Exothermic Dark Matter for XENON1T Excess

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Abstract

Motivated by the recent excess in the electron recoil from XENON1T experiment, we consider the possibility of exothermic dark matter, which is composed of two states with mass splitting. The heavier state down-scatters off the electron into the lighter state, making an appropriate recoil energy required for the Xenon excess even for the standard Maxwellian velocity distribution of dark matter. Accordingly, we determine the mass difference between two component states of dark matter to the peak electron recoil energy at about \(2.5\text{ keV}\) up to the detector resolution, accounting for the recoil events over \(E_R = 2 - 3\text{ keV}\), which are most significant. We include the effects of the phase-space enhancement and the atomic excitation factor to calculate the required scattering cross section for the Xenon excess. We discuss the implications of dark matter interactions in the effective theory for exothermic dark matter and a massive \(Z'\) mediator and provide microscopic models realizing the required dark matter and electron couplings to \(Z'\).

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1 Introduction

The nature of dark matter has been a long standing mystery for astrophysics and particle physics. Weakly Interacting Massive Particles (WIMPs) have been searched for in various direct detection, cosmic-ray as well as collider experiments, and the indirect probe or constraint from the early Universe and the intensity frontier has expanded the WIMP paradigm beyond weak scale.

Quite recently, a tantalizing hint has been announced for the potential dark matter signals from the electron recoil events in the recoil energy, \( E_R = 1 - 10 \text{ keV} \), from XENON1T experiment [1]. The origin of the Xenon excess has been pondered over by particle physicists as seen from a lot of the already published articles on the arXiv for the last few days [2–6]. A simple explanation with the solar axion or the neutrino magnetic dipole moment has been put forward from the XENON1T collaboration, but both cases are inconsistent with the star cooling constraints, because the electron coupling to the axion or the neutrino magnetic dipole moment required for the Xenon excess exceeds them by the order of magnitude [1].

In this article, we consider the possibility that exothermic dark matter transits from the heavy state to the light state in the event of the scattering with the electron. This possibility was already discussed in a different context in the literature with a motivation to explain the annual modulation signal at DAMA/LIBRA [7]. In this scenario, the down-scattering of dark matter makes the recoiled energy of the electron much larger than the one inferred from the elastic scattering between the non-relativistic dark matter in the standard halo model and the electron. We discuss the details of the kinematics of exothermic dark matter and calculate the scattering cross section for the Xenon events by including the phase-space enhancement for inelastic scattering and the atomic excitation factor for a small momentum transfer between dark matter and electron. We infer the required scattering cross section for dark matter as a function of dark matter mass at a fixed recoil energy near \( E_R = 2 - 3 \text{ keV} \) up to the detector resolution. We remark that there was a recent discussion on the monochromatic electron recoil spectrum in the case of \( 3 \rightarrow 2 \) inelastic scattering between dark matter and electron [8].

We also provide the model-independent discussion with a massive \( Z' \) mediator on the effective model parameters explaining the Xenon excess and the dark matter relic density. We develop it for microscopic origins of the dark matter transition as well as the electron coupling, based on the \( Z' \)-portal and the vector–like lepton portal.

A similar idea has been discussed in Ref. [5] while this article is being finalized. Our results agree with theirs and complement with the detailed dynamics of exothermic dark matter such as the phase-space enhancement, the constraints from dark matter relic density, and concrete microscopic models realizing the scenario.
2 Exothermic dark matter and electron recoil

In this section, we begin with the kinematics for exothermic dark matter in the case of down-scattering off the electron. Then, we calculate the event rate for the electron recoil in the Xenon atoms by including the phase-space enhancement and the atomic excitation factor.

2.1 Kinematics for exothermic dark matter

We first consider the inelastic scattering between dark matter and electron, $\chi_1 e \rightarrow \chi_2 e$, where two dark fermions, $\chi_1$ and $\chi_2$, have a small mass difference, $\Delta m = m_{\chi_1} - m_{\chi_2} > 0$, which is the exothermal condition for the transition in the dark matter states. We assume that only $\chi_1$ or both $\chi_1$ and $\chi_2$ account for the observed relic density for dark matter. The up-scattering process, $\chi_2 e \rightarrow \chi_1 e$, is also possible if kinematically allowed in the tail of the dark matter velocity distribution, for instance, at $v \sim 0.1$ for $\Delta m \sim \text{keV}$, so it is would be effectively forbidden. Thus, we focus on the process, $\chi_1 e \rightarrow \chi_2 e$, in the following.

Then, we obtain the exact expression for the electron recoil energy as

$$E_R = \Delta m + E_0 \left[ 1 - \frac{m_{\chi_2} \mu_2^2}{m_{\chi_1}} \left( \frac{\sqrt{1 + \kappa}}{m_{\chi_2}} + \frac{1}{m_e} \right)^2 \right] + \frac{\mu_1^2 v^2}{m_e} \sqrt{1 + \kappa} (1 - \cos \theta)$$

where $v$ is the initial velocity of dark matter $\chi_1$,

$$\kappa = 1 + \frac{\mu_2 m_{\chi_1}}{\mu_1^2} \cdot \frac{\Delta m}{E_0}, \quad E_0 = \frac{1}{2} m_{\chi_1} v^2.$$  \hspace{1cm} (2.1)

and $\theta$ is the scattering angle in the center of mass frame, $\Delta m = m_{\chi_1} - m_{\chi_2}$, and $\mu_1, \mu_2$ are the reduced masses of the dark matter-electron system before and after scattering, given by $\mu_1 = m_e m_{\chi_1} / (m_e + m_{\chi_1})$ and $\mu_2 = m_e m_{\chi_2} / (m_e + m_{\chi_2})$, which become $\mu_1 \simeq \mu_2 \simeq m_e$ for $m_{\chi_1}, m_{\chi_2} \gg m_e$. And the exact formula for the 3-momentum transfer is also given by

$$q^2 = \mu_1^2 v^2 \left[ \left( 1 + \frac{\Delta m}{m_e} \right)^2 + 1 + \kappa - \frac{2 \mu_2}{\mu_1} \left( 1 + \frac{\Delta m}{m_e} \right) \sqrt{1 + \kappa \cos \theta} \right].$$ \hspace{1cm} (2.2)

