Terrestrial Probes of Electromagnetically Interacting Dark Radiation

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We study the possibility that dark radiation, sourced through the decay of dark matter in the late Universe, carries electromagnetic interactions. The relativistic flux of particles induces recoil signals in direct detection and neutrino experiments through its interaction with millicharge, electric/magnetic dipole moments, or anapole moment/charge radius. Taking the DM lifetime as 35 times the age of the Universe, as currently cosmologically allowed, we show that direct detection (neutrino) experiments have complementary sensitivity down to $\epsilon \sim 10^{-11}$ ($10^{-12}$), $d_\chi/\mu_\chi \sim 10^{-5} \mu_B$ ($10^{-13} \mu_B$), and $a_\chi/b_\chi \sim 10^{-2}$ GeV$^{-2}$ ($10^{-8}$ GeV$^{-2}$) on the respective couplings. Finally, we show that such dark radiation can lead to a satisfactory explanation of the recently observed XENON1T excess in the electron recoil signal without being in conflict with other bounds.

I. INTRODUCTION

The identity of dark matter (DM) in the Universe remains to be a mystery. While the gravitational manifestations of DM are numerous and well studied, the connection between DM and the Standard Model (SM) of particles and fields is unknown. A large amount of effort and resources have been invested in the attempts to detect the non-gravitational interaction of SM particles with DM. Thus far, these efforts have resulted in stringent limits on the strength of interaction for certain types of DM. In particular, DM in the form of individual particles with masses comparable to the masses of SM particles has been constrained [1]. While many of these searches were targeting weakly interacting massive particles (or WIMPs), over time it has become clear that the sensitivity of the existing experiments extends beyond WIMP-nucleon scattering and beyond the electroweak scale for DM masses, see e.g. [2][7].

In parallel to the experimental developments, the last two decades brought a more general view on DM physics. Early on (e.g. with the example of supersymmetric WIMP relics [8]) it was understood that the DM may not come “in isolation”, but be, in fact, a member of a more generic dark sector. This sector could comprise additional heavy particles charged under the SM gauge group (e.g. weak scale supersymmetry), or alternatively may include the mediators of interaction or carriers of a “dark force”, as well as new very light degrees of freedom known as dark radiation (DR). Light particles from dark sectors would almost invariably have to have a small coupling to the SM. Dark forces and light mediators have received significant attention in the literature, both in the cosmological and laboratory settings [9]-[13]. In comparison, dark radiation has mostly been studied in the context of its contribution to the cosmological expansion rate, parametrized by $N_{\text{eff}}$. One interesting aspect of DR (for a set of representative ideas see Refs. [14]-[26]) is that it can be created in a non-thermal way, and therefore differ in energy from the characteristic energies of the quanta of the cosmic microwave background. In this paper we will discuss some observable signatures of interacting DR in the regime when $\omega_{\text{DR}} > \omega_{\text{CMB}}$ holds for the respective typical energies.

The goal of the present paper is to consider a class of dark radiation models that interacts with the SM via a single photon exchange. We will consider DR originating from the late decays of DM particles with mass $m_X > \text{keV}$. This range will allow to probe such scenarios using DM direct detection experiments, and for $\omega_{\text{DR}} > 200 \text{keV}$, using sensitive underground neutrino experiments. The first goal of this project is to map the sensitivity of the best existing experiments vs. the strength of DR electromagnetic form factors controlling the interaction and available DR fluxes. The latter depend on the mass and lifetime of decaying progenitor particles.

The second goal of this paper is to consider DR and the multitude of possible electromagnetic form factors as candidates for the explanation of the recently reported signal excess in the XENON1T experiment [27]. The excess, consistent with the injection of $O(2-3)$ keV electromagnetic energy, has numerous candidate explanations. The collaboration itself has tried to connect it with DR coming from the Sun, in form of the axions and/or neutrinos with electromagnetic dipole interaction. Our goal is then to generalize it to a number of possible form factors, in the situation when the DR radiation flux is maximized by employing the DM→DR decay. It is unclear if the explanation of this excess can be achieved with DR without being in conflict with other measurements and constraints.

