

Detecting Boosted Dark Matter from the Sun with Large Volume Neutrino Detectors

Joshua Berger,¹ Yanou Cui,^{2,3} and Yue Zhao⁴

¹*SLAC National Accelerator Laboratory, Menlo Park, CA 94025, USA*

²*Perimeter Institute for Theoretical Physics, Waterloo, Ontario N2L 2Y5, Canada*

³*Maryland Center for Fundamental Physics, University of Maryland, College Park, MD 20742, USA*

⁴*Stanford Institute of Theoretical Physics, Physics Department,
 Stanford University, Stanford, CA 94305, USA**

We study novel scenarios where thermal dark matter (DM) can be efficiently captured in the Sun and annihilate into boosted dark matter. In models with semi-annihilating DM, where DM has a non-minimal stabilization symmetry, or in models with a multi-component DM sector, annihilations of DM can give rise to stable dark sector particles with moderate Lorentz boosts. We investigate both of these possibilities, presenting concrete models as proofs of concept. Both scenarios can yield viable thermal relic DM with masses $O(1)$ - $O(100)$ GeV. Taking advantage of the energetic proton recoils that arise when the boosted DM scatters off matter, we propose a detection strategy which uses large volume terrestrial detectors, such as those designed to detect neutrinos or proton decays. In particular, we propose a search for proton tracks pointing towards the Sun. We focus on signals at Cherenkov-radiation-based detectors such as Super-Kamiokande (SK) and its upgrade Hyper-Kamiokande (HK). We find that with spin-dependent scattering as the dominant DM-nucleus interaction at low energies, boosted DM can leave detectable signals at SK or HK, with sensitivity comparable to DM direct detection experiments while being consistent with current constraints. Our study provides a new search path for DM sectors with non-minimal structure.

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*Electronic address: jberger@slac.stanford.edu, ycui@perimeterinstitute.ca, zhaoyue@stanford.edu

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I. INTRODUCTION

The evidence for the existence of particulate Dark Matter (DM) [1–3] is extremely compelling. All of the robust signals seen thus far are, however, gravitational and do not pin down the detailed properties of DM. In order to proceed, it is imperative to discover possible non-gravitational interactions of DM with other particles. One discovery strategy is to look for scattering of ambient DM off of SM nuclear targets. This is well motivated by the paradigm of Weakly Interacting Massive Particle (WIMP) DM, where thermal annihilation to SM states predicts DM relic abundance, and generally implies an appreciable DM-nucleon scattering rate based on crossing-symmetry. The majority of halo DM today is expected to be deeply non-relativistic, having a typical velocity of order 10^{-3} . The momentum transfer in collisions of such DM with hadronic matter is only of order a few keV, requiring specialized instruments with low thresholds for the detection of recoils. Such experiments are powerful probes of DM interactions, particularly in scenarios where DM sector is minimal and has s-wave, Spin Independent (SI) interactions with nuclei, which are enhanced due to the coherence of the interactions over the entire nucleus. In non-minimal scenarios, especially when the SI interaction is insignificant, other detection strategies could offer significantly enhanced sensitivity to DM interactions.

One compelling possibility is to look for “boosted DM” that is not of thermal, cosmological origin, but rather is generated at late times, as introduced in [4–6]. Several well-motivated classes of models, such as multi-component DM models and models with non-minimal stabilization symmetries, can generate boosted DM as a product of annihilation or decay in nearby clumps of DM. The relevant processes have forms such as multi-component annihilation $\psi_i\psi_j \rightarrow \psi_k\psi_\ell$ [7–9], semi-annihilation $\psi_i\psi_j \rightarrow \psi_k\phi$ (where ϕ is a non-DM state) [7, 10–13], 3 \rightarrow 2 self-annihilation [14–16], or decay transition $\psi_i \rightarrow \psi_j + \phi$. Annihilation of DM is of particular interest, as it is required in the generic scenario that the DM abundance is set by thermal freeze-out.

In this paper, we focus on two concrete scenarios: models with semi-annihilation of one DM species ψ charged under a Z_3 symmetry, and models with a two-component DM sector with species ψ_A and ψ_B having masses $m_A > m_B$ and ψ_A being the dominant DM component (here, ψ and $\psi_{A,B}$ need not be fermions). In the example of semi-annihilation of Z_3 DM, the DM thermal relic abundance is set by the annihilation process

$$\psi\psi \rightarrow \bar{\psi}\phi, \quad (1)$$

where ϕ is a lighter dark sector state that may decay away. At present times, the non-relativistic ψ undergoes the same annihilation process in the galactic halo or in the Sun if it is captured there. Assuming $m_\phi \ll m_\psi$, the final state $\bar{\psi}$ is produced with a Lorentz boost factor $\gamma = 5/4$. In the example of two-component DM sector, thermal relic abundance of dominant DM ψ_A is set by:

$$\psi_A\bar{\psi}_A \rightarrow \psi_B\bar{\psi}_B. \quad (2)$$

The same annihilation in the present day produces final state ψ_B with Lorentz factor $\gamma = m_A/m_B$.

In the earlier work [5] on boosted DM detection, one of us (YC) and collaborators focused on the above two-component DM model as an example, assuming ψ_A has no direct couplings to SM. In the particular scenario considered in that work, ψ_B had interactions with both electrons and quarks, making electrons the more sensitive scattering target for detection provided that the mediator of these interactions was sufficiently light. If interactions with electrons are suppressed or if the mediator is heavier than $O(10 \text{ MeV})$, then the detection via interactions with quarks is important. These interactions can also enhance the flux of boosted DM by generating a large rate of DM capture and subsequent annihilation in the Sun. For a broad range of parameters, the flux from the Sun will dominate over the flux from the galactic center, making it possible to have observable signals with scattering cross sections of weak scale size or even smaller. Both DM and mediator masses could occupy a wide range of $O(1 - 100) \text{ GeV}$. This paper therefore addresses the scattering off of hadrons as a means of detection of boosted DM.

The large boost factor of this flux of DM opens new avenues for its detection via hadronic recoil. The low threshold requirement for detection of non-relativistic DM is relaxed and it becomes more effective to use much larger detectors that are sensitive only to more energetic recoils. The boost factors obtained in semi-annihilating and two-component DM models are typically non-negligible, but modest, as seen above. The typical momentum transfer is $O(\text{GeV})$ for boosted DM masses larger than a few GeV. The largest experiment of sensitive to hadronic recoils at this energy is

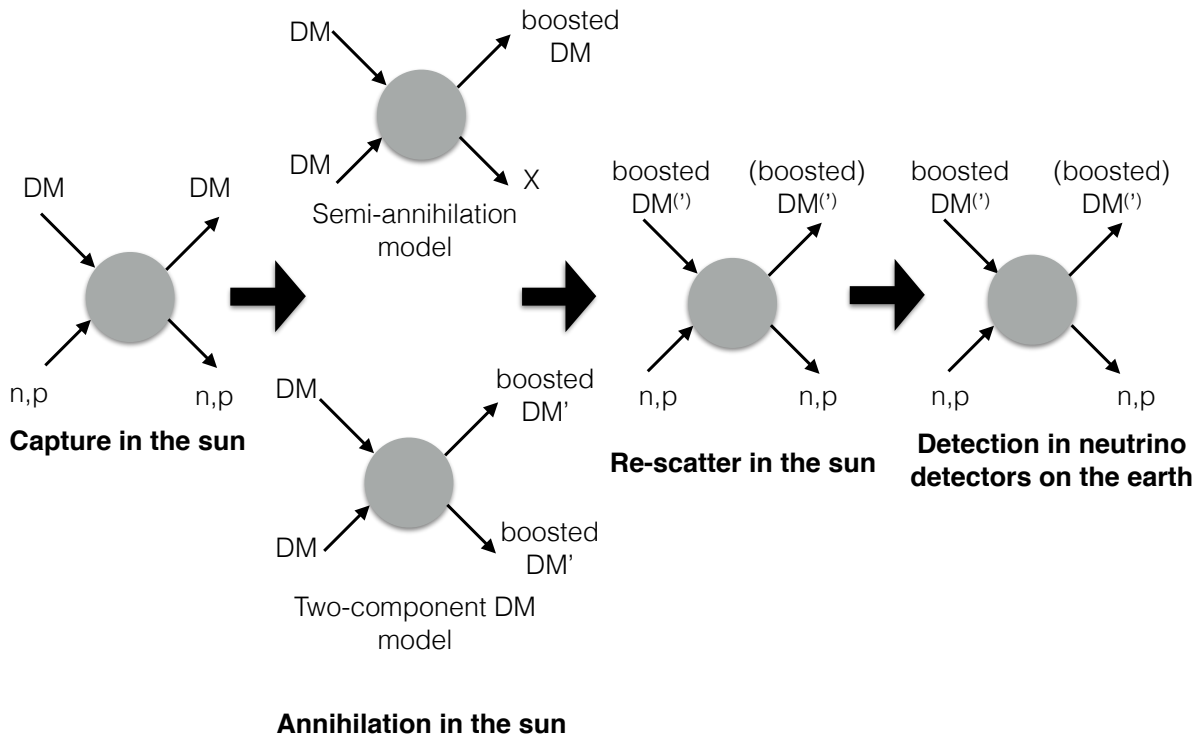


FIG. 1: The chain of processes leading to boosted DM signal from the Sun. DM' denotes the lighter DM in the two-component DM model. X is a lighter dark sector particle that may decay away. Details of the two example models, semi-annihilating DM and two-component DM, are given in Section II.

currently Super-Kamiokande (SK) [17]. Several planned and proposed experiments could have sensitivity, including not only Hyper-Kamiokande (HK) [18], but liquid Argon based detectors such as l LAr TPC and GLACIER [24, 25]. Other large volume detectors such as IceCube/PINGU/MICA [19–21], KM3NeT [22], and ANTARES [23] are best suited to looking for higher energy recoils. Note that these experiments were designed to look for neutrinos and/or proton decay, but can be repurposed to find boosted DM. This is not surprising as interactions of boosted DM with hadronic matter are similar in structure to neutral-current interactions of high energy neutrinos.