If $\Delta m = 0$, the standard formulas for elastic scattering are recovered as $E_R = \frac{\mu_1^2 v^2}{m_e} (1 - \cos \theta) = \frac{q^2}{2m_e}$.

It is also illuminating to express the momentum transfer for down-scattering of dark matter from the energy conservation in terms of the DM velocity, the recoil energy, the mass difference, and the angle $\zeta$ between the dark matter velocity $\vec{v}$ and the momentum transfer $\vec{q}$, as follows,

$$q_{\pm} = m_{\chi_1} v \cos \zeta \pm \sqrt{m_{\chi_1} (m_{\chi_1} \cos^2 \zeta - \Delta m) v^2 - 2 m_{\chi_2} (E_R - \Delta m)}. \hspace{1cm} (2.3)$$
Then, the range of the momentum transfer is given by the following: for $E_R > \Delta m$,

$$q_\pm = m_{\chi_1}v \pm \sqrt{m_{\chi_1}^2 v^2 - 2(E_R - \Delta m)}; \quad (2.5)$$

for $E_R < \Delta m$,

$$q_\pm = \pm m_{\chi_1}v \pm \sqrt{m_{\chi_1}^2 v^2 - 2(E_R - \Delta m)}. \quad (2.6)$$

Then, for $\Delta m = 0$, we can recover the values for elastic scattering from eq. (2.5) as $q_\pm = m_{\chi_1}v \pm \sqrt{m_{\chi_1}^2 v^2 - 2m_{\chi_1}E_R}$.

For $\Delta m \ll m_e \ll m_{\chi_1}$, we can approximate eq. (2.2) as

$$\kappa \simeq \frac{\Delta m}{2m_e v^2} \simeq 2.2 \times 10^4 \left(\frac{220 \text{ km/s}}{v}\right)^2 \left(\frac{\Delta m}{3 \text{ keV}}\right). \quad (2.7)$$

Then, for $\kappa \gg 1$, the electron recoil energy and the momentum transfer become simplified to

$$E_R \simeq \Delta m - \sqrt{\Delta m}m_e v^2 \left(1 + \frac{m_e}{2m_{\chi_1}} \sqrt{\kappa}\right) + \sqrt{\kappa} m_e v^2 (1 - \cos \theta)$$

$$q^2 \simeq 2m_e\Delta m \left(1 - \frac{2}{\sqrt{\kappa}} \cos \theta\right). \quad (2.8)$$

Therefore, $q^2 \simeq 2m_eE_R$ holds as in the case with $\Delta m = 0$, but either the recoil energy or the momentum transfer depend little on the dark matter velocity and the scattering angle. The above result is consistent with eqs. (2.5) or (2.6), and it shows a tiny momentum transfer, for which the atomic excitation factor becomes important [3, 5, 9, 10], as will be discussed later.

On the other hand, for $\Delta m \ll m_e \ll m_{\chi_1}$ and $\kappa \ll 1$, we have

$$E_R \simeq \Delta m + m_e v^2 (1 - \cos \theta), \quad (2.10)$$

$$q^2 \simeq 2m_e^2 v^2 (1 - \cos \theta). \quad (2.11)$$

In this case, the recoil energy is bounded by $2m_e v^2 \simeq 5.5 \times 10^{-4} \text{ keV} \gg \Delta m$ for $v \sim 220 \text{ km/s}$, which is too small to account for the Xenon experiment. Thus, we focus on the regime with $\kappa \gg 1$ in the following discussion, which is sufficient for explaining the Xenon excess.

### 2.2 The event rate for electron recoil

We begin with the general expression for the event rate per target mass [11], given by

$$dR = \frac{\rho_{\chi_1} v}{m_{\chi_1} m_T} d\sigma f_1(v) dv \quad (2.12)$$
where \( m_T \) is the target nucleus mass and \( f_1(v) = \frac{4v^2}{v_0^2\sqrt{\pi}} e^{-v^2/v_0^2} \) with \( v_0 = 220 \text{ km/s} \) for the Maxwellian velocity distribution of dark matter and \( \int_0^\infty f_1(v) dv = 1 \), and \( \rho_{\chi_1} \) is the local energy density of dark matter, which is given by \( \rho_{\chi_1} = 0.4 \text{ GeV/cm}^3 \) if \( \chi_1 \) occupies the full dark matter.

