\bar{B}^1 \rightarrow \tau^+\tau^-$, since every $\tau$ decay results in a tau neutrino that has about an even chance of exiting the Sun as a muon neutrino.

In Fig. 1 we show the muon flux at the surface of the Earth from $B^1\bar{B}^1$ annihilations in the Sun for $m_{B^1} = 600$ and 1000 GeV, respectively. For the purposes of comparison, the spin-dependent cross section was fixed at $\sigma_{H,SD} = 10^{-6}$ pb for both masses. The majority of events result from the indirect source $B^1\bar{B}^1 \rightarrow \tau^+\tau^-$, when the $\tau$’s decay producing a muon neutrino either directly or through flavor oscillations.

Using the neutrino energy spectrum, the event rate expected at an existing or future neutrino
Figure 2: The number of events per year in a detector with effective area equal to one square kilometer. Contours are shown for $r_{q_R} = 0.1$, 0.2, and 0.3. The $r_{q_R} = 0.3$ is shown merely for comparison, since this mass ratio is larger than would be expected from the one-loop radiative correction calculations of the KK mode masses. The relic density of the $B^1$'s lies within the favored range $\Omega_{B^1} h^2 = 0.16 \pm 0.04$ for the solid sections of each line. The relic density is smaller (larger) for smaller (larger) LKP masses [12].

telescope can be calculated. This is shown in Fig. 2 for a detector with an effective area of 1 km$^2$, such as IceCube. Each line corresponds to a different value of $r_{q_R}$. We should emphasize that the expected size of the one-loop radiative corrections from (6)–(7) predict $0.1 \lesssim r_{q_R} \lesssim 0.2$ for $10/R \lesssim \Lambda \lesssim 100/R$. For this range, a kilometer scale neutrino telescope would be sensitive to a $B^1$ with mass up to about 1 TeV. The relic density of the $B^1$ varies from low to high values from left to right in the graph. The range of mass of the $B^1$ that gives the appropriate relic density $\Omega_{B^1} h^2 = 0.14 \pm 0.04$ to be dark matter was estimated from [12] and shown in the figure by the solid sections of the lines. Note that coannihilation with KK quarks is not included in Fig. 2 however this is not likely to significantly affect our results for $r_{q_R} \gtrsim 0.1$ as shown in the figure [12]. Generally the largest uncertainties in our calculations arise from astrophysical sources, such as the local dark matter halo density, which we estimate to be within about a factor of two. Combining the expected size of the one-loop radiative corrections with a relic density appropriate for dark matter, we find that IceCube should see between a few to tens of events per year.

For detectors with smaller effective areas one simply has to scale the curves down by a factor $A/(1 \text{ km}^2)$ to obtain the event rate. In particular, for the first generation neutrino telescopes
including AMANDA, ANTARES, and NESTOR, with effective areas of order 0.1 km$^2$, the event rate could be as high as a few events per year for a $B^1$ mass at the lower end of the solid line region.

A discovery of WIMP annihilations in the Sun would inevitably require a careful analysis of the signal over background. The background for this class of experiments consists of atmospheric neutrinos \[27\] and neutrinos generated in cosmic ray interactions in the Sun’s corona \[28\]. In the direction of the Sun (up to the angular resolution of a neutrino telescope), tens of events above 100 GeV and on the order of 1 event per year above 1 TeV, per square kilometer are expected from the atmospheric neutrino flux. Fortunately, for a very large volume detector with sufficient statistics, this background is expected to be significantly reduced, and possibly eliminated. Furthermore, this rate could be estimated based on the rate from atmospheric neutrinos, a level of about a few events per year. The final background is then further reduced by selecting on judiciously chosen angular and/or energy bins. Neutrinos generated by cosmic ray interactions in the Sun’s corona, however, cannot be reduced in this way. This irreducible background is predicted to be less than a few events per year per square kilometer above 100 GeV.

5 Conclusions

The prospects for indirect detection of $B^1$ dark matter is very promising at kilometer scale neutrino telescopes. Using the one-loop radiative corrections to the KK mass spectrum, we showed that if the $B^1$ lies in the mass range in which it has an acceptable present-day relic density to be the dark matter of the universe, then a 1 km$^2$ neutrino telescope is expected to detect between a few to tens of $B^1$ annihilation events in the Sun per year. This mass range of the $B^1$ is between about 600 to 800 GeV where coannihilations with right-handed KK leptons plays a significant role. The relatively large signal relies on the Sun reaching equilibrium between $B^1$ capture and annihilation, which we explicitly verified for this range of $B^1$ masses.

Although we have focused on the $B^1$ mass range that results in the appropriate dark matter relic density as currently consistent with cosmology, it is straightforward to extrapolate to other masses. In fact, there are two particle physics effects that are expected to lead to a lowering of the $B^1$ mass for a fixed relic density. The first effect is the inclusion of coannihilation with left-handed KK leptons. The second effect, in a six dimensional model, is that there are typically multiple LKPs corresponding to multiple conserved parities. In both of these cases the lowering of the mass of the LKP(s) for a fixed relic density implies a larger event rate at neutrino telescopes. We are therefore optimistic that future detectors will find or exclude this fascinating possibility.

Acknowledgments

We thank Géraldine Servant and Tim Tait for several valuable discussions and for reading the manuscript. We also thank John Beacom for informing us of Ref. \[23\] and for enlightening
discussions concerning its effects on our results. G.D.K. thanks the Aspen Center for Physics and the T-8 group of Los Alamos National Laboratory for hospitality where part of this work was completed. This work was supported in part by a DOE grant No. DE-FG02-95ER40896 and in part by the Wisconsin Alumni Research Foundation.

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