Compared to the boosted DM signal from DM annihilation in the GC as studied in [5], the signal from the Sun involves more processes, including the capture, evaporation, annihilation, and rescatter (slow-down) of the DM particles in the Sun, as well as the scattering of boosted DM off nuclei in the terrestrial detectors. In addition, there is the apparent challenge that an appreciable DM-nucleon scattering rate giving rise to detectable signals is likely to be ruled out by the existing DM DD bounds on the thermal non-relativistic component of DM. In our work, we have taken all the above complexities into consideration, and found that there are large classes of models with reasonable range of parameter space giving good detection prospect at SK and HK, while being compatible with conventional DM detection limits. The key feature of these viable models is that the DM scattering off nuclei is dominantly Spin-Dependent (SD) and/or has a velocity-dependence, such as v^2 in the non-relativistic limit. We also note that the signal reach can be further improved by using ionization-based liquid argon neutrino detectors where Cherenkov threshold is irrelevant. In Fig. 1 we illustrate the chain of processes involved in giving rise to the boosted DM signal in which we are interested.

A search for an excess of proton recoils pointing toward the Sun as we propose here is a new search that has not been performed by the SK collaboration. In the search for neutrinos from WIMP DM annihilation in the Sun, electron or muon tracks along the direction of the Sun, dominantly from charged-current interaction have been studied [26–28], while similar proton recoils from neutral-current (NC) interaction were not investigated due to the relatively larger scattering angle. Nonetheless, as discussed earlier, boosted DM only scatter off SM via NC-type of interactions. Unless very light mediator is involved, a proton recoil signal would typically have much larger rate compared to a single-ring electron signal and is the primary search channel.

We would also like to reiterate other broad motivations for investigating boosted DM signals, as have been discussed in more detail in [5]. First, annihilation into dark sector particles, as occurs in both our two-component DM or semi-annihilating DM examples, evades strong constraints from DM detection experiments in a natural way, while still

allowing for a thermal freeze-out origin of DM. This is complementary to other variations such as models where DM annihilates into dark radiation or to dark states that decay back to the SM, as discussed in Ref. [29–34] for example. Second, these studies of boosted DM demonstrate how the expected phenomenology, and search strategies for a non-minimal DM sector or single component DM with non-minimal stabilization symmetry can be very different from those for conventional single-component Z_2 DM models. Non-minimality is already a familiar fact for the SM matter sector, where protons and electrons stabilized by separate B - and L -number symmetries. It is quite reasonable that DM sector also has beyond the simplest, minimal content. Earlier work on phenomenology from multi-component DM sector include [35–42], on DM sector with non-minimal stabilization symmetry such as Z_3 include [43–47].

The organization of this paper is as follows. In Sec. II we present examples of semi-annihilating DM and two-component DM models in detail, starting from effective operators for DM interactions with quarks and then developing UV completions. In Sec. III we analyze the various DM processes in the Sun including capture, evaporation, annihilation and rescattering, and eventually determine the flux of boosted DM incident on the Earth. Then in Sec. IV we discuss the detection rate of boosted DM at large volume detectors using the proton Cherenkov ring signal. We assess the discovery prospects at SK and its upgrade HK in Sec. V, commenting on relevant constraints on these particular models. Finally we conclude in Sec. VI with discussions of other possibilities. We include more details in the appendices.

II. MODELS

In this section, we present the concrete examples of the two classes of models on which we focus in this paper, semi-annihilating DM models and two-component DM model. These models are related by the fact that DM particles appear in the final state of annihilation processes, opening the possibility that the final state DM is boosted.

In the most concise version of a semi-annihilation model, there is only one specie of dark matter particle. There is thus a direct relation between solar capture rate and the detection rate at neutrino detectors such as SK. Furthermore, there is only a small yet generic range ($\gamma = 1 - 1.25$) for the boost factor of final state DM particle in the semi-annihilation models, which as we will see, interestingly falls in the sweet-spot for detecting a proton track signal.

In the two component DM models, different particles are involved in solar capture and SK detection processes, leading to more general kinematic possibilities at the expense of a larger number of parameters. The mass ratio between two DM particles in the two component DM model controls the boost factor of DM particle in the final states, which will be important in determining the signatures at SK and other detectors. We demonstrate below that the preferred mass ratio range for obtaining an observable signal ranges approximately from 1.1 to 2.2. If the mass ratio is too low, then the DM particle in the final state is not boosted and the recoiling proton does not generate Cherenkov light. If the ratio is too high, the interaction between the boosted DM and the protons in the detector is dominantly inelastic and the rate of single Cherenkov ring events, which are most easily distinguished from background events, is suppressed.

In studying these signatures, however, we parameterize both classes of models in terms of a small number of phenomenological parameters which are relatively insensitive to the details of the complete models we present in this section. These models serve as an important proof-of-principle that complete models can be constructed, as well as a motivation for the forms of interactions that we will introduce. Other models yielding similar signatures may be possible, but the study of these other possibilities is beyond the scope of this work.

A. Semi-annihilating DM models

There are several simple ways to construct a semi-annihilating DM model. For example, introducing either a Z_3 symmetry for DM particles or a near-degenerate spectrum in the dark sector along with a stabilizing symmetry can lead to such behavior. In this paper, we focus only on the simplest version of semi-annihilation model where we assume a single DM particle χ charged under a Z_3 symmetry, which protects χ from decay. The DM χ can either be a scalar or a fermion. Taking the approximation that DM particles in the initial states are almost at rest, the boost factor of DM particle in the final state is

$$\gamma_\chi = \frac{(5m_\chi^2 - m_\psi^2)}{4m_\chi^2}, \quad (3)$$

where m_ψ is the mass of the lighter unstable particle ψ in the final state of semi-annihilation process, ψ can either be a scalar or fermion. γ_χ ranges from 1 to 1.25. For simplicity, we assume $m_\psi \ll m_\chi$. Then the boost factor is near maximal. Lowering γ_χ by a small amount does not induce a big change in our conclusion. Further we will show in

later content that a boost factor around 1.25 is within the preferred region for detection at SK and similar Cherenkov light detectors.

The unstable particle in the final state of semi-annihilation ψ is neutral under Z_3 , and can decay to SM particles. The decay products and decay lifetime are highly model dependent. On the other hand, only the mass of ψ can affect the boost factor of the DM particle in the final state. Thus focusing on the boosted DM particle provides the most model-independent way to study this class of DM models.

We emphasize that in the simplest version of semi-annihilation model, DM particles in the initial and final states are the same. The scattering cross-section between DM particle and nucleon, $\sigma_{\chi,N}$, determines both DM solar capture rate and the interaction probability when a DM particle passes through the region of SK. This reduces the number of effective free parameters in the model.

It is well-known that there are two classes of DM-nucleus interactions: spin-dependent or spin-independent (SI). In traditional DD experiments, if the scattering process is governed by SI operators, DM can scatter with the whole nucleus coherently due to the low velocity of DM in the local halo, leading to strong sensitivity to DM interactions. On the other hand, if the scattering process is dominantly SD, the constraints from DD are much weaker.

The situation is very different in the process we are considering. Since the Sun is mainly constituted of hydrogen, the coherent enhancement is absent for SI operators. There is also a fraction of helium in the Sun, but the coherence effect is not large enough to make order of magnitude difference on the capture rate for SI and SD operators. On the other hand, in SK we focus on the DM-nucleon scattering process which can kick a proton out from the nuclei in water and generate a Cherenkov ring. The momentum transfer in such process is generically larger than few hundreds of MeV, which makes the coherence effects negligible. Unlike the ordinary DM DD experiments, the reach for SI operators will be comparable to that of SD operators. Since SD operators are much less constrained than SI operators in DD experiments, we choose our benchmark operators to be SD. The reach limits for SI operators should not be dramatically different.

Furthermore, the velocity of DM particles captured and thermalized in the Sun is much smaller than that of a boosted DM particle from DM annihilation. In the non-relativistic limit, $\sigma_{\chi,N}$ may have non-trivial dependence on relative velocity or momentum transfer. In some cases, such dependence may induce dramatic enhancements to the cross section for interactions between boosted DM and nucleons compared to that for collisions between the thermal non-relativistic component of DM and nuclei. In the rest of this section, we are going to focus on two scenarios. In the first, we assume that the leading term for $\sigma_{\chi,N}$ in non-relativistic limit has v^0 dependence, while in the second, we assume v^2 dependence.

1. v^0 operator

In this section, we focus on the semi-annihilating model where the scattering cross section between DM and nucleons has v^0 dependence in non-relativistic limit. To make our discussion more concrete, we take a benchmark operator for detailed calculations. A typical operators of this kind is

$$\mathcal{O}_{SD,v^0} = \frac{1}{M^2} \bar{\chi} \gamma^5 \gamma^\mu \chi \bar{q} \gamma_\mu \gamma^5 q, \quad (4)$$

where χ is the DM particle, assumed to be a Dirac fermion. χ_L and χ_R are the chiral components of χ , such that $\chi = (\chi_L, \chi_R^\dagger)$. Thus χ_L and χ_R have Z_3 charges as $e^{i2\pi/3}$ and $e^{i4\pi/3}$ respectively. The scattering induced by this operator is SD.