We will adopt a fairly minimal scheme, where the progenitor decay, $X \rightarrow \bar{\chi}_i \chi_j$, is sourcing DR in form of $\chi_i$’s, which we assume to be a Dirac fermion. The Lagrangians specifying the $\chi_i$ interactions with the photon $A_\mu$ or its field strength tensor $F_{\mu\nu}$ in ascending order of their di-
mensionality read,
\[ \mathcal{L}_{\text{dim}=4} = e \epsilon \chi^\gamma \mu X A_\mu, \]
\[ \mathcal{L}_{\text{dim}=5} = \frac{1}{2} \mu_\chi \bar{\chi} \sigma^{\mu \nu} \chi F_{\mu \nu} + \frac{i}{2} d_\chi \bar{\chi} \sigma^{\mu \nu} \gamma^5 \chi F_{\mu \nu}, \]
\[ \mathcal{L}_{\text{dim}=6} = -a_\chi \bar{\chi} \gamma^5 \lambda \partial^\nu F_{\mu \nu} + b_\chi \bar{\chi} \gamma^5 \lambda \partial^\nu F_{\mu \nu}. \]
Here \( e \epsilon \) is the millicharge \((mQ)\), \( \mu_\chi \) and \( d_\chi \) are the magnetic and electric dipole moments (MDM and EDM), and \( a_\chi \) and \( b_\chi \) are the anapole moment and charge radius interaction (AM and CR), respectively. If \( \chi \) were a single Majorana fermion, then only the \( a_\chi \) interaction is allowed. If \( \chi \) is instead a complex scalar, \( \mu_\chi \), \( d_\chi \) and \( a_\chi \) will vanish. For a real scalar such form factors simply do not exist. Thus, we consider DR with the Dirac fermion case to have the most variety of phenomenological consequences. In the past, various aspects of (effective) electromagnetic couplings of dark sector particles (including constraints from direct detection, beam dump limits, SM precision observables, colliders and stellar constraints) were explored in a number of publications [28–43]; for the mQ interaction, see the recent review [13] and references therein.

The paper is organized as follows: In Sec. [II], we compute the DR flux from decaying DM (DDM), followed by an overview of experiments in Sec. [III]. In Sec. [IV] we derive the expected electron recoil (ER) or nuclear recoil (NR) event rate by the DR. In Sec. [V] we demonstrate the constraints and the forecasts of sensitivity on the parameter space and present the fit to the XENON1T excess. Finally, in Sec. [VI] we conclude and give outlooks.

II. DARK RADIATION FLUX

In this section, we collect the ingredients for the calculation of the DR flux from DDM. For simplicity, we consider a two-body decay \( X \rightarrow \chi \chi \) and assume a single decay channel for \( X \). In that case, \( X \chi \chi \) coupling can be traded for the lifetime of \( X \). The expected Galactic energy differential flux is given by
\[ \frac{d\phi_{\text{gal}}^{\text{gal}}}{dE_\chi} = \frac{e^{-t_0/\tau_X}}{m_X \tau_X} \frac{dN_X}{dE_\chi} R_\odot \rho_\odot \langle D \rangle, \]
where \( t_0 \) is the age of the universe, \( m_X \) and \( \tau_X \) are the DM mass and lifetime, \( R_\odot \simeq 8.33 \) kpc is the distance between the Sun and the Galactic center, \( \rho_\odot = 3 \times 10^3 \text{keV/cm}^3 \) is the local DM energy density and \( \langle D \rangle \simeq 2.1 \) is the averaged \( D \)-factor assuming a NFW profile [41]. For simplicity, we assume a 100% decaying fraction of DM; if this is not the case, the formulas are to be dressed with the DDM fraction in an obvious way. The DR injection spectrum is monochromatic,
\[ \frac{dN_X}{dE_\chi} = 2 \delta \left( E_\chi - \frac{m_X}{2} \right), \]
with a negligible spread by the parent DM velocity dispersion.
galactic DR flux (solid lines) originating from DDM with \(m_X = 50, 100, 500\) keV; we apply a 2% Gaussian smearing on the monochromatic Galactic flux for visualization. The fluxes are compared to the solar neutrino flux (solid black line) taken from [54, 57]. Below 10 keV, we include the contribution from pi meson decay, photoproduction, and bremsstrahlung from [58]. As can be seen, both galactic and extragalactic DR fluxes are, in magnitude, in roughly the same ballpark as the solar neutrino flux.