We define the total scattering cross section as

\[
\sigma_e = \int_{q_-}^{q_+} \frac{dq}{dq^2} \frac{d\sigma(q = 1/a_0)}{dq^2} dq^2, \tag{2.13}
\]

where the differential cross section in the integrand is evaluated at \( q = 1/a_0 \), and write down the differential cross section in terms of the integrated atomic excitation factor, \( K_{\text{int}}(q) \), as follows,

\[
\frac{d\sigma}{dq^2} = \frac{\sigma_e}{q_+^2 - q_-^2} K_{\text{int}}(E_R) P^2(v) \tag{2.14}
\]

where \( P^2(v) \) is the phase space factor, which is unity for the elastic scattering, and

\[
K_{\text{int}}(E_R) = \int_{q_-}^{q_+} a_0^2 q' dq' K(E_R, q') \tag{2.15}
\]

with \( a_0 = \frac{1}{\alpha m_e} \) being the Bohr radius and \( K(E_R, q') \) being the atomic excitation factor. We note that the total recoil energy is also deposited significantly near \( E_R \sim \text{keV} \) to ionize the electrons bound to the Xenon atoms, and the atomic excitation factor can be important for a small momentum transfer \([3,5,9,10]\).

Then, we obtain the differential event rate per target mass as

\[
\frac{dR}{dE_R} = \frac{\sigma_e \rho_{\chi_1}}{m_{\chi_1} m_T} K_{\text{int}}(E_R) \int_{v_{\text{min}}}^{\infty} \frac{v P^2(v)}{q_+^2 - q_-^2} f_1(v) dv \tag{2.16}
\]

where \( v_{\text{min}} \) is the minimum velocity of dark matter required for a given recoil energy \( E_R \). As a result, we get the event rate per detector as

\[
R_D = M_T \int_{E_T}^{\infty} \frac{dR}{dE_R} dE_R \tag{2.17}
\]

where \( E_T \) is the detector threshold energy and \( M_T \) is the fiducial mass of the detector, given by \( M_T \sim 4.2 \times 10^{27} (M_T/\text{tonne}) m_T \) for Xenon.

Now we apply the general result in eq. (2.16) for the case with down-scattering dark matter. We take \( \kappa \gg 1 \), for which the recoil energy is appreciable. In this case, we obtain the phase space factor \( P^2(v) \) in eq. (2.14) as

\[
P^2(v) \simeq \sqrt{1 + \frac{2\Delta m}{\mu_1 v^2}} \simeq \sqrt{\frac{2\Delta m}{m_e}} \frac{1}{v}. \tag{2.18}
\]
Then, using eq. (2.16) with \( q^2 \approx 2m_e \Delta m \left( 1 \pm \frac{2}{\sqrt{\kappa}} \right) \) from eq. (2.9), we obtain the differential event rate for \( E_- < E_R < E_+ \) with \( E_\pm = \Delta m \left( 1 \pm \frac{2}{\sqrt{\kappa}} \right) \) as

\[
\frac{dR}{dE_R} \approx \frac{\sigma_e \rho_{\chi_1}}{2m_e m_{\chi_1} m_T} K_{\text{int}}(E_R) \int_{v_\text{min}}^{v_\text{max}} \frac{f_1(v)}{v} dv \theta(E_R - E_-) \theta(E_+ - E_R) \tag{2.19}
\]

where \( v_\text{min} = 0, \ v_\text{max} = \sqrt{\frac{2\Delta m}{m_e}} \) at \( \kappa = 1 \). For \( E_+ - E_- \ll E_\pm \), we can approximate \( \theta(E_R - E_-) \theta(E_+ - E_R) \approx (E_+ - E_-) \delta(E_R - \Delta m) \). Therefore, we can rewrite eq. (2.19) as

\[
\frac{dR}{dE_R} \approx \frac{(2\Delta m)}{m_e}^{1/2} \frac{\sigma_e \rho_{\chi_1}}{m_{\chi_1} m_T} K_{\text{int}}(E_R) \delta(E_R - \Delta m) \int_{0}^{v_\text{max}} f_1(v) dv \tag{2.20}
\]

We note that for \( \Delta m \gg 1.3 \times 10^{-4} \text{keV} \), we have \( v_\text{max} \gg v_0 \), resulting in \( \int_{0}^{v_\text{max}} f_1(v) dv \approx 1 \). Therefore, we find that there is no Boltzmann suppression due to an enhancement factor \( P^2(v) \) in eq. (2.18), as compared to the case with elastic scattering.

Consequently, from eq. (2.17) with eq. (2.20) and \( F^2(E_R, \Delta m) \approx 1 \), we get the total event rate per Xenon detector with \( m_T = m_{Xe} \) as

\[
R_D \approx \left( \frac{M_T \sigma_e \rho_{\chi_1}}{m_{\chi_1} m_T} \right) \left( \frac{2\Delta m}{m_e} \right)^{1/2} K_{\text{int}}(\Delta m) \approx 50 \left( \frac{M_T}{\text{tonne-yr}} \right) \left( \frac{K_{\text{int}}(\Delta m)}{19.4} \right) \left( \frac{\rho_{\chi_1}}{0.4 \text{GeV cm}^{-3}} \right) \times \left( \frac{\sigma_e/m_{\chi_1}}{1.6 \times 10^{-44} \text{cm}^2/\text{GeV}} \right) \left( \frac{\Delta m}{2.5 \text{keV}} \right)^{1/2} \tag{2.21}
\]

where we have used the normalization for the integrated atomic excitation factor from Ref. [5].

For comparison to the experimental data, the mono-energetic event rate can be convoluted with the detector resolution by

\[
\frac{dR_D}{dE_R} = \frac{R_D}{\sqrt{2\pi} \sigma} e^{-\frac{(E_R - \Delta m)^2}{(2\sigma^2)}} \alpha(E) \tag{2.22}
\]

where \( \sigma \) is the detector resolution, which varies between 20% at \( E = 2 \text{keV} \) and 6% at \( E = 30 \text{keV} \), and \( \alpha(E) \) is the signal efficiency [1]. For \( E_R = 2 - 10 \text{keV} \), the signal efficiency is given by \( \alpha(E) \approx 0.7 - 0.9 \) [1].