Now let us first UV complete the effective operator in eq. (4). This can be achieved by introducing a gauge boson Z' which couples axially to both χ and quarks. The Lagrangian can be written as

$$\mathcal{L} \supset -i\chi_L^\dagger \bar{\sigma}^\mu D_\mu \chi_L - i\chi_R^\dagger \bar{\sigma}^\mu D_\mu \chi_R - iQ^\dagger \bar{\sigma}^\mu D_\mu Q - iu^{c\dagger} \bar{\sigma}^\mu D_\mu u^c - id^{c\dagger} \bar{\sigma}^\mu D_\mu d^c \quad (5)$$

We assume the axial gauge group is spontaneously broken, and Z' is massive. The covariant derivatives in eq. (5) includes canonical couplings of Z' . To introduce axial current coupling, we require Q has same charge as u^c and d^c . χ_L and χ_R also share the same $U(1)'$ charge. The anomaly of this $U(1)'$ can be canceled by introducing extra charged particles at higher mass scale. For simplicity, we take the charges to have the form $q_Q = q_{u^c} = q_{d^c} = q_{SM}$ and $q_{\chi_L} = q_{\chi_R} = q_{DM}$. After integrating out Z' , the effective operator eq. (4) is generated with

$$\frac{1}{M^2} = \frac{g_{Z'}^2 q_{SM} q_{DM}}{m_{Z'}^2}. \quad (6)$$

Noted that integrating out the longitudinal part of Z' will induce another sub-leading operator, which we drop in the following discussion.

To place χ in the context of a semi-annihilating DM model, we introduce another lighter Dirac fermion ψ , which is neutral under Z_3 . We take its chiral components to be (ψ_L, ψ_R^\dagger) . Both ψ_L and ψ_R have the same $U(1)'$ charge as χ_L and χ_R . The Lagrangian can be written as

$$\mathcal{L} \supset y_1 \phi \chi_L \chi_R + y_2 \phi (\psi_L \psi_L + \psi_R \psi_R) + y_3 \phi \psi_L \psi_R + \frac{\lambda_1}{m^2} (\chi_L \chi_L) (\chi_R^\dagger \psi_R^\dagger) + \frac{\lambda_2}{m^2} (\chi_R \chi_R) (\chi_L^\dagger \psi_L^\dagger), \quad (7)$$

where ϕ is a scalar field charged under $U(1)'$. Its condensation breaks the $U(1)'$ and gives mass to both χ and ψ . The non-renormalizable terms can be easily UV completed by introducing another complex scalar field which is also charged under Z_3 . Its detailed properties are not important, so we do not write down the full UV completion explicitly. ψ is not protected by Z_3 symmetry. After $U(1)'$ is broken, it can decay to SM particles.

Next, we determine the cross-section for DM-nucleon interactions using the effective operator in eq. (4). Note that the non-relativistic scattering processes relevant to the signal considered are in the Sun, which is made up nearly entirely of hydrogen, which has a single proton as a nucleus, and of helium, which has 0 spin. Boosted DM will resolve the individual nucleons in heavier nuclei. Therefore, scattering off of nucleons is the only relevant process for the process chain shown in Fig. 1. The velocity of the boosted DM particle in the final state is nearly 0.6, so the momentum transfer cannot be larger than roughly 2 GeV. Therefore effective operator is always a reasonable approximation to describe such scattering as long as the mass of Z' is larger than few GeV. Nevertheless, we present the full form of the cross-section in terms of the phenomenological parameters introduced below in Appendix B. For simplicity, we take the approximation of $m_\chi \gg m_p$. If $m_\chi \lesssim 5$ GeV, most DM particles captured by the Sun will evaporate after thermal equilibrium is reached. More details on evaporation will be presented in later sections.

There is, however, a form factor correction relative to the proton scattering in DD experiments, as well as possible isospin dependence of the interactions. For all SD interactions, the relevant form factor is the axial form factor, which is known from scattering neutrinos off of protons and neutrons at these energies. The Q^2 dependence, as in [48], is thus

$$F(Q^2) = \frac{1}{(1 + \frac{Q^2}{M_A^2})^2}, \quad (8)$$

where M_A is an empirical scale measured to be 1.03 GeV, and $\sigma_{\text{DD}} \propto F(Q^2)^2$.

The total scattering cross-section under the above assumptions is

$$\sigma_{\chi, N} = \frac{3m_N^2 m_\chi^2}{\pi M^4 (m_\chi + m_N)^2} \left(\sum_q \Delta_q \right)^2 F(Q^2), \quad (9)$$

where $\sigma_{\chi, N}$ is the scattering cross-section per nucleon N , which is the same for protons and neutrons assuming isospin-preserving interactions. Neglecting corrections from the form factor, if we take $\sigma_{\text{DD}, p}$ to be 10^{-38}cm^2 , M is around 400 GeV, where we assumed isospin-respecting coupling to quarks, and the spin factors Δ_q are presented in Appendix A.

For a low suppression scale, one may also worry about the collider constraints. Leptons may be neutral under $U(1)'$, which helps in evading constraints from LEP. Such leptophobic Z' is well motivated and has been studied extensively. Thus the only relevant constraints are monojet and dijet resonance searches. The monojet constraints can be alleviated by reducing the $U(1)'$ gauge coupling or $U(1)'$ charge. To keep the effective suppression scale M fixed, the mass of Z' is also lowered. If the mediator mass is smaller than twice DM mass, then the pair of DM are produced through off-shell Z' in the collider. Both smaller coupling constants and Z' being off-shell helps to avoid the monojet constraints. For example, if we take $m_{Z'}$ to be 10 GeV while fixing the $U(1)'$ charge to be order one, to get $M = 400$ GeV, the $U(1)'$ gauge coupling needs to be around 0.025. With a jet cut at 250 GeV in 8 TeV LHC, the monojet cross section is much less than 0.1 pb, well below the current constraints from such searches. Dijet resonance searches are very weak when $m_{Z'}$ is very light. For a review, see [49]. Below 200 GeV mass (with sensitivity down to ~ 80 GeV), the best constraint is from UA2, and the gauge coupling is only constrained to be smaller than 1.7 at 200 GeV. Our model is thus unconstrained from such searches as well at the moment.

The conventional DM annihilation channel into the SM quarks must also be present due to crossing symmetry. If this channel dominates over semi-annihilation, our signal is suppressed and indirect detection may place strong bounds on the model. We therefore estimate the conventional DM annihilation cross section in our UV completion, i.e. eq. (5). Let us focus on the first generation of quarks and take massless quark approximation. The total annihilation cross section can be written as

$$\langle \sigma_{\text{ann}}^{\chi\chi \rightarrow \bar{q}q} v \rangle = \frac{g^4 q_{SM}^2 q_{DM}^2}{2\pi} \frac{m_\chi^2 v^2}{(4m_\chi^2 - m_{Z'}^2)^2 + \Gamma_{Z'}^2 m_{Z'}^2}. \quad (10)$$

The conventional annihilation cross section is thus p-wave suppressed. A generic choice of parameters, this process is thus subdominant to semi-annihilation and is unconstrained by indirect detection experiments.

One subtlety in this model should be addressed at this point. Since we assign axial charge to first generation quarks, the SM Yukawa interaction is not invariant under $U(1)'$ gauge symmetry. Setting $q_{SM} = q_{DM}$, then one can introduce a non-renormalizable term to generate an up quark mass (with a similar term for the generating the down quark mass),

$$\mathcal{L}_{\text{Yukawa}} = \frac{y_u}{M_Y} \phi H Q u^c. \quad (11)$$

Small masses for first generation quarks can be generated after both ϕ (which acts like a ‘‘flavon’’ here) and H get VEVs. There could also be conventional DM annihilation process through s-channel ϕ boson exchange. However, it is p-wave and light quark mass suppressed. We do not consider this channel further.

The specific structure of this model is largely irrelevant for the signal studied in this paper. We thus use a phenomenological parametrization to describe the interactions of DM in this model. The relevant interactions are with nuclear matter. At low energies, the cross-section for this interaction corresponds to the DD interaction in eq. (9). We can thus eliminate the charges and couplings from our description in favor of the DD cross-section, such that the parameters of the model are taken to be the masses of χ and Z' , as well as the DD cross-section σ_{DD} that would be seen in conventional non-relativistic DD experiments. We emphasize once more that there is only a mild dependence on the Z' mass for the parts of parameter space considered provided that $m_{Z'}$ is larger than a few GeV. The full differential cross-section for DM-nucleon interactions in terms of these parameters is presented in Appendix B and is used in all further calculations for such models.

2. v^2 operator

In this section, we study a semi-annihilating DM model with an operator inducing a v^2 dependence in the non-relativistic scattering between DM and nucleon. Our benchmark operator is written as

$$\frac{i}{M^2} (\chi^\dagger \partial_\mu \chi - \partial_\mu \chi^\dagger \chi) \bar{q} \gamma^5 \gamma^\mu q, \quad (12)$$

where χ is the DM particle and is a scalar field in this model. As with the operator in eq. (4), this is a dimension 6 operator that can be generated by integrating out a massive gauge boson under which both SM quark and DM particle are charged. One can UV complete this operator via the interaction Lagrangian

$$\mathcal{L} \supset D_\mu \chi^\dagger D^\mu \chi - i Q^\dagger \bar{\sigma}^\mu D_\mu Q - i u^{c\dagger} \bar{\sigma}^\mu D_\mu u^c - i d^{c\dagger} \bar{\sigma}^\mu D_\mu d^c. \quad (13)$$

We once more introduce a spontaneously broken $U(1)'$ gauge group, with massive gauge boson Z' . Thus the covariant derivatives in eq. (13) also include canonical couplings of Z' . To introduce axial current coupling while forbidding vector current coupling to quarks, we require left-handed quark Q to have same charge as u^c and d^c . The anomaly of this $U(1)'$ can be canceled by introducing extra charged particles at higher mass scale. For simplicity, we take $q_Q = q_{u^c} = q_{d^c} = q_{SM}$. After integrating out Z' , the effective operator eq. (12) is generated with

$$\frac{1}{M^2} = \frac{g_{Z'}^2 q_\chi q_{SM}}{m_{Z'}^2}. \quad (14)$$

The additional contribution induced by integrating out longitudinal component of Z' is again negligible.