### III. EXPERIMENTS

In this section, we outline the considered experiments and the way to derive constraints and forecasts of sensitivity on the parameter space. For \(E\) in the \(\mathcal{O}(\text{keV})\) energy ballpark, we consider the scattering of DR in the XENON1T detector. Neutrino experiments such as Borexino, Super-Kamiokande (SK) as well as the future Hyper-Kamiokande (HK) and Deep Underground Neutrino Experiment (DUNE) have larger energy threshold, MeV–GeV range, and we consider DR-electron scattering for which the solar neutrinos \((E_R < 30\text{ MeV})\) and atmospheric neutrinos \((E_R > 30\text{ MeV})\) become the main background. For Borexino, we consider DR-proton scattering in addition.

#### A. XENON1T

The XENON1T detector, located underground at the Gran Sasso laboratory, is a dual-phase time projection chamber with liquid and gaseous xenon. The registered signals include prompt scintillation (S1) and secondary scintillation from ionization (S2). In a recent analysis [27], an excess of events was identified in the S1 data at \(\mathcal{O}(\text{keV})\). Although poorly understood backgrounds exist [27, 67], the possibility that this signal is due to new physics has been entertained abundantly, see [27, 68–72]. The excess is not in conflict with an earlier S2-only analysis by the experiment [59]. Here we derive both, the favored region for the anomaly and the constraints on the parameter space using the S1+S2 and S2-only data. Details on the limit-setting procedure can be found in [70] which we follow here; see also [74].

#### B. Borexino

The Borexino experiment features a liquid scintillator-based detector with 280 ton fiducial mass, primarily designed to measure solar neutrinos in the quasi-elastic scattering signal with electrons [75]. We use the latest data from the CNO neutrino search of phase-III [60, 61] of the experiment, with an exposure of 209.4 ton-yr. Between the threshold energy 320 keV and 2640 keV, the observed event rate and the best-fit background plus solar neutrino-induced rate are reported for each energy bin. Note that the standard neutrino events are a background in our consideration. The detection efficiency is assumed to be unity. We derive 95% C.L. limits using the CL\(_x\) method [76].

For heavier progenitor masses, we further consider the proton recoil signal in the Borexino detector. Here, we adopt Birk’s law to account for the energy quenching in the organic scintillator,

\[
E_{\text{vis}} = \int_0^{E_R} \frac{dE}{1 + k_B dE/dx},
\]

where \(E_{\text{vis}}\) is the visible energy, \(k_B \approx 0.01\text{ cm/MeV}\) is Birk’s constant and \(dE/dx\) is the stopping power which we compute using the SRIM computer package; see also [77]. For the scintillator pseudocumine C\(_9\)H\(_{12}\) with a mass density of \(\rho = 0.88\text{ g/cm}^3\), the stopping power for protons is roughly \(dE/dx \sim \mathcal{O}(100)\text{ MeV/cm}\), albeit energy-dependent. For electrons, \(dE/dx \sim \mathcal{O}(10^{-3})\text{ MeV/cm}\) so that we are allowed to neglect the energy quenching since \(dE/dx \ll k_B^{-1}\). We utilize the same data and method presented above to derive the corresponding constraint. See Fig. 2 for a demonstration of the event rate from different operators and the Borexino data.

#### C. Super-Kamiokande

Super-Kamiokande (SK) is a neutrino experiment with a water-based Cherenkov detector located 2.7 km underground in Japan. The fiducial mass is 22.5 kton.
that the DR-induced events
the various signal strengths can be derived by requiring
its data-taking time [63]. The ensuing 90% C.L. limits on
The estimated background is 3993

corresponding efficiency is taken from [62].

N

satisfy,

\[ N_{\text{sig}}^{\text{DR}} \leq \text{Max}[0, N_{\text{obs}} - N_{\text{bkg}}] + 1.28\sqrt{N_{\text{obs}}}. \] (7)

**IV. EVENT RATE**

### A. Scattering on bound electrons

For small progenitor mass, the resulting DR is low-
energetic enough that we need to account for bound state
effects in the DR-electron scattering and resulting atomic
ionization process. Combining the DR flux from Sec. III
and the differential cross section given in App. A, the
differential event rate for scattering with the electrons is

\[
\frac{dR}{dE_R} = \kappa N_T \epsilon(E_R) \int_{q_{-}}^{q_{+}} dq \int_{p_{\chi}}^{p_{\chi,\text{max}}} dp_{\chi} \frac{d\sigma_{\chi}}{E_{\chi} dE_{\chi}} d\chi_d q dE_R,
\] (8)