The XENON1T excess is most significant from the electrons at \( E_R = 2 - 3 \text{keV} \) with the detector resolution being about \( \sigma = 0.4 \text{keV} \), so we take the recoil energy of the mono-energetic electron in our model to be \( E_R \approx \Delta m \approx 2.5 \text{keV} \). Moreover, from \( \alpha(E) \approx 0.8 \) at \( E_R \approx 2.5 \text{keV} \), we need to rescale the total event rate per detector in eq. (2.21) by a factor 0.8. Therefore, taking into account the total exposure in XENON1T for SR1, which is 0.65 tonne-yr [1], and for \( \rho_{\chi_1} = 0.4 \text{GeV cm}^{-3} \), we can get about 50 events near \( E_R = 2 - 3 \text{keV} \) for the XENON1T electron recoil events for \( \sigma_e/m_{\chi_1} \approx 3.2 \times 10^{-44} \text{cm}^2/\text{GeV} \).
3 The effective theory for exothermic dark matter

We continue to discuss the effective theory for exothermic dark matter in the presence of a massive $Z'$ mediator and constrain the parameter space for the $Z'$ couplings and the mass parameters from the Xenon excess. For completeness, we also provide the formulas for dark matter annihilation cross sections in the effective theory and comment on the compatibility of the Xenon excess with the correct relic density.

3.1 The effective interactions and the Xenon excess

We consider two Majorana dark matter fermions, $\chi_1$ and $\chi_2$, with different masses, $m_{\chi_1} > m_{\chi_2}$, and a massive dark gauge boson $Z'$ with mass $m_{Z'}$. We take the effective Lagrangian with $Z'$ couplings to dark fermions, electron and electron neutrino, in the following form,

$$\mathcal{L}_{\text{eff}} = (g Z' Z'_{\mu} \bar{\chi}_2 \gamma^\mu (v_\chi + a_e \gamma^5) \chi_1 + \text{h.c.}) + g Z' Z'_\mu \bar{e} (v_e + a_e \gamma^5) e + g Z' Z'_\mu \bar{\nu} \gamma^\mu (v_\nu + a_\nu \gamma^5) \nu \tag{3.1}$$

where $v_i, a_i$ with $i = \chi, e, \nu$ are constant parameters. In the next subsection, we will show a microscopic model for the above effective interactions. For $\Delta m < m_{Z'}$ and $\Delta m < 2m_e$, there is no tree-level decay process for the dark matter fermion $\chi_1$.

For $a_e \neq 0$, the effective vertex interaction for one $Z'$ and two photons, $Z'_{\mu} (q) - A_\lambda (p_1) - A_\sigma (p_2)$, with $q = q_1 + q_2$, is induced by the electron loops, taking the following form for $q^2 \ll 4m_e^2$ [14],

$$\Gamma^{\mu\lambda\sigma} (q, q_1, q_2) \simeq \left( \epsilon^{q_1, \lambda\sigma\mu} - \epsilon^{q_2, \lambda\sigma\mu} \right) \cdot \frac{a_e e^2 g_{Z'}}{4\pi^2} \left( 1 + \frac{q^2}{12m_e^2} \right). \tag{3.2}$$

Then, the heavy state of dark matter can decay into the lighter state of dark matter and two photons, with the decay rate given by

$$\Gamma (\chi_1 \rightarrow \chi_2 \gamma\gamma) \simeq \frac{a_e^2 (v_\chi^2 + a_\chi^2) e^4 g_{Z'}^2 (\Delta m)^7}{3072\pi^7} \frac{m_{\chi_1}^2 m_{Z'}^4}{m_{\chi_1}^4 m_{Z'}^4}. \tag{3.3}$$

Therefore, in this case, the lifetime of the dark fermion $\chi_1$ is much longer than the age of the Universe for perturbative effective couplings. However, the diffuse X-ray background [15] puts the bound on the lifetime of the dark fermion $\chi_1$ to $\tau_{\chi_1} > 10^{24}$ sec, which gives rise to

$$|a_e| g_{Z'} \sqrt{v_\chi + a_\chi} < 0.3 \left( \frac{2.5 \text{ keV}}{\Delta m} \right)^{7/2} \left( \frac{m_{\chi_1}}{0.3 \text{ GeV}} \right) \left( \frac{m_{Z'}}{1 \text{ GeV}} \right)^2. \tag{3.4}$$

As a result, a nonzero perturbative axial vector coupling to the electron is safe from the X-ray bounds for exothermic dark matter and light $Z'$ mediator. As we will discuss in the next section, some microscopic models with vector-like leptons can induce a suppressed axial vector coupling to the electron.
Moreover, there is another loop process for the three-photon decay channel, \( \chi_1 \rightarrow \chi_2 + 3\gamma \), but the corresponding decay rate is highly suppressed by \( \Gamma \propto (\Delta m)^3/(m_{Z'}^4 m_\chi^8) \) [5], thus being consistent with X-ray bounds [15].