To add semi-annihilation processes to this scenario, we assume scalar χ is stabilized by Z_3 symmetry. We introduce another lighter scalar ϕ , which we take to be neutral under Z_3 . A proper choice of $U(1)'$ charge for ϕ leads to the following interaction Lagrangian:

$$\mathcal{L} \supset m_\chi^2 \chi^\dagger \chi + \frac{m_\phi^2}{2} \phi^\dagger \phi + \lambda_1 (\chi^3 \phi^\dagger + \chi^{\dagger 3} \phi). \quad (15)$$

It is assumed that semi-annihilation dominates over conventional annihilation and thus determines the DM relic abundance. ϕ can decay promptly after $U(1)'$ is broken. Its decay products are model dependent and we do not study them further.

We proceed to calculate the DM-nucleon cross-section in this model. Again, we note that in semi-annihilation scenario, it is always a reasonable approximation to use the effective operator to calculate the cross-section for

mediator masses of at least a few GeV. Z' couples axially to quarks in this model as well and one can thus apply the same form factor as eq. (8) to estimate the elastic scattering between DM and nucleon.

The total cross-section can be approximated by

$$\sigma_{\chi,N} = \frac{m_N^2}{2\pi M^4} v^2 \left(\sum_q \Delta_q \right)^2 F(Q^2) \quad (16)$$

which is suppressed by DM velocity squared, as we expect. As a benchmark, if we take $\sigma_{\chi,N}$ to be 10^{-40} cm² and $v = 10^{-3}$, the suppression scale is estimated to be 30 GeV.

As with the v^0 operator as discussed in previous section, one can lower $m_{Z'}$ at the same time as the coupling constant while keeping the suppression scale M fixed. For example, if we take $m_{Z'}$ to be 5 GeV, the coupling constant should be around 0.2. Taking DM mass as 10 GeV and assuming Z' couples universally to all quarks, the monojet cross section at 8 TeV LHC with 250 GeV jet cut is only around 0.04 pb. Further, a coupling constant as small as 0.2 is also safe from dijet resonance constraints. Thus collider searches are not yet sensitive to this UV model.

As in v^0 scenario, one may be worried whether such low suppression scale induces a large cross section for conventional DM annihilation into SM quarks. To estimate the ordinary DM annihilation cross section, we focus on the first generation of quarks and work in the massless quark limit. The annihilation cross section can be written as

$$\langle \sigma_{\text{ann}}^{\chi\chi \rightarrow q\bar{q}v} \rangle = \frac{g^4 q_{DM}^2 q_{SM}^2}{\pi} \frac{2m_\chi^2 v^2}{(4m_\chi^2 - m_{Z'}^2)^2 + \Gamma_{Z'}^2 m_{Z'}^2} \quad (17)$$

This annihilation cross section is again p-wave suppressed. A generic choice of parameters gives a small annihilation cross section for this standard DM annihilation channel. Thus the boosted DM flux from the Sun will not be reduced by the existence this channel. Also this model is safe from indirect detection constraints.

Finally, as in the v^0 case, this model can be parameterized by the phenomenological parameters m_χ , $m_{Z'}$, and σ_{DD} , with only mild dependence on $m_{Z'}$. The full differential cross-section for DM-nucleon interactions in terms of these parameters is presented in Appendix B and is in future calculations.

B. Two-Component DM models

In a two-component DM model, there are at least two components of stable particles, ψ_A and ψ_B . We assume throughout that ψ_A is the dominant component of DM. The DD constraints to B can be negligible if B has a sufficiently suppressed relic abundance. In these models, the solar capture rate is controlled by $\sigma_{A,p}$, the scattering cross-section of A with protons in the Sun, while the boosted DM detection rate is controlled by $\sigma_{B,p}$, the scattering cross-section of B to knock out protons in the target. As in the the semi-annihilation scenario discussed above, there could still be subtleties coming from non-trivial velocity dependence. However, since we have already decoupled $\sigma_{A,p}$ and $\sigma_{B,p}$, the cross section enhancement when DM is boosted can be partially mimicked by imposing $\sigma_{B,p} \gg \sigma_{A,p}$. Thus in the following discussion we will focus on a benchmark model where both $\sigma_{A,p}$ and $\sigma_{B,p}$ have v^0 dependence at leading order.

We consider two Majorana fermion DM ψ_A, ψ_B , where $m_A > m_B$. ψ_A is the major DM, while ψ_B is a subdominant component. This is natural if ψ_B has a larger thermal annihilation cross section. For both ψ_A and ψ_B , we consider the same type of SD DM-nucleon scattering operator as used in the v^0 dependent semi-annihilation model, i.e. eq. (4):

$$\mathcal{O}_{SD,v^0} = \frac{1}{M^2} \bar{\chi} \gamma^5 \gamma^\mu \chi \bar{q} \gamma_\mu \gamma^5 q.$$

Note that the fact that ψ_A and ψ_B are Majorana fermions automatically eliminates the possibility of operator $\bar{\chi} \gamma^\mu \chi$. Because of this, the quark-side of coupling does not have to be purely axial to easily evade direct detection constraints, as the following operator is SI but v^2 -suppressed [50, 51]:

$$\mathcal{O}_{SD,v^2} = \frac{1}{M^2} \bar{\chi} \gamma^5 \gamma^\mu \chi \bar{q} \gamma_\mu q. \quad (18)$$

For simplicity, we focus on the \mathcal{O}_{SD,v^0} operator for the two-component model.

Such an operator can be generated by a similar UV completion to that of Lagrangian in eq. (5). Both ψ_A, ψ_B are charged under $U(1)'$ and have Majorana masses which may result from $U(1)'$ symmetry breaking. The relevant model Lagrangian can be written in 4-component notation as

$$\begin{aligned} \mathcal{L} \supset & \bar{\psi}_A (i\partial_\mu + \frac{1}{2} q_{AG} Z'_\mu) \gamma^\mu \gamma_5 \psi_A + \bar{\psi}_B (i\partial_\mu + \frac{1}{2} q_{BG} Z'_\mu) \gamma^\mu \gamma_5 \psi_B + \bar{\psi}_q [i\partial_\mu + (g_{q,V} + g_{q,A} \gamma_5) Z'_\mu] \psi_q \\ & - \frac{1}{2} m_A \bar{\psi}_A \psi_A - \frac{1}{2} m_B \bar{\psi}_B \psi_B - \frac{1}{2} m_{Z'}^2 Z'^\mu Z'_\mu, \end{aligned} \quad (19)$$

where g is the $U(1)'$ gauge coupling, q_A , q_B are the $U(1)'$ charges of ψ_A , ψ_B , and ψ_q are the SM quarks. As discussed earlier, for simplicity, we take $g_{q,V} \rightarrow 0$, which can be realized in a UV construction in the same way as in the semi-annihilation models. In order for ψ_B to annihilate away efficiently leaving a suppressed relic abundance relative to that of ψ_A , we further assume $m_B > m_{Z'}$, which also helps alleviate the potential constraint from monojet searches at colliders, as discussed earlier.

Both the cross sections of ψ_A and ψ_B for scattering off nucleons are relevant to determining the size of the signal from this model. Boosted ψ_B from ψ_A annihilation can be highly relativistic during subsequent scatterings if $m_A/m_B \gg 1$. In cases where the ψ_B is sufficiently boosted, it is imperative to consider the full form of the scattering cross-section for this model as presented in Appendix B. Note that this model has additional phenomenological parameters as the masses and effective DD cross-sections for ψ_A and ψ_B are different in general. We parameterize this model in terms of the mass and cross-section for ψ_A , as well as the ratio of the masses of the Z' and ψ_B to ψ_A and the ratio of the ψ_B effective direct detection cross-section to that of ψ_A .

In the non-relativistic limit, the total scattering cross section with nucleons for $\chi = \psi_A, \psi_B$ is

$$\sigma_{\chi,N}^{v \rightarrow 0} = \frac{3g_\chi^2 g_q^2 m_\chi^2 m_N^2}{\pi m_{Z'}^4 (m_\chi + m_N)^2} \left(\sum_q \Delta_q \right)^2. \quad (20)$$

Various collider constraints are evaded when the mediator is light and the couplings to quarks are not too large.