where \( \kappa \) is the exposure of the experiment, \( N_T \) is the
time of detection efficiency, and \( \epsilon(E_R) \) is the
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from DDM that includes both, the galactic and extragalactic components. The minimum $\chi$-momentum for a given recoil energy $E_R$ and momentum transfer $q$ is

$$\chi_\text{min} = \frac{q}{\Delta E q} \left[ x + \frac{\Delta E^2}{q^2} \right] \left[ x + \frac{4m_n^2}{q^2} \right],$$  \hspace{1cm} (9)$$

where $x = 1 - \Delta E^2 / q^2$ with $\Delta E = E_R + |E_B^{n,l}|$ being the deposited energy and $E_B^{n,l}$ is the binding energy of the bound state orbital $(n,l)$. The upper boundary of the $\chi$ integration is given by $\chi_{\text{max}} = \chi_n$. The integration boundaries of $q$ are given by

$$q_+ = \chi_{\text{max}}^\text{max} + \sqrt{(m_X - 2\Delta E)^2 - 4\chi_n^2}, \hspace{1cm} q_- = \Delta E.$$

To obtain the total event rate, we sum up the contributions from all kinematically available $(n,l)$ shells.

### B. Scattering on free particles

For larger progenitor mass, the O(MeV–GeV) ER signals induced by DR are best probed in the large-volume neutrino experiments mentioned in Sec. III. For such recoil energies, the initial electron can be considered as a free particle. With the recoil cross section given in Appendix B the total differential event rate reads

$$\frac{dN}{dE_R} = \kappa N_T \varepsilon(E_R) \int_{\chi_{\text{min}}}^{\chi_{\text{max}}} p_x \frac{d\chi}{E_x} \frac{d\phi_x}{d\chi} \frac{d\sigma}{dE_R}.$$  \hspace{1cm} (11)$$

Here, the lower integration boundaries of $p_x$ is given through $p_x^{\text{min}} = \sqrt{(E_x^{\text{min}})^2 - m_X^2}$ with

$$E_x^{\text{min}} = E_R^2 + \frac{1}{2m_e} \sqrt{m_e(E_R + m_e)(E_Rm_e + 2m^2)},$$

and the upper boundary as before. The expected number of events is given by

$$N_{\text{sig}} = \int_{E_{\text{th}}}^{E_{\text{max}}} dE_R \frac{dN}{dE_R},$$

where $E_{\text{th}}$ is the threshold recoil energy. The maximal recoil energy $E_{\text{max}}$ is either given by the energy range of the experimental data or half of the progenitor mass.

For large enough $m_X$, the DR is energetic enough to generate $O(\text{keV–MeV})$ NR in direct detection and neutrino experiments. The framework for NR is the same as scattering on free electrons discussed above, but the recoil cross section become target-dependent; see [58] for detailed formulas of the nuclear recoil cross section.

### A. Constraints on the effective couplings

We show the resulting constraints (shaded regions) and forecasts of sensitivity (lines) for millicharged DR in the left panel of Fig. 3, for MDM/EDM interactions in the right panel of Fig. 3 and for AM/CR interactions in Fig. 4. We also show the XENON1T excess favored region, with details on the fitting procedure given in Sec. V B. Previous constraints derived from the anomalous energy loss in red giant stars (RG) and SN1987A cooling are included for comparison [39, 81], which apply when $m_X$ is smaller than the plasma frequency in the stellar environment: $m_X \lesssim 10 \text{keV} (20 \text{MeV})$ for RG (SN1987A); see also [40, 82]. For $m_Q$, we note that there exist additional bounds from galaxy cluster magnetic fields [83] and the timing of radio waves [84]. However, both of them scale with the DR mass, thus they are not included in the figures.

Due to the energy dependence in the cross sections, the experiments with higher threshold are more important for higher dimensional operators. We see that current SK and future HK and DUNE can all provide better sensitivity than current best limit from the stellar energy loss for dimension 5 and 6 operators, assuming $\tau_X = 35t_0$. For $m_Q$, the improvement of sensitivity between experiments that probe free electron scattering and XENON1T is not so notable compared to higher dimensional operators.

The constraints derived in this paper for $m_X = 0$ also apply to electromagnetic form factors of neutrinos if they play the role of DR. On the other hand, if only taking the SM weak interactions of neutrinos, direct detection and neutrino experiments put bounds on the mass and lifetime of the progenitor $X$; see, e.g., Ref. [17].