If the neutrino couplings to \( Z' \) are nonzero, the dark matter fermion \( \chi_1 \) would decay into a neutrino pair via the off-shell \( Z' \) gauge boson, which is bounded by the lifetime of dark matter for explaining the XENON1T electron recoil excess. The decay rate of the dark fermion \( \chi_1 \) is given [13] by

\[
\Gamma(\chi_1 \rightarrow \chi_2 \nu \bar{\nu}) \simeq N_\nu G_F^2 (\Delta m)^5 / 30 \pi^3 (v_\chi^2 + 3a_\chi^2)(v_\nu^2 + a_\nu^2)
\]

where \( N_\nu \) is the number of neutrinos coupled to \( Z' \), and \( G_F' \equiv g_{Z'}^2 / (\sqrt{2} m_{Z'}^2) \). Then, for \( \Delta m = 2.5\text{ keV} \) and \( v_\chi = -a_\chi = \frac{1}{2} \), the lifetime of the dark fermion \( \chi_1 \) is longer than the age of the Universe, as far as

\[
G_F' \sqrt{N_\nu (v_\nu^2 + a_\nu^2)} < 2.4 \times 10^{-6} \text{ GeV}^{-2}.
\]

Nonetheless, neutrino experiments such as Super-Kamiokande [16] constrain the lifetime of dark matter to \( \tau_{\chi_1} > 10^{24} \text{ sec} \) [17], similarly to the X-ray bounds, thus being much stronger than the above bound. As will be shown in the next section, the effective neutrino couplings are induced in the case of \( Z' \) portal with a gauge kinetic mixing, but they can be sufficiently small to be consistent with both the neutrino bounds and the Xenon excess.

For \( m_e, m_{\chi_1}, m_{Z'} \gg q \sim m_e \Delta m, m_e \ll m_{\chi_1} \), and \( \Delta m \ll m_{\chi_1} \), the total scattering cross section for \( \chi_1 e \rightarrow \chi_2 e \), up to the phase space factor \( P^2(v) \) in eq. (2.18), is given by

\[
\sigma_e \simeq \frac{v_\chi^2 v_e^2 g_{Z'}^2 m_{\chi_1}^2}{\pi m_{Z'}^2}
\]

where \( e \) is the electromagnetic coupling. Here, we chose the \( Z' \) couplings to be consistent with the dilepton bounds from BaBar, \( |v_e| g_{Z'} \lesssim 10^{-4} e \) for \( 0.02 \text{ GeV} < m_{Z'} < 10.2 \text{ GeV} \) [18], or the bound from mono-photon + MET from BaBar [19], \( |v_e| g_{Z'} \lesssim (4 \times 10^{-4} - 10^{-3}) e \) for \( m_{Z'} < 8 \text{ GeV} \). Thus, we need to have nonzero vector couplings to both dark matter and electron for the scattering cross section without velocity suppression. In order to explain the XENON1T electron recoil events near \( E_R = 2 - 3 \text{ keV} \) in our model, we take \( \Delta m \simeq 2.5 \text{ keV} \) and the required scattering cross section gives rise to the following useful formula,

\[
\left(\frac{v_\chi g_{Z'}}{0.6}\right)^2 \left(\frac{v_e g_{Z'}}{10^{-4} e}\right)^2 \left(\frac{1 \text{ GeV}}{m_{Z'}}\right)^4 \left(\frac{0.3 \text{ GeV}}{m_{\chi_1}}\right)^2 \left(\frac{\Omega_{\text{DM}}}{\Omega_{\chi_1}}\right) \simeq 1
\]

where \( \Omega_{\text{DM}} \) is the observed total abundance of dark matter and \( \Omega_{\chi_1} \) is the abundance of the dark fermion \( \chi_1 \). Therefore, light dark matter and \( Z' \) mediator are favored by the explanation of the Xenon excess with exothermic dark matter. As we scale up the \( Z' \) gauge coupling, we can take a larger value of \( m_{Z'}^4 m_{\chi_1} \) in order to maintain the number of the electron recoil events.
3.2 Dark matter annihilation and relic density

Dark matter fermions $\chi_1$ and $\chi_2$ can co-annihilate into a pair of electrons as well as into a pair of $Z'$ gauge bosons if kinematically allowed. Then, taking $\Delta m \ll m_{\chi_i}$ and ignoring the lepton masses, the total annihilation cross section for $\chi_1 \bar{\chi}_2 \rightarrow e^- e^+$ and $\chi_1 \bar{\chi}_1, \chi_2 \bar{\chi}_2 \rightarrow Z' Z'$ is given by

$$\langle \sigma v \rangle = \frac{1}{2} \langle \sigma v \rangle_{\chi_1 \bar{\chi}_2 \rightarrow e^- e^+} + \frac{1}{2} \langle \sigma v \rangle_{\chi_1 \bar{\chi}_2 \rightarrow Z' Z'}$$

with

$$\langle \sigma v \rangle_{\chi_1 \bar{\chi}_2 \rightarrow e^- e^+} = \frac{g_{Z'}^4 \left[ \chi_1 \right]}{\pi} \left[ v_e^2 + a_e^2 + N_\nu (a_\nu^2 + a_\nu^2) \right] \frac{m_{\chi_1}^2}{(m_{Z'}^2 - 4m_{\chi_1}^2)^2 + \Gamma_{Z'}^2 m_{Z'}^2},$$

$$\langle \sigma v \rangle_{\chi_1 \bar{\chi}_1, \chi_2 \bar{\chi}_2 \rightarrow Z' Z'} = \frac{g_{Z'}^4}{4\pi} \left[ v_{\chi}^4 + a_{\chi}^4 + 2v_{\chi}^2 a_{\chi}^2 \left( 4 \frac{m_{\chi_1}^2}{m_{Z'}^2} - 3 \right) \right] \frac{m_{\chi_1}^2}{(m_{Z'}^2 - 2m_{\chi_1}^2)^2} \left( 1 - \frac{m_{Z'}^2}{m_{\chi_1}^2} \right)^{3/2}$$

We note that the contributions coming from $a_\chi$ to the annihilation cross section are $p$-wave suppressed. Since we need $v_\chi \neq 0$ for explaining the Xenon electron excess, the $p$-wave annihilations are sub-dominant.