The thermal relic abundance of ψ_A is dominantly determined by $\psi_A \psi_A \rightarrow \psi_B \psi_B$ annihilation via A' exchange, with the cross section given as follows:

$$\langle \sigma_{\text{ann}}^{AA \rightarrow BB} v \rangle = \frac{g_A^2 g_B^2}{12\pi} \frac{\sqrt{m_A^2 - m_B^2}}{m_A (4m_A^2 - m_{Z'}^2)^2 + \Gamma_{Z'}^2 m_{Z'}^2} \left[3m_B^2 + v^2 \cdot \frac{24m_B^2 m_A^4 + m_A^2 m_{Z'}^2 (-6m_B^2 + m_{Z'}^2) - m_{Z'}^4 m_B^2}{m_{Z'}^4} \right], \quad (21)$$

where we can see that the s-wave part is suppressed by m_B^2/m_A^2 , therefore we have kept the potentially non-negligible p-wave contribution as well. The relic abundance of ψ_A takes the standard form expected for WIMP DM

$$\Omega_A \simeq 0.2 \left(\frac{3 \times 10^{-26} \text{ cm}^3/\text{s}}{\langle \sigma_{AA \rightarrow BB} v \rangle} \right). \quad (22)$$

We present a couple of benchmark parameter points at which the observed DM abundance corresponds to the thermal relic abundance of ψ_A :

$$\begin{aligned} & \{m_A = 150 \text{ GeV}, m_B = 100 \text{ GeV}, m_{Z'} = 50 \text{ GeV}, g_A = 3 \cdot 10^{-3}, g_B = 0.3\} \\ & \{m_A = 60 \text{ GeV}, m_B = 30 \text{ GeV}, m_{Z'} = 20 \text{ GeV}, g_A = 0.08, g_B = 0.4\}. \end{aligned}$$

The thermal annihilation cross section of ψ_B by annihilating into Z' is given by

$$\langle \sigma_{\text{ann}}^{BB \rightarrow Z' Z'} v \rangle|_{v \rightarrow 0} = \frac{g_B^4}{2\pi} \frac{m_B^2 - m_{Z'}^2}{(m_{Z'}^2 - 2m_B^2)^2} \sqrt{1 - \frac{m_{Z'}^2}{m_B^2}}. \quad (23)$$

The annihilation of ψ_A, ψ_B into SM quarks is helicity suppressed and/or can be suppressed by assuming $g_q < g_{A,B}$. Computing the thermal relic abundance of ψ_B is more complicated than that for ψ_A as $\psi_A \psi_A \rightarrow \psi_B \psi_B$ can be important as well. Simple analytic estimates can be obtained depending on parameter region. These were discussed intensively in [5] where the concept of ‘‘balanced freeze-out’’ was introduced for the region where ψ_B freezes out much later than ψ_A . We do not repeat the discussion here. Considering that there can be other channels beyond the minimal model which can sufficiently deplete the abundance of ψ_B , in this paper we do not elaborate the relic abundance calculation of ψ_B and related direct/indirect detection bounds.

III. BOOSTED DM FLUX FROM THE SUN

In this section, we determine the flux Φ of DM particles from the Sun. The flux can be written as¹

$$\Phi = \frac{AN^2}{4\pi \text{AU}^2}, \quad (24)$$

¹ It is common in the literature to break the full annihilation rate into the annihilation rate per DM pair A and the number of DM particles N as written here.

where AN^2 is related to the annihilation rate of DM captured in the Sun, Γ_A by $\Gamma_A = \frac{1}{2}AN^2$, and AU is an astronomical unit, the distance from the Sun to the Earth. The annihilation rate, in turn, is effectively given by the capture rate C in the part of parameter space we are interested in. The primary goals of this section are thus to calculate C and to determine the region of parameter space where $AN^2 = C$.

A. DM Capture Rate by the Sun

For any model in which DM scatters off of nucleons, the capture rate can be written as

$$C = \int dV du \sigma_{\chi,H}(w \rightarrow v)|_{v < v_{\text{esc}}} \frac{w^2}{u} n_{\chi}(r) n_H(r) f(u), \quad (25)$$

where dV is a volume element of the Sun, $\sigma_{\chi,H}(w \rightarrow v)|_{v < v_{\text{esc}}}$ is the total cross-section for DM to scatter to a velocity below the escape velocity, u is the velocity of the DM particle if it were far away from the Sun, w is the actual velocity of the DM $\sqrt{u^2 + v_{\text{esc}}^2}$, $n_{\chi}(r)$ is the number density profile of DM particles near the Sun, $n_H(r)$ is the number density of hydrogen nuclei [58], and $f(u)$ is the local DM velocity distribution. In the Appendix C, we determine the details of the pieces of (25).

Given these quantities, the integral over the velocity and volume can be performed numerically. The case where DM has velocity suppressed interactions with matter actually has a roughly factor of 20 enhancement of its capture rate for a fixed effective DD cross-section. Since the DM has fallen into the Sun's gravitational potential, it scatters at higher velocity in the Sun than it does in a detector near the surface of the Earth. In the analysis below, we determine the capture rate for each parameter point, but in order to gain an intuition of the orders of magnitude involved, we present the numerical results for the parameter point with $m_{\chi} = 100$ GeV and $\sigma_{\text{DD}} = 10^{-42}$ cm². Here, χ is the DM particle in semi-annihilating models and ψ_A in two component models. The mediator mass has negligible effect since effective operator is an excellent approximation. We thus leave it to be specified later. In the cases where the elastic DM-nucleon scattering cross-section is not v suppressed, we find

$$C = 2.0 \times 10^{20} \text{ sec}^{-1}. \quad (26)$$

If the cross-section is suppressed by v^2 , then we find

$$C = 5.1 \times 10^{21} \text{ sec}^{-1}. \quad (27)$$

The dependence on the mass and cross-section can roughly be parameterized in the large DM mass limit as

$$C(m_{\chi}, \sigma_{\chi,N}) \approx C(100 \text{ GeV}, 10^{-42} \text{ cm}^2) \left(\frac{\sigma_{\text{DD}}}{10^{-42} \text{ cm}^2} \right) \left(\frac{100 \text{ GeV}}{m_{\chi}} \right)^2 \quad (28)$$

for $m_{\chi} \gg 1$ GeV and for $C(100 \text{ GeV}, 10^{-42} \text{ cm}^2)$ given by eq.(26) or eq.(27), depending on the model being considered.

B. Capture-loss Equilibrium in the Sun

After the formation of the Sun, DM begins to be captured by the process described in Sec. IIIA. On the other hand, there are two dominant processes that reduce the amount of DM in the Sun, annihilation and evaporation. After a long time of accumulation of DM particles in the Sun, an equilibrium state may be reached such that

$$AN^2 = C - EN, \quad (29)$$

where E is the rate per DM particle at which DM evaporates from the Sun and N is the number of DM particles captured in the Sun. In a regime where $E \approx 0$, which we demonstrate below is generic, the annihilation rate is given by the capture rate. After DM particle is captured by the Sun, it will soon reach thermal equilibrium with the Sun. The DM distribution can be characterized by a thermal radius, which is given by [52]

$$r_{\text{th}} = \left(\frac{3T}{2\pi m_{\chi} G_N \rho_c} \right)^{1/2} = 0.01 R_{\odot} \left(\frac{T}{1.2 \text{ keV}} \right)^{1/2} \left(\frac{100 \text{ GeV}}{m_{\chi}} \right)^{1/2}, \quad (30)$$

where G_N is Newton's constant, ρ_c is the core density of the Sun, R_\odot is the solar radius, and $T = 1.2\text{keV}$ is the Sun's core temperature. Thus one can calculate the time needed for the DM to reach an equilibrium between capture and annihilation, $\tau_{\text{eq}} = 1/\sqrt{C \cdot A}$, and compare it with the age of the Sun [52]:

$$\frac{t_\odot}{\tau_{\text{eq}}} = 10^3 \left(\frac{C}{10^{25} \text{ sec}^{-1}} \right)^{1/2} \left(\frac{\langle \sigma_{\text{ann}} v \rangle}{3 \times 10^{-26} \text{ cm}^3 \text{ sec}^{-1}} \right)^{1/2} \left(\frac{0.01 R_\odot}{r_{\text{th}}} \right)^{3/2}. \quad (31)$$

As long as $\frac{t_\odot}{\tau_{\text{eq}}} > 1$, equilibrium is reached by the present day. In this case, the DM flux will be independent of the DM annihilation cross section. For an annihilation cross-section close to the value which gives correct thermal relic abundance, the Sun will have reached capture–annihilation equilibrium by now. Even in models where the relic abundance is determined by p-wave annihilation, as long as there is a non-negligible s-wave component, equilibrium can still be reached in the Sun. For the models we will be discussing, the equilibrium condition will always be satisfied for the region of interest in parameter space.

A full treatment of DM evaporation in the case without velocity suppression is found in [53]. Here, we apply simple arguments to estimate the evaporation rate in the v^2 case. Evaporation occurs when an energetic nucleus on the tail of the (local) solar Boltzmann distribution collides with a slower DM particle and imparts a velocity larger than the escape velocity. Typically, the DM velocity will be much less than the nucleus velocity and we approximate it to be at rest. Under this approximation, the evaporation rate per unit volume in the Sun is given by

$$\frac{dE}{dV} = \int_{(m_p + m_\chi)v_{\text{esc}}/2m_p}^{\infty} du \sigma_{\chi, H}(0 \rightarrow v)|_{v > v_{\text{esc}}} n_\chi(r) n_H(r) u_H f(u_H), \quad (32)$$

where μ is the proton-DM reduced mass, $\sigma(0 \rightarrow v)|_{v > v_{\text{esc}}}$ is the cross-section for a proton of velocity u_H to impart a velocity above the escape velocity, n_χ is the local captured DM number density, n_H the local hydrogen density, and f is the velocity distribution of hydrogen in the Sun. The individual pieces are again presented in the Appendix.C. The resulting rate can be integrated. It is found that in this approximation, the term EN in (29) can be neglected for DM masses above 4 GeV. To be conservative given the approximations we make in this determination, we consider only DM masses larger than 5 GeV and neglect evaporation effects entirely in models with either v^0 and v^2 behavior.