### B. XENON1T excess

In light of the recent excess in the $O(\text{keV})$ recoil energy range observed by XENON1T [27], we also explore the possibility of explaining the excess with DR. assuming...
the background modelling is correct. This lines up with several other new physics scenarios and their constraints that have been investigated in this context [24, 68]. Moreover, PandaX-II reports for its own data that it is both, consistent with a new physics contribution as well as with a fluctuation of background [65]. Thus the observational status of an excess in XENON1T remains unclear at the moment.

In Fig. [5] we show the best-fit event rate induced by DR and the data in the energy range [0, 10] keV in two fitting scenarios, including and excluding the first bin. By excluding the first bin, the recoil spectrum can better fit to the peak of excess, but at the expense of significantly overshooting the first bin, (See also related discussions in Refs. [72, 80]). When the first bin is included in the fit, the second bin cannot be filled but the overall fit is still satisfactory, similar to the anomalous neutrino magnetic dipole moment explanation [27, 68]. The corresponding best-fit parameters and $\chi^2$/dof are given in Table [11]. We observe that higher-dimensional operators yield improved fits to the excess, as their recoil spectra are less peaked at low $E_R$. We also note that the best-fit coupling of dimension 5 operator is consistent with

$$\phi_X^{\text{best}} \times (\mu_{\nu}^{\text{best}})^2 \simeq \phi_{\nu}^{\text{solar}} \times (\mu_{\nu}^{\text{best}})^2,$$

although the free electron approximation is adopted in [68].

For massless DR, the favoured parameter space for the excess is excluded by stellar energy loss constraints, such as red giant stars for dimension 4 and 5 operators and SN1987A for the dimension 6 operators, shown in Fig. [5] and taken from [39, 51]; see also [40, 52]. However, stellar energy loss is effective only when $\chi$ production is kinematically allowed. Taking DR with $m_\chi = 70$ keV, the constraints from the stellar energy loss are alleviated, while, at the same time, leaving the fits to the XENON1T excess to remain unchanged. This is owed to the relativistic nature of considered DR. Finally, we consider the constraint from the measured number of relativistic degrees of freedom $N_{\text{eff}}$, as $\chi$ particles are also populated in the early universe through plasmon decay and electron-positron annihilation [39, 40, 51]. As shown in Fig. [6] for mQ (left panel), MDM/EDM (middle panel) and AM/CR (right panel), there remains allowed parameter space for explaining the XENON1T excess for dimension 4 and 6 operators. For dimension 5 operators, the viable parameter space is covered by the SN1987A bound.
The emerging DR flux from DDM is then probed in underground rare-event searches. For $m_X \lesssim 1$ MeV direct detection experiments offer the best sensitivity with their ability of registering keV-scale energy depositions and below. Heavier progenitors are better probed with neutrino experiments, as $\chi$-induced events leave MeV-scale signals. For concreteness, in this work we have chosen a benchmark value of $\tau_X = 35\tau_0$ with the bulk of DM still to decay in the distant future.

The scattering of $\chi$ on electrons is the most important signal channel. We demonstrate that the recent (S1+S2) data from the XENON1T experiment yields $\epsilon \lesssim 2 \times 10^{-11}$ at $m_X \simeq 10$ keV, and $d_\chi, \mu_\chi \lesssim 2 \times 10^{-9} \mu_B$ as well as $a_\chi, b_\chi \lesssim 2 \times 10^{-2} \text{GeV}^{-2}$ at $m_X \simeq 100$ keV. In addition, we find that it is also possible to reach a satisfactory fit to the reported excess of events seen in the XENON1T data at few keV energy. The fit improves by increasing the dimensionality of the operator, as the lowest energy bin in the data prohibits too strong of an IR-biased signal. The AM/CR interaction thereby yields the best fit. The DR mass-dependence is relatively mild in those drawn conclusions as these particles retain their (semi-)relativistic nature except at the very kinematic edge $2m_\chi \simeq m_X$. However, stellar and cosmological constraints critically depend on $m_\chi$. By choosing a benchmark value of $m_\chi = 70$ keV we demonstrate that a XENON1T explanation remains intact for $m_Q$ and for the dim-6 AM and CR operators, evading bounds from the anomalous energy loss inside RG stars, of the protoneutron star of SN1987A and from the cosmological $N_{\text{eff}}$ limit.