For light dark matter with sub-GeV mass, once $\chi_1 \bar{\chi}_1, \chi_2 \bar{\chi}_2 \rightarrow Z' Z'$ is open, the resultant annihilation cross section would be too large for a sizable $g_{Z'}$ to account for the correct relic density. Thus, in this case, we can take $m_{\chi_1} < m_{Z'}$ such that the annihilation of the dark matter fermion $\chi_1$ into a pair of $Z'$ is forbidden at zero temperature, but it is open in the tail of the Boltzmann distribution at a finite temperature during freeze-out [12]. Then, from the detailed balance condition for the forbidden channels, the effective annihilation cross section for the forbidden channels, $\chi_1 \bar{\chi}_1, \chi_2 \bar{\chi}_2 \rightarrow Z' Z'$ becomes

$$\langle \sigma v \rangle_{\chi_1 \bar{\chi}_1, \chi_2 \bar{\chi}_2 \rightarrow Z' Z'} = \left( \frac{n_{eq}^{\chi_1}}{n_{eq}^{\chi_1}} \right)^2 \langle \sigma v \rangle_{Z' Z' \rightarrow \chi_1 \bar{\chi}_1, \chi_2 \bar{\chi}_2}$$

with

$$\langle \sigma v \rangle_{Z' Z' \rightarrow \chi_1 \bar{\chi}_1, \chi_2 \bar{\chi}_2} = \frac{4g_{Z'}^4}{9m_{Z'}^2} \left[ v_{\chi}^4 + a_{\chi}^4 \left( \frac{m_{Z'}^2 + 2m_{\chi_1}^2}{m_{Z'}^2 + m_{\chi_1}^2} \right) + 2v_{\chi}^2 a_{\chi}^2 \left( \frac{3m_{Z'}^2 - m_{\chi_1}^2}{m_{Z'}^2 + m_{\chi_1}^2} \right) \right] \left( 1 - \frac{m_{Z'}^2}{m_{\chi_1}^2} \right)^{3/2}$$

Here, for $\Delta m \ll m_{\chi_1}$, the Boltzmann suppression factor can be approximated to

$$\left( \frac{n_{eq}^{\chi_1}}{n_{eq}^{\chi_1}} \right)^2 \simeq \left( \frac{n_{eq}^{\chi_1}}{n_{eq}^{\chi_1}} \right)^2 \simeq \frac{9}{4} \left( \frac{m_{Z'}}{m_{\chi_1}} \right)^3 e^{-2(m_{Z'} - m_{\chi_1})/T}.$$

The forbidden channels are important for obtaining the correct relic density for light dark matter, because the strong annihilation cross section can be compensated by the Boltzmann suppression factor [12].

As a consequence, the dark matter number density is given by $n_{DM} = n_{\chi_1} + n_{\chi_2}$ with $n_{\chi_1} \simeq n_{\chi_2}$, so the corresponding relic abundance is determined as

$$\Omega_{DM} h^2 = 0.12 \left( \frac{10.75}{g_*(T_f)} \right)^{1/2} \left( \frac{x_f}{20} \right) \left( \frac{4.3 \times 10^{-9} \text{GeV}^{-2}}{x_f \int_{x_f}^{\infty} x^{-2} \langle \sigma v \rangle} \right)$$

(3.15)
where \( x_f = m_{\chi_1}/T_f \) at freeze-out temperature. Therefore, parametrizing the effective annihilation cross section by \( \langle \sigma v \rangle = \frac{\alpha_{\text{eff}}}{m_{\chi_1}^2} \), we can achieve a correct relic density, provided that

\[
m_{\chi_1} \simeq 150 \text{ MeV} \left( \frac{\alpha_{\text{eff}}}{10^{-5}} \right). \tag{3.16}\]

As a result, we can get the correct relic density by taking a small effective coupling \( \alpha_{\text{eff}} \) from small SM couplings, \( v_e, v_\nu \), or due to the Boltzmann-suppression, \( \alpha_{\text{eff}} \sim g_{Z'}^2 e^{-x_f(m_{Z'} - m_{\chi_1})} \), for \( g_{Z'} = 0.6 \) and \( \frac{m_{Z'}}{m_{\chi_1}} - 1 \simeq 0.4 \), being consistent with the constraint from the Xenon excess with \( \Omega_{\chi_1} \simeq \frac{1}{4}\Omega_{\text{DM}} \) in eq. (3.8).

In the above discussion, we focused on the annihilation channels in the minimal scenario for the Xenon excess. However, if the aforementioned annihilation channels with \( Z' \) interactions are not sufficient for a correct relic density, due to small \( Z' \) couplings to the SM, we can also consider the dark matter self-interactions for dark matter annihilation, in particular, as in the case of SIMP dark matter [20, 21] where sub-GeV light dark matter is a natural outcome of the \( 3 \rightarrow 2 \) annihilations with strong self-interactions.

### 4 Microscopic models

In this section, we propose a microscopic model for exothermic dark matter by taking two left-handed dark fermions, \( \psi_1 \) and \( \psi_2 \), with opposite charges, +1 and −1, under the dark \( U(1)' \) symmetry. We also introduce a dark Higgs \( \phi \) with charge −2 under the \( U(1)' \). We assume that all the SM particles are neutral under the \( U(1)' \), but dark matter can communicate with the SM through 1) the gauge kinetic mixing, \( \sin \xi \), or 2) the mixing between electron and an extra vector-like lepton. We first discuss the dark matter interactions and proceed to derive the effective interactions for the electron in each case of portal models.