C. Rescattering in the Sun

As the DM particle travels from the center of the Sun where it is produced in annihilation processes, it may rescatter off of solar nucleons and lose velocity. This alters the effective detection cross-section. In this paper, we use a conservative estimate for this effect. At larger cross-sections, there is a larger flux and thus a larger detection rate for boosted DM. On the other hand, there is an increasingly significant loss of energy as the DM exits the Sun. Since only DM with sufficiently large energy can scatter protons to momenta above the Cherenkov threshold, most of the scattered DM will not be detected if the energy loss is too great. Since both the collision rate with the detector target and the energy loss as the DM exits the Sun scale as σ_{DD}^2 , there is a detailed interplay between these effects. In this paper, we use the very conservative approximation that there is no sensitivity to models for which the mean energy of the DM escaped from the Sun is insufficient to provide a large enough Cherenkov detection rate. For larger cross-sections, fluctuations become important in determining the mean detection cross-section and there is likely sensitivity up to significantly higher cross-sections. A full determination of the mean detection cross-section is beyond the scope of this work.

The easiest way to compute this mean energy at the exit of the Sun is to work with Mandelstam variables. The probability that a DM particle interacts with Mandelstam variable t between t and $t + \Delta t$ while traveling between r and $r + dr$ from the center of the Sun is given by

$$dP = n_N(r) \frac{d\sigma}{dt}(s, t) dt dr, \quad (33)$$

where n_N is the number density of nucleons in the Sun². Recall that in the rest frame of the nucleons, $s = m_\chi^2 + m_N^2 + 2E_\chi m_N$, $t = 2m_N(E'_\chi - E_\chi)$. Therefore the change in energy is proportional to the change in Mandelstam s , which is given by

$$ds = tdP, \quad (34)$$

² In this calculation, we assume isospin is a valid approximation such that the scattering cross-sections off protons and neutrons are equal.

where Mandelstam t is the change in s for a collision parameterized by s and t . The mean energy is directly related to the mean value of s . To determine this mean, we solve

$$\int_{s_0}^{(s)} \frac{ds}{\int dt t(d\sigma/dt)} = \int_0^{R_\odot} dr n_N, \quad (35)$$

where $s_0 \equiv s(v_0)$, v_0 is the constant velocity of boosted DM particles coming right out of the annihilation process, and R_\odot is the solar radius. We do not claim to have sensitivity to models where the detection rate assuming DM particles incident on the detector with energy $\langle E_\chi \rangle$, which is directly related to $\langle s \rangle$, is too small to generate a sufficient number of signal events.

IV. DETECTION OF BOOSTED DM

A. Detection Mechanism for Signals

Neutrino experiments have well established techniques to detect recoils of energetic charged particles and thus can be repurposed to detect boosted DM particles via recoiling protons or electrons from DM-matter collision. The flux of boosted DM from the Sun is small compared to, say, that of the non-relativistic halo DM, so that a large volume detector is required. Two representatives of the largest active neutrino experiments (including their near future upgrade/extension) are IceCube/Deepcore/PINGU and SK/HK, which both use photo-multiplier tubes to detect Cherenkov light emitted during collisions with the target, and are potentially good candidates for boosted DM search. In this work, we focus on detecting Cherenkov protons instead of electrons from DM-matter collisions in the detector for the following reasons. First, as discussed earlier, for neutral-current type of interaction, scattering off protons has a larger rate than off electrons, except for models where very t -channel light mediator is involved [5]. Second, since the DM solar capture and rescatter rates are determined by the interaction between DM and nucleon, focusing on detecting proton signal at terrestrial experiments avoids introducing further model dependence in the DM-electron coupling.

The detection of proton recoils with momentum larger than around 2 GeV becomes problematic. In this regime, scattering becomes dominantly inelastic, leading to multi-rings events where the direction information is lost. Additionally, protons above 2 GeV in momentum have $> 50\%$ chance of producing pions as they travel through the detector, whose decay would give an extra electron-like ring. Finally, high energy elastic collisions may be more difficult to distinguish from muon recoils. As illustrated in [55], a Cherenkov ring of a recoiling proton is similar to that of a muon, but different from that of an electron which has blurred edge due to electromagnetic showers. Proton with sufficiently low momenta, less than a few GeV, are likely to be stopped within the detector due to its strong nuclear interactions, which causes the Cherenkov light emission from the recoiling proton to stop abruptly.

Due to the high energy threshold of IceCube, $\gtrsim 100$ GeV, it is not suitable to detect proton tracks in the a few GeV range. It is expected to be challenging even at its low energy extension, PINGU, with a few GeV as energy threshold[54]. In this paper we focus on determining the current limits and expected sensitivity for SK and HK in the single-ring proton track channel. Planned Liquid Argon Time-Projection Chamber experiments [24, 25] may also have sensitivity to proton recoils, though a study of this prospect is beyond the scope of this work.

For the reasons outlined above, when studying detection at SK and HK, we consider only recoiling proton momenta below 2 GeV as in [48]. Since the target material is water, which has a proton Cherenkov momentum of 1.07 GeV [48], there is only sensitivity for collisions which yield proton momenta at least this energetic.

Given the detection mechanism of searching for single Cherenkov rings from protons recoiling in water, we proceed to determine the effective cross-section of the detector and thereby obtain a estimated prediction for the number of expected signal events. This effective cross-section is determined both by the detection efficiency and acceptance and by the short distance scattering of the DM off of the target protons.

The typical recoil spectrum for protons above the Cherenkov threshold in the regions of parameter space considered in this work are similar to those of the background neutrinos. Given this similarity, we estimate the efficiency for detecting a DM particle to be given by 70%, which is the estimated efficiency for detecting atmospheric neutrinos via their single proton Cherenkov signature[55].

B. Background Reduction

The main background for our signal is from atmospheric neutrinos with NC interaction, which are nearly isotropic across the sky, and are the aim of the current searches for proton tracks at SK. Additional backgrounds to detection

include those that fake atmospheric neutrino neutral current scattering. As outlined in [55], these are primarily charged pion and muon scattering events from cosmic rays. Our signal can be distinguished from these backgrounds based on the following discriminators:

- *Angular information:*

Whereas the atmospheric neutrino background is nearly isotropic, the incoming boosted DM is coming nearly entirely from the direction of the Sun. The signal may be enhanced compared to the atmospheric neutrino and other backgrounds by cutting on the angle of the recoiling proton with respect to the Sun. For a boosted DM velocity of $v = 0.6$, assuming $m_{DM} \gg m_p$, the maximal angle between the incoming DM particle and the recoiled proton is 40.9° from the Sun. The angular resolution for the recoiling protons is 2.8° which is not an effective limiting factor in this case. We find that optimal s/\sqrt{b} is obtained for a cut on the proton recoil angle equal to the maximal recoil angle, $\theta_{\max} = 40.9^\circ$ for $v = 0.6$, around the Sun. This cut is optimized for the semi-annihilation DM spectrum, but is also approximately correct for two-component models in the region of parameter space to which our signal provides the greatest sensitivity. Beyond the requirement that the proton momentum fall between 1.07 GeV and 2 GeV, the signal acceptance for this cut is essentially 1. The background acceptance, on the other hand, is reduced by roughly ratio of the solid angle covered by the search to 4π ,

$$\eta_{\text{bkg}} \approx \frac{1}{2}(1 - \cos \theta_{\max}) = 0.122. \quad (36)$$

In our analysis, we use the most recent SK proton track data that is publicly available, which includes the runs of SK-I and SK-II, up to the year of 2009 [55]. With this data set, the background for our search is expected to be 49.6 in a lifetime of 2287.8 days over the full detector acceptance[55]. After applying the above acceptance cut, we would expect 6.05 background events.

In addition to eliminating a significant fraction of the background, an angular restriction would leave a large solid angle of side band, which would have no signal contamination. This region could be used to normalize the background and determine the detection efficiency, thereby eliminating most of the systematic uncertainty in the measurement of the signal region. For this reason, for our projected sensitivity calculations, we do not consider any systematic uncertainty, assuming that systematics would be subdominant to statistical uncertainties.

- *Absence of μ^\pm, e^\pm excesses:*

Neutrino background giving rise to proton tracks via NC interaction also lead to corresponding electron and muon signals via charged-current (CC) interaction with comparable rate. But those accompanying channels are absent for boosted DM as it only has NC type of interactions (as discussed the e^- signal from DM- e^- NC scattering can in generally suppressed, and has a different correlation to the proton signals compared to the case of neutrinos). This feature can only distinguish boosted DM signal from possibly proton tracks generated by the neutrinos from conventional WIMP DM annihilation in the Sun, which cannot be reduced by the directional information we discussed earlier.

Further background discrimination may be possible by exploiting finer information such as the energy or angular distribution of the scattered proton or by using multi-variate analysis. Here we take a basic cut-and-count approach based on simple variables.