For progenitor masses $m_X \gtrsim 1$ MeV Borexino has the best sensitivity reaching $\epsilon \lesssim 10^{-12}$, $d_\chi, \mu_\chi \lesssim 3 \times 10^{-12} \mu_B$ and $a_\chi, b_\chi \lesssim 2 \times 10^{-6} \text{GeV}^{-2}$ at $m_X \simeq 1$ MeV. These limits rely on the detailed modeling of Borexino backgrounds and its solar neutrino-induced events. The limits are eventually surpassed by the ones from SK, once DR induces electron recoils above the solar neutrino endpoint energies. Best sensitivity is attained for $m_X \simeq 100$ MeV with $\epsilon \lesssim 4 \times 10^{-13}$, $d_\chi, \mu_\chi \lesssim 10^{-13} \mu_B$ and $a_\chi, b_\chi \lesssim 4 \times 10^{-9} \text{GeV}^{-2}$. Finally, we also provide forecasts for HK and DUNE, with relatively mild expected improvements.

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Appendix A: Recoil cross section on bound electron

For scattering on bound electrons, the differential cross section for massive DR that may either be relativistic or
non-relativistic reads \[^{[4]}\]

\[
\frac{d\sigma}{d\Omega_{ER}}|_{(n,l)} = \frac{\sigma_e}{8\mu_e^2 E_R \rho_X E_X} \int d\Omega_R \, q |f_{n,l}(\vec{q})|^2 \times |F_X(q, E_X)|^2 ,
\]

where \(E_R\) is the electron recoil energy, \(v\) is the relative velocity, \(\mu_e\) is the reduced mass between \(\chi\) and \(e^-\), \(q\) is the momentum transfer, \(d\Omega_R\) is the solid angle element of the final state electron and \(|f_{n,l}(\vec{q})|^2\) is the atomic form factor for the atomic state \((n, l)\) \[^{[3 \quad 89]}\]. In this work, we use the numerical result of \(|f_{n,l}(\vec{q})|^2\) that was obtained from an atomic calculation described in \[^{[70]}\]. We normalize our results to the non-relativistic effective scattering cross section \(\sigma_e\) on a free electron, evaluated at a typical atomic momentum transfer \(q_0 \simeq \alpha m_e\) and at vanishing kinetic energy, \(E_X = m_X\).

\[
\sigma_e = \frac{\mu_e^2}{16\pi m_X^2 m_e^2} |\mathcal{M}(q = q_0, E_X = m_X)|^2 ,
\]

where \(|\mathcal{M}|^2\) is the squared amplitude, averaged over the initial state spins and summed over the final state spins. We assume \(m_e\) is much larger than the momentum transfer \(q\) and the deposited energy \(\Delta E\), which is valid when we consider the scattering with bound electrons in this work. We then expand \(|\mathcal{M}(q, E_X)|^2\) in \(q/m_e = O(\alpha)\) and are careful to additionally retain the leading terms in a velocity-expansion that become relevant in the non-relativistic limit \(v \ll \alpha\). For mQ, EDM, CR the leading order terms in both expansions coincide; for MDM and AM, the leading order term in \(q/m_e\) is velocity suppressed in the non-relativistic limit, and we add the term that is not velocity suppressed but of higher order in \(q/m_e\). We find

\[
mQ : \tilde{\sigma}_e = \epsilon^2 \pi \alpha^2 \frac{16m_e^2m_e^2}{q_0^2(m_X + m_e)^2} , \quad (A3)
\]

\[
MDM : \tilde{\sigma}_e = \mu_X^2 \alpha^2 \frac{m_X^2}{(m_X + m_e)^2} , \quad (A4)
\]

\[
EDM : \tilde{\sigma}_e = \frac{4m_X^2m_e^2}{q_0^2(m_X + m_e)^2} , \quad (A5)
\]

\[
AM : \tilde{\sigma}_e = a_X^2 \alpha^2 \frac{m_X^2}{(m_X + m_e)^2} , \quad (A6)
\]

\[
CR : \tilde{\sigma}_e = b_X^2 \alpha^2 \frac{4m_X^2m_e^2}{(m_X + m_e)^2} , \quad (A7)
\]

where \(\epsilon\) is the millicharge of \(\chi\) in units of the elementary charge \(e\), \(\mu_X\) and \(d_X\) are the magnetic and electric dipole moment of \(\chi\), and \(a_X\) and \(b_X\) are the anomole moment and charge radius coupling of \(\chi\). The listed non-relativistic effective scattering cross sections agree with the ones found in \[^{[3 \quad 38]}\].