The Lagrangian for the dark sector is given, as follows,

\[
\mathcal{L} = -\frac{1}{4} F_{\mu \nu}' F'^{\mu \nu} + |D_\mu \phi|^2 - V(\phi, H) \\
+ i \bar{\psi}_1 \gamma^\mu D_\mu \psi_1 + i \bar{\psi}_2 \gamma^\mu D_\mu \psi_2 \\
- m_\psi \psi_1 \psi_2 - y_1 \phi \psi_1 \psi_1 - y_2 \phi^* \psi_2 \psi_2 + \text{h.c.} \tag{4.1}\]

where \( F'_{\mu \nu} = \partial_\mu Z'_\nu - \partial_\nu Z'_\mu \), \( B_{\mu \nu} \) is the field strength tensor for the SM hypercharge, the covariant derivatives are \( D_\mu \phi = (\partial_\mu + 2ig_{Z'} Z'_\mu) \phi \), \( D_\mu \psi_1 = (\partial_\mu - ig_{Z'} Z'_\mu) \psi_1 \), \( D_\mu \psi_2 = (\partial_\mu + ig_{Z'} Z'_\mu) \psi_2 \), \( m_\psi \) is the Dirac mass for dark fermions, and \( y_{1,2} \) are the Yukawa couplings for dark fermions. Here, \( V(\phi, H) \) is the scalar potential for the singlet scalar \( \phi \) and the SM Higgs.

After the dark Higgs gets a VEV as \( \langle \phi \rangle = v_\phi \), the \( Z' \) gauge boson receives mass \( m_{Z'} = 2\sqrt{2}g_{Z'} v_\phi \), and there appears a mass mixing between \( \psi_1 \) and \( \psi_2 \). Then, diagonalizing the mass matrix for the dark fermions, we get the mass eigenvalues and the mixing matrix, as follows,

\[
m^2_{\chi_{1,2}} = m^2_\psi + 2(y_1^2 + y_2^2) v_\phi^2 \pm 2 \sqrt{(y_1^2 - y_2^2)^2 v_\phi^4 + (y_1 + y_2)^2 v_\phi^2 m^2_\psi}, \tag{4.2}\]
and
\[
\begin{pmatrix}
\chi_1 \\
\chi_2
\end{pmatrix} = \begin{pmatrix}
\cos \theta & -\sin \theta \\
\sin \theta & \cos \theta
\end{pmatrix} \begin{pmatrix}
\psi_2 \\
\psi_1
\end{pmatrix} \tag{4.3}
\]
with
\[
\sin 2\theta = -\frac{4(y_1 + y_2)v_\phi m_\psi}{m_{\chi_2}^2 - m_{\chi_1}^2}. \tag{4.4}
\]

For simplicity, we take \(y_1 = y_2\), then the mass eigenvalues become \(m_{\chi_1,2} = m_\psi \pm 2y_1 v_\phi\) and the mixing angle is given by \(\theta = \frac{\pi}{4}\). Then, the mass difference can be small as far as \(2|y_1|v_\phi \ll m_\psi\). There was a recent discussion on the microscopic model for exothermic dark matter with a complex scalar field where the breaking of dark \(U(1)'\) makes a small mass splitting between the real scalar fields [5].

For instance, for \(m_\psi \sim 1\) GeV, we need \(y_1 \sim 1.5 \times 10^{-6}\). As a result, including the \(Z'\)-portal couplings to the SM fermions for a small gauge kinetic mixing, we summarize the \(Z'\) gauge interactions as follows,
\[
\mathcal{L}_{DM} = g_{Z'} Z'_\mu \left( \bar{\chi}_1 \gamma^\mu P_L \chi_2 + \bar{\chi}_2 \gamma^\mu P_L \chi_1 \right). \tag{4.5}
\]
Then, we obtain the effective dark matter couplings in the Lagrangian (3.1) as
\[
v_\chi = -a_\chi = \frac{1}{2}. \tag{4.6}
\]

As a result, we can realize the transition interactions between two states of dark matter via the \(Z'\) mediator, that are necessary for explaining the Xenon excess.

### 4.1 \(Z'\)-portal

In the presence of a gauge kinetic mixing,
\[
\mathcal{L}_{\text{kin-mix}} = -\frac{1}{2} \sin \xi B_{\mu \nu} F^{\mu \nu}, \tag{4.7}
\]
the mixing between \(Z'\) and \(Z\) gauge bosons gives rise to the \(Z'\) gauge interactions to the SM as
\[
\mathcal{L}_{\text{eff,1}} = -\varepsilon Z'_\mu \left( \bar{e} \gamma^\mu e + \frac{m_{Z'}}{2e W^2 m_Z^2} \bar{\nu} \gamma^\mu P_L \nu \right) + \cdots \tag{4.8}
\]
where the ellipse denotes the electromagnetic and neutral current interactions for the rest of the SM fermions. Therefore, there are not only electron couplings but also neutrino couplings, although the latter being further suppressed by \(m_{Z'}/m_Z^2\) [22]. Consequently, we can identify the effective couplings in the Lagrangian (3.1), as follows,
\[
v_e = -\frac{\varepsilon g_{Z'}}{g_{Z'}}, \quad a_e = 0, \quad v_\nu = -a_\nu = -\frac{\varepsilon m_{Z'}^2}{4e W^2 g_{Z'} m_Z^2}. \tag{4.9}
\]
Then, the lifetime of the dark fermion $\chi_1$ is given by

$$\tau_{\chi_1} = \frac{1}{\Gamma(\chi_1 \rightarrow \chi_2 \nu \bar{\nu})} = \left(10^{-4}e\right)^2 \left(\frac{2.5 \text{ keV}}{\Delta m}\right)^5 8.9 \times 10^{24} \text{ sec}, \quad (4.10)$$

which can be consistent with the bound from neutrino experiments that we discussed in the previous section.