We now turn to determining the short distance contribution to the effective detector cross-section. Given a DM particle that has exited the Sun with velocity v , corresponding to a Mandelstam s in collisions with protons at rest, the effective cross-section for the DM to be detected by the detector can be written as

$$\Sigma(s) = \epsilon(s)Z\sigma_{\text{Cher}}(s), \quad (37)$$

where ϵ is the detection efficiency, Z is the number of protons in the detector and σ_{Cher} is the cross-section for a DM particle to scatter a proton within the accepted momentum range as described above. Since the typical momentum transfer is very high, the proton binding to other nucleons is negligible. As discussed above, we take $\epsilon(s)$, which does in principle depend on the incident DM velocity, to be a constant 70%. For SK, we have $Z = 6.8 \times 10^{33}$, while for HK, the planned target has $Z = 1.7 \times 10^{35}$. The threshold s for Cherenkov detection is

$$s_{\min} = \sqrt{m_p(2m_\chi^2 E_{p,\text{Cher}} + 2m_p m_\chi^2 + m_p p_{\text{Cher}}^2)} + m_p E_{p,\text{Cher}} + m_\chi^2. \quad (38)$$

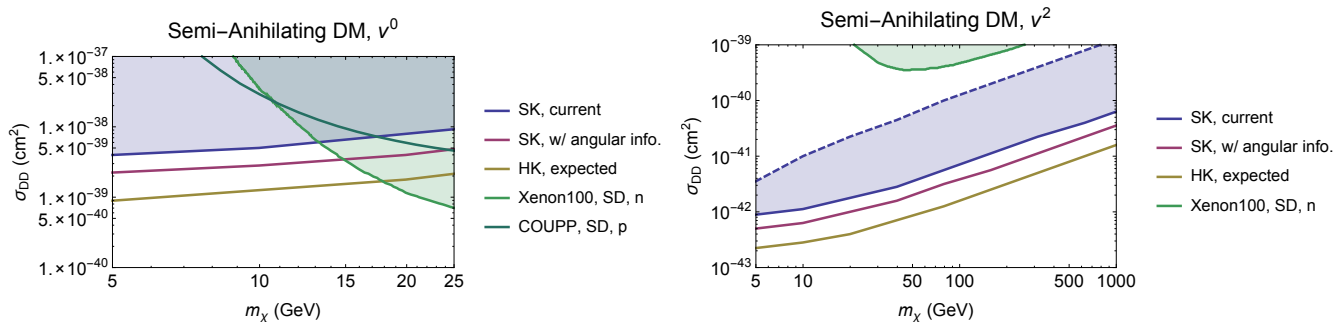


FIG. 2: Limits on the parameter space of semi-annihilation models with low energy operators O_{SD, v^0} (left) and O_{SD, v^2} (right). The lines labeled as “SK, current,” “SK, w/ angular info.,” and “HK, expected” are obtained as described in V. The dashed line right above the “SK, current” line indicates the cross-section above which rescattering lowers the mean velocity such that the detection cross-section at the mean velocity is too low to be seen. This is not a hard cutoff, but rather a conservative estimate, as described in the Sec. IV. The lines labeled as “Xenon100, SD, n” and “COUPP, SD, p” are derived from Refs. [56, 57]. The models are parameterized by their effective DD cross-section and the DM mass.

V. RESULTS

The flux and effective cross-section calculated in Sections.III and IV can be combined to approximate the total number of events expected in the detector as

$$N_{\text{sig}} = \Phi \Sigma(\langle s \rangle) \Delta t, \quad (39)$$

where Δt is the lifetime of the experiment. We conduct the following three different analyses to evaluate the sensitivity of SK and HK for the boosted DM search as we proposed.

In the first analysis, using the CL_s method we determine the 95% exclusion region implied by the current SK analysis published in [55], which combines runs I and II of the SK experiment with lifetime 2287.8 days and assumes signal coming from the full solid angle of the sky. The number of proton recoil events observed in SK runs I and II is 38, with an expected atmospheric neutrino and other process background totaling 49.6 events [55]. The background is assumed to follow a log-normal distribution centered at 49.6 events with a 20% width, which accounts for estimated systematic and theory uncertainties. We marginalize over the background distribution. Given the data observation and expected background, as well as a signal efficiency of 70%, models which predict more than 20.7 events before folding in the detection efficiency are currently excluded.

The second analysis determines the expected sensitivity if the analysis were extended to include runs III and IV of SK, as well as making use of the direction of proton recoils relative to the Sun, as described in Section.IV. Finally, we determine the expected sensitivity of the HK experiment assuming a lifetime equal to that of runs I-IV of the SK experiment and again making use of the proton recoil direction. The lifetime of runs I-IV is 4438.2 days. In these two analyses we determine 2σ sensitivity region assuming the signal is present. The expected backgrounds for the full SK and HK datasets, after applying a cut on the recoil angle with respect to the Sun, are 23.5 and 587 events respectively. Note that since we are focusing on a small angular region for these analyses, we expect systematic uncertainties to be greatly reduced by looking at a side-band away from the direction of the Sun, such that statistic uncertainty dominates for background estimation. We find that a 2σ excess can be claimed if 15.5 events and 230 events respectively are predicted before folding in the detection efficiency.

From the above limits, we determine the model parameter space for which N_{sig} events are excluded by current SK data and for which there would be sensitivity using the full data at SK or HK and applying a cut on the recoil angle of the proton. The results for the semi-annihilating models are presented in Fig. 2. It is emphasized once more that these results are essentially independent of the specific structure of the model, beyond the Lorentz structure of the DM-quark scattering amplitude. The results for two benchmark two-component DM models are presented in Fig. 3. While these results do introduce some model dependence via the ratio of couplings and masses, rough scaling of the limits holds as

$$\sigma_{\text{DD}, A}^{\text{lim}}(m_B/m_A, g_B/g_A) \sim \sigma_{\text{DD}, A}^{\text{lim}}(2/3, 5) \left(\frac{g_B/g_A}{5} \right)^2, \quad (40)$$

provided that $1.1 \lesssim m_A/m_B \lesssim 2.2$. Beyond this mass range, (40) may not be a good approximation, but, particularly for smaller mass ratios, there should still be a significant signal. Such models would predict even more energetic

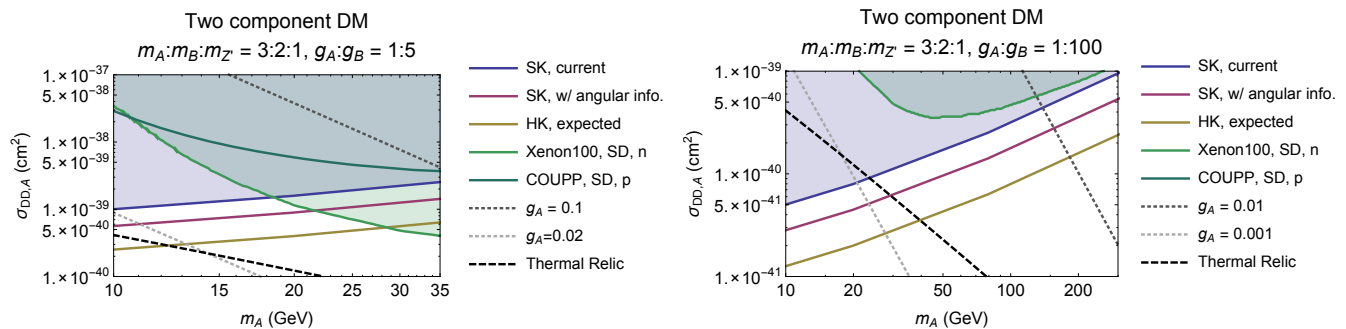


FIG. 3: Limits on the parameter space of two-component models for two different benchmark parameter choices. The lines labeled as “SK, current,” “SK, w/ angular info.,” and “HK, expected” are obtained as described in V. The lines labeled “Xenon100, SD, n” and “COUPP, SD, p” are derived from Refs. [56, 57]. The models are parameterized by the effective DD cross-section and the DM mass for the heavier DM particle ψ_A , which is assumed to make up nearly 100% of cosmic DM. The thermal relic lines assume $\Omega_A = 0.23$ and are derived using eq. 22. For reference, lines of corresponding to models with fixed coupling g_A are shown as well.

incident DM particles in the detector and the resulting proton collisions would be beyond the regime where the approximations we make hold. A full Monte Carlo simulation study including inelastic scattering is likely required to determine the constraints. Note that a larger g_A/g_B ratio tends to give more viable, reachable parameter space. Even for the large choice of $g_A/g_B = 100$ as in the right plot of Fig. 3, g_B is still well within perturbative region. Nonetheless such large coupling hierarchy may be unnatural if both ψ_A and ψ_B couples to the same Z' directly. One may resolve this by assuming ψ_A couples to another $U(1)'$ gauge boson which mixes with the Z' with which ψ_B couples directly. Alternatively, ψ_A, ψ_B may couple to quarks through two separate Z' s entirely, with $\psi_A\psi_A \rightarrow \psi_B\psi_B$ generated by an separate operator. These all require going beyond the minimal model, yet are still reasonable possibilities. Lastly, we emphasize that the reason we propose the large ratio of g_B/g_A is to induce a sizable enhancement of $\sigma_{DD,B}$ respect to $\sigma_{DD,A}$. This can also be achieved by assuming ψ_A -proton and ψ_B -proton scattering are induced by different effective operators. Thus the a large ratio between $\sigma_{DD,B}$ and $\sigma_{DD,A}$ can be easily achieved in other scenarios.

VI. CONCLUSION

In this paper, we investigate a new search channel for boosted DM. As opposed to earlier related work [5], which focused on “secluded” DM where observable signals are only possible from DM annihilation in the galactic center, here we consider here the alternative case where the DM also has appreciable interactions with SM quarks so that it can be effectively captured in the Sun, where it annihilates and produces boosted DM that can be detected at large volume neutrino detectors. Such annihilation also determines the thermal relic abundance of the DM, which realizes thermal WIMP paradigm in an alternative way. We study both semi-annihilating DM and two-component DM as examples. In our simple example models, this detection mechanism provides additional sensitivity beyond conventional direct detection and collider searches in models with relatively light mediators that couple to light quarks and dominantly generate spin-dependent DM-nucleus scattering. Viable alternative models are possible and may be worth investigating.

We propose a new search based on elastic scattering induced proton tracks pointing towards the direction of the Sun, which is typically the primary search channel for DM and mediator with weak scale masses. In particular, we studied the sensitivity limits for boosted DM at Cherenkov-light based neutrino detector such as SK and its upgrade HK. The existing SK proton track data already has sensitivity to some of the parameter space, while the region that could be probed would be moderately enlarged by using an analysis that takes into account directional information with the full SK data and would be even more enlarged by a search at HK. Future large-volume liquid Argon neutrino detector based on ionization signals may significantly extend the sensitivity. As already studied in [5], the single ring electron channel may be complementary to or even more important than the proton track channel when a mediator mass much lower than weak scale is involved. It is also possible to have complementary indirect detection signals, for instance from ϕ decay in semi-annihilation model with $\chi\chi \rightarrow \chi^\dagger\phi$ or Z' decay in $\psi_B\psi_B \rightarrow Z'Z'$ in two-component model. We leave an investigation of these interesting possibilities to future work.