The dark matter form factor is defined as

\[
|F_X(q, E_X)|^2 = \frac{|\mathcal{M}(q, E_X)|^2}{|\mathcal{M}(q = q_0, E_X = m_X)|^2} , \quad (A8)
\]

with the concrete expressions for the respective effective operators given by

\[
mQ : |F_X(q, E_X)|^2 = \frac{E_X^2 q_0^4}{m_X^2 q^4} , \quad (A9)
\]

\[
MDM : |F_X(q, E_X)|^2 = \frac{4m_X^2 (E_X^2 - m_X^2)}{q^2 m_X^2} + 1 , \quad (A10)
\]
EDM : \[ |F_{\chi}(q, E_\chi)|^2 = \frac{E_\chi^2 q_0^2}{m_\chi^2 q^2}, \] (A11)

AM : \[ |F_{\chi}(q, E_\chi)|^2 = \frac{4m_\chi^2 (E_\chi^2 - m_\chi^2) + q^2 m_\chi^2}{q_0^2 m_\chi^2}, \] (A12)

CR : \[ |F_{\chi}(q, E_\chi)|^2 = \frac{E_\chi^2}{m_\chi^2}. \] (A13)

Unlike in the direct detection literature that is concerned with chiefly non-relativistic scatterings, the form factors defined here carry an additional dependence on \( E_\chi \) as is generally the case for relativistic scattering processes. When taking the non-relativistic limit \( E_\chi \approx m_\chi \), we retrieve the non-relativistic dark form factors found in the literature [3, 38]. In addition, in the relativistic limit, \( E_\chi \gg m_\chi \), the helicity suppression introduced by \( \gamma^5 \) drops out, and the respective dim-5 and dim-6 form factors become equivalent. The dark form factors presented here are applicable across the entire kinematic regime.

For massless DR, \( m_\chi = 0 \), the differential cross section can be written as,

\[
\frac{d\sigma_v}{d\Omega_{\vec{p}_f q}|_{n,l}} = \frac{\hat{\sigma}_e}{8E_R p_{\chi}} \int d\Omega_{\vec{p}_f q}|_{n,l} |F_{\chi}(q, p_{\chi})|^2. \quad (A14)
\]

It should be noted that \( \hat{\sigma}_e \) in (A2) is ill-defined for \( m_\chi \rightarrow 0 \), but in the product \( \hat{\sigma}_e |F_{\chi}(q, E_\chi)|^2 \) the mass-dependence cancels out. Hence, in practice, keeping with the usually adopted convention for the definition of \( \hat{\sigma}_e \) does not pose any obstruction.

**Appendix B: Recoil cross section on free particle**

In agreement with the previous work on the dark sector-photon interactions [38], we list here for completeness the recoil cross sections for scattering on free electrons,

\[
\begin{align*}
\text{mQ} : \quad & \frac{d\sigma}{dE_R} = e^2 \frac{\pi \alpha m_e (E_\chi^2 + 2E_\chi^2 - E_R(2E_{\chi}m_e + m_\chi^2 + m_\chi^2))}{(E_\chi^2 - m_\chi^2)^2 E_R^2 m_\chi^2}, \\
\text{MDM} : \quad & \frac{d\sigma}{dE_R} = \mu_\chi^2 \frac{(E_R - 2m_e)m_\chi^2 - 2(E_R - E_\chi)E_{\chi}m_e}{2(E_\chi^2 - m_\chi^2)E_R m_e}, \\
\text{EDM} : \quad & \frac{d\sigma}{dE_R} = a_2^2 \frac{2E_\chi m_e(E_R - E_R m_e)}{2(E_\chi^2 - m_\chi^2)E_R m_e}, \\
\text{AM} : \quad & \frac{d\sigma}{dE_R} = a_2^2 \frac{m_e (E_R - E_R(2E_\chi + m_e) + 2E_\chi^2) + m_\chi^2 (E_R - 2m_e)}{E_\chi^2 - m_\chi^2}, \\
\text{CR} : \quad & \frac{d\sigma}{dE_R} = b_2^2 \frac{E_\chi^2 m_e - E_R(2E_\chi m_e + m_\chi^2 + m_\chi^2) + 2E_\chi^2 m_e}{E_\chi^2 - m_\chi^2},
\end{align*}
\]

For the general expressions that are applicable for the analogous scattering on nuclei, see App. E in [38].
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