### 4.2 Vector-like lepton portal

We introduce an extra vector-like charged lepton $E$ which has charge $-2$ under the $U(1)'$ but is singlet under the $SU(2)_L$. Then, the mixing between the SM right-handed electron and the vector-like lepton is given by

$$\mathcal{L}_{VL} = -M_E \bar{E}E - (y_E v \bar{E} e_R + \text{h.c.}) \quad (4.11)$$

As a consequence, the mass matrix for the electron and the vector-like lepton takes

$$M_e = \begin{pmatrix} m_e & 0 \\ y_E v & M_E \end{pmatrix}. \quad (4.12)$$

Then, the mass eigenvalues are given by

$$m_{f_1,2}^2 = \frac{1}{2} \left( m_e^2 + M_E^2 + y_E^2 v^2 \mp \sqrt{(m_e^2 + y_E^2 v^2 - M_E^2)^2 + 4y_E^2 v^2 M_E^2}\right). \quad (4.13)$$

On the other hand, the mixing angles for the right-handed electrons and the left-handed electrons are given [23], respectively, by

$$\sin(2 \theta_R) = -\frac{2y_E v \phi M_E}{m_{f_1}^2 - m_{f_2}^2},$$

$$\sin(2 \theta_L) = \frac{m_e}{m_{f_1} m_{f_2}} \sin(2 \theta_R). \quad (4.15)$$

Therefore, for $m_e, y_E v \phi \ll M_E$ and $(y_E v \phi / M_E)^2 \ll (m_e / M_E)$, the mass eigenvalues are approximated to $m_{f_1} \sim m_e$ and $m_{f_2} \sim M_E$, and the mixing angles become $\theta_R \sim \frac{2y_E v \phi}{M_E}$ and $\theta_L \sim \frac{m_e}{M_E} \theta_R$. Given the experimental bound on the vector-like charged lepton from LEP and LHC, $M_E \gtrsim 100 \text{ GeV}$, we have $\frac{m_e}{M_E} \lesssim 5 \times 10^{-6}$, so we can ignore the mixing for the left-handed electrons. But, for $m_{f_1} \sim m_e$, the mixing angle for the right-handed electrons is bounded by $\theta_R \lesssim \sqrt{\frac{m_e}{M_E}}$, thus $\theta_R$ can be as large as $2.2 \times 10^{-3}$. On the other hand, we note that the suppressed mixing angle $\theta_L$ for the left-handed electron is consistent with electroweak precision data.

Consequently, we get the following effective interactions for $Z'$,

$$\mathcal{L}_{\text{eff},\text{II}} = -2g_{Z'} Z'_\mu \left( \bar{E} \gamma^\mu P_R E + \theta_R^2 \bar{e} \gamma^\mu P_R e - \theta_R \bar{E} \gamma^\mu P_R e - \theta_R \bar{e} \gamma^\mu P_R E \right)$$
\[- \frac{g}{2cW} Z_\mu \left( \bar{e}_L \gamma^\mu P_L e + \theta_L \bar{E} \gamma^\mu P_L e + \theta_L \bar{\ell} \gamma^\mu P_L E + \theta_L^2 \bar{\ell} \gamma^\mu P_L E \right) \]
\[- \frac{g}{\sqrt{2}} \theta_L \bar{E} \gamma^\mu P_L \nu W^+_{\mu} + \text{h.c.} \]  

(4.16)

Therefore, we can identify the effective couplings in the Lagrangian (3.1), as follows,

\[ v_e = a_e = -\theta^2_R, \quad v_\nu = a_\nu = 0. \]  

(4.17)

As a result, we obtain the necessary electron coupling to \( Z' \) for explaining the Xenon excess through the small mixing between the SM right-handed electron and the vector-like lepton. In the model with vector-like lepton portal, however, there is no direct coupling between \( Z' \) and neutrinos. Even the \( Z' - \ell - E \) vertex with \( E \) decaying into the SM particles does not make the dark fermion \( \chi_1 \) to decay, provided that \( \Delta m < 2m_\ell \) is chosen.

5 Conclusions

We proposed exothermic dark matter to explain the recent electron excess reported by XENON1T experiment. Even for a small dark matter velocity, as known from the standard Maxwellian velocity distribution of dark matter, we achieved the appropriate recoil electron recoil energy at about 2.5 keV by considering the down-scattering of the heavier dark matter state off the electron into the lighter state. Thus, we showed that about 50 recoil events over \( E_R = 2 - 3 \) keV, which are most significant, can be explained in this scenario up to the detector resolution.

Including the effects of the phase-space enhancement and the atomic excitation factor, we derived the required scattering cross section for the Xenon excess to be about \( \sigma_e \sim 10^{-44} \text{cm}^2 \) for sub-GeV light dark matter. We took the effective theory approach for exothermic dark matter with a massive \( Z' \) mediator and discussed the implications of the Xenon excess for dark matter interactions to \( Z' \) and the dark matter relic density. We also provided microscopic models with \( Z' \) portal and vector-like portal, realizing the required dark matter and electron couplings to \( Z' \), while the heavier state of dark matter is long-lived enough to satisfy the bounds from the X-ray or the neutrino experiments.

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