Boosted DM is generic in scenarios with multiple DM components or single DM specie with non-minimal stabilization symmetry. Detecting boosted DM can be crucial in uncovering such features of DM sector, and in some cases can even be a smoking gun signal. As we have seen, even in the models where appreciable DM-nucleon scattering rate is

present to facilitate solar capture of DM, current or near future DD limit can still be far from the parameter region that can be probed by dedicated boosted DM search. It is particularly intriguing that large-volume neutrino detectors can be repurposed to search for boosted DM. These signals provide further motivation for a thorough examination of new search strategies for DM sector with non-minimal structure.

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Appendix A: Isospin Dependence of Spin-Dependent DM-nucleon Scattering

The isospin dependence of DM-nucleon scattering can be parameterized as

$$\sigma_{\chi,N} = \sigma_{\chi p} \frac{[A_u \Delta_u^{(n)} + A_d \Delta_d^{(n)} + A_s \Delta_s^{(n)}]^2}{[A_u \Delta_u^{(p)} + A_d \Delta_d^{(p)} + A_s \Delta_s^{(p)}]^2}, \quad (\text{A1})$$

where $A_{u,d,s}$ parameterize the relative sizes of the couplings of DM to the quark flavors and $\Delta_{u,d,s}^{(N)}$ are given by [60]

$$\Delta_u^{(p)} = 0.78, \quad \Delta_d^{(p)} = -0.48, \quad \Delta_s^{(p)} = -0.15 \quad (\text{A2})$$

and

$$\Delta_u^{(n)} = \Delta_d^{(p)}, \quad \Delta_d^{(n)} = \Delta_u^{(p)}, \quad \Delta_s^{(n)} = \Delta_s^{(p)}. \quad (\text{A3})$$

Note that if the coupling to quarks is universal, then this factor is 1. For simplicity and concreteness, we assume this case.

Appendix B: Parametrization of the DM-nucleon Cross-section

Beyond the discussion in Section II, we parametrize all of the relevant cross-sections using the effective DD cross-section σ_{DD} , rather than the couplings. This parameter is related to the Lagrangian parameters by

$$\sigma_{\text{DD}} = \frac{3\mu_{\chi,N}^2}{\pi M^4} \left(\sum_q \Delta_q \right)^2, \quad (\text{B1})$$

for models corresponding to the operator in eq. (4) with χ denoting the DM particle in semi-annihilating models and $\chi = \psi_A, \psi_B$ in two component models. M is the suppression scale of the corresponding operator. For the model corresponding to eq. (12), the relation is

$$\sigma_{\text{DD}} = \frac{m_N^2}{2\pi M^4} v_0^2 \left(\sum_q \Delta_q \right)^2, \quad (\text{B2})$$

where v_0^2 is the mean squared velocity of DM in the local halo.

In terms of these cross-section ‘‘parameters,’’ the full form of the scattering cross-section for DM-nucleon interactions for v^0 -like interactions is

$$\frac{d\sigma_{\chi,N}}{dt} = \frac{\sigma_{\text{DD}}}{24\mu_{\chi,N}^2} \cdot \frac{m_{Z'}^4}{(t - m_{Z'}^2)^2} \frac{t^2 + 2t(2E_\chi m_p - m_p^2 - m_\chi^2) + 8m_p^2(E_\chi^2 + 2m_\chi^2)}{\lambda(s, m_p^2, m_\chi^2)} F(-t)^2, \quad (\text{B3})$$

where E_χ is the energy of incoming DM χ in the lab frame, and in the lab frame $t = 2m_p(m_p - E_p)$, $s = m_\chi^2 + m_p^2 + 2E_\chi m_p$, E_p is the energy of the scattered proton, $\lambda(x, y, z) \equiv x^2 + y^2 + z^2 - 2xy - 2yz - 2zx$, and F is the form factor given in eq. (8). For v^2 -like interactions, the full form of the differential cross-section is

$$\frac{d\sigma_{\chi,N}}{dt} = -\frac{\sigma_{\text{DD}}}{8\mu_{\chi,N}^2 v_0^2} \frac{m_{Z'}^4}{(t - m_{Z'}^2)^2} \frac{t(m_\chi + m_N)^2}{m_N^2 p_\chi^2} F(-t), \quad (\text{B4})$$

where p_χ is the lab-frame momentum of χ .

Appendix C: Detailed Determination of the Capture and Evaporation Rates

In this Appendix, we calculate the rates C and E presented in (25) and (32) respectively. We proceed factor by factor in the integrands.

A DM particle χ that has velocity w at solar radius r and scatters off of a nucleus n in the Sun is captured if its final velocity v in the solar frame is less than the escape velocity v_{esc} at the radius of the scattering. This cross-section is given by

$$\sigma_{\chi,p}(w \rightarrow v)|_{v < v_{\text{esc}}} = \int_{-1}^{c_{\theta,\text{max}}} dc_\theta \frac{d\sigma_{\chi,p}}{dc_\theta}, \quad (\text{C1})$$

where c_θ is the cosine of the CM scattering angle. The minimum scattering angle is given by

$$c_{\theta,\text{max}} = 1 - \frac{m_n m_\chi}{p_{\chi,\text{CM}}^2} \left(\frac{1}{\sqrt{1 - v_{\text{esc}}^2}} - \frac{1}{\sqrt{1 - w^2}} \right), \quad (\text{C2})$$

where $p_{\chi,\text{CM}}$ is the momentum of the DM particle in the CM frame. At the low momentum transfers in the collisions in the Sun, the DM scatters coherently off of the nuclei in the Sun. Since the scattering is SD in the models we consider and helium is spin 0, scattering occurs exclusively on hydrogen nuclei, which are protons. The escape velocity is determined from solar model [58].

The DM velocity going into these collisions is determined from the local DM velocity distribution. The velocity at distances far from the Sun is approximately given by a Boltzmann distribution [59]

$$f(u) = \sqrt{\frac{6}{\pi}} \frac{u}{v_G \bar{v}} \exp\left(-\frac{3}{2} \frac{u^2 + v_G^2}{\bar{v}^2}\right) \sinh\left(\frac{3u v_G}{\bar{v}^2}\right), \quad (\text{C3})$$

where v_G is the velocity of the Sun in the Milky Way and \bar{v}^2 is the local mean squared velocity of DM. As the DM falls into the gravitational potential of the Sun, it gains speed such that, by conservation of energy,

$$w(r) = \sqrt{u^2 + v_{\text{esc}}^2(r)} \quad (\text{C4})$$

at distance r from the center of the Sun. In addition, the density of DM in the Sun gets a Sommerfeld enhancement from falling into the Sun of w/u .

The local number density of DM far from the Sun is taken to be $m_\chi n_\chi = 0.3 \text{ GeV}/\text{cm}^3$. The number density of hydrogen atoms at a distance r from the center of the Sun is again determined from the solar model [58].

Putting all of these pieces together, we find that the rate for DM capture in a volume element dV near a distance r from the center of the Sun and coming from velocity between u and $u + du$ at distance far from the Sun is given by

$$dC = dV du (\sigma_{\chi,H}(w \rightarrow v)|_{v < v_{\text{esc}}} w n_H) \left(\frac{w}{u} n_\chi\right) f(u), \quad (\text{C5})$$

where the first factor is the number of interactions at r per DM particle and the second factor is the number density of DM including the Sommerfeld enhancement. To obtain (25), we integrate over the velocities u for which scattering to $v < v_{\text{esc}}$ is possible and over the volume of the Sun. The upper limit on u such that capture is possible at a radius r is given by

$$u < \frac{2\sqrt{m_\chi m_p} v_{\text{esc}}}{m_\chi - m_p}. \quad (\text{C6})$$

for $m_\chi > m_p = m_H$.

The determination of the evaporation follows similar arguments. The primary differences are in the detailed kinematics. For instance, the minimum scattering angle is given by

$$c_{\theta,\max} = 1 + \frac{m_\chi^2}{p_{\chi,\text{CM}}^2} \left(1 - \frac{1}{\sqrt{1 - v_{\text{esc}}^2}} \right). \quad (\text{C7})$$

There is no Sommerfeld enhancement, but the DM density is given by the captured DM density. If the DM undergoes a sufficient number of scatters before it annihilates, then its number density is thermalized and given by

$$n_\chi(r) = N \frac{\exp(-M\Phi(r)/T_W)}{\int_0^{R_\odot} dr 4\pi r^2 \exp(-M\Phi(r)/T_W)}, \quad (\text{C8})$$

where $\Phi(r)$ is the gravitational potential of the Sun at r and T_W is the thermalized DM temperature. The thermalized DM temperature is an averaged solar temperature, which we take to be $T(r)$, the temperature of the Sun at r .

The distribution of hydrogen velocities follows a thermal distribution at r given by

$$f(u_H) = \left(\frac{m_p T(r)}{2\pi} \right)^{3/2} \exp\left(-\frac{m_p u_H^2}{2T(r)}\right). \quad (\text{C9})$$

In order to induce evaporation, the hydrogen velocity must be sufficiently large to kick the DM to a velocity above the escape velocity. In the non-relativistic limit, this minimal velocity is given by

$$u_H > \frac{(m_p + m_\chi)v_{\text{esc}}}{2m_p}. \quad (\text{C10})$$

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