

Direct-detection constraints on inelastic dark matter with a scalar mediator

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We phenomenologically calculate direct-detection constraints on inelastic dark matter (DM) for a scalar portal scenario with leptophilic couplings. The p-wave velocity suppression of the annihilation cross section of scalar-mediated inelastic Dirac DM implies the opening of viable regions of DM parameter space in the MeV-GeV mass range. Xenon-based experiments can provide a constraints on scalar-mediated inelastic fermion dark matter for sub-MeV mass splitting, via endothermic and exothermic spin-independent DM-electron scattering. To estimate the relevant constraints, we use public data from the XENON1T, PandaX-4T, and LUX-ZEPLIN liquid-xenon experiments that measure ionization electron signals.

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I. INTRODUCTION

In recent decades, a broad range of cosmological and astrophysical observations has indicated that roughly 85% of the matter content of the Universe cannot be accounted for by the known particle content of the Standard Model [1]. The existence of dark matter is inferred from gravitational phenomena on multiple scales, including galactic rotation curves, gravitational lensing, galaxy-cluster dynamics, and the cosmic microwave background [2]. By contrast, no nongravitational interaction between dark matter and standard model fields has yet been established experimentally.

The lack of a confirmed nongravitational signal implies that the particle properties of dark matter remain largely unconstrained beyond its gravitational effects on the visible sector. This situation motivates a systematic exploration of dark-matter scenarios capable of accounting for the observed phenomena attributed to dark matter. As a result, deriving phenomenological constraints on the parameter space of dark-matter models from current and future experiments helps to narrow the range of viable dark-matter candidates.

Light dark matter in the mass range from 1 MeV to 1 GeV has attracted considerable attention following a series of sufficiently stringent constraints on heavy dark-matter scenarios [3–5]. If dark matter in this regime were once in

thermal equilibrium with the Standard Model plasma in the early Universe, reproducing the observed relic abundance typically requires an efficient depletion mechanism [6–8]. Such a mechanism is often provided by portal interactions mediated by new fields of different spin. In such scenarios, the dark sector can communicate with the Standard Model through mediators of spin-0 (e.g., light hidden Higgs bosons) [9–17], spin 1 (e.g., sub-GeV dark photons) [18–29], and spin-2 (e.g., massive dark gravitons) [30–36].

The inelastic dark-matter paradigm was introduced in the context direct-detection experiments where transitions between nearly degenerate states can qualitatively reshape the kinematics of scattering [37]. Originally proposed to explain the anomalous modulation signal reported by the DAMA Collaboration [38], inelastic dark-matter paradigm has since developed into a theoretically compelling scenario for sub-GeV thermal dark matter.

The realization of inelastic dark matter was developed in terms of inelastic fermions, which consist of two states with a mass splitting and dominant off-diagonal interactions [39]. The key feature phenomenology arises from off-diagonal interactions between the dark-matter ground state, χ_1 with mass m_{χ_1} , and an excited state, χ_2 with mass $m_{\chi_2} \gtrsim m_{\chi_1}$. For sufficiently small mass splittings, Boltzmann suppression of the heavier state sets in only after freeze-out and therefore has little impact on the total dark-matter relic abundance [40,41]. However, for sufficiently large mass splittings, the abundance of the heavier dark-matter state begins to be suppressed already by the time of freeze-out [42,43]. The relevant parameter space can be excluded by the accelerator based experiments [44–48].

Direct-detection experiments aim to measure rare energy depositions produced by interactions between Galactic-halo dark-matter particles and detector target material [49].

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In particular, dark-matter particles can scatter off nuclei or electrons in the target material of terrestrial detectors. The nuclear-recoil channel is the canonical probe for weak-scale dark matter; however, its sensitivity deteriorates rapidly for sub-GeV dark matter, since the recoil energies transferred to nuclei typically fall below experimental threshold [50]. Therefore, scattering on atomic electrons provides a well-motivated probe of light dark matter [51–54].

In liquid-xenon detectors, this motivates low-threshold analyses based on ionization-only electronic-recoil data, which have become probes of sub-GeV dark matter. However, in the light dark-matter mass range, the bound-state nature of the initial electron must be taken into account [55]. A realistic treatment of dark-matter-electron scattering requires the shell structure and detector-specific ionization response to be taken into account explicitly [56].

In this work, we derive phenomenological direct-detection constraints on the parameter space of thermal, light inelastic fermion dark-matter models. We consider a leptophilic scalar portal, for which direct scattering off bound electrons provides the dominant channel, whereas the Migdal contribution does not occur at the tree level. We use publicly available electron-ionization data from the XENON1T, PandaX-4T, and LUX-ZEPLIN (LZ) experiments.

For the inelastic scattering process $\chi_i + e^- \rightarrow \chi_f + e^-$, we use a following definition of the mass splitting: $\delta = m_f - m_i$. This definition allows the mass splitting to be either positive or negative. Additionally, we introduce the relative mass splitting in the following form $\Delta = \delta/m_{\chi_1}$, which is normalized to the lighter dark-matter mass.

We note that the strengthening of the signal in the down-scattering regime, $\chi_2 T \rightarrow \chi_1 T$, is a characteristic consequence of inelastic dark-matter kinematics, that can similarly arise for mediators of different spin [57]. In particular, Ref. [58] considers two nearly degenerate scalar dark-matter states coupled through a vector mediator. The analysis above was specifically focused on the transition from the heavier state to the lighter one for a fixed absolute mass splitting chosen to account for the XENON1T electron-recoil excess [59].

Next, direct detection of light exothermic inelastic dark matter (iDM) was investigated in Ref. [60] through scattering off atomic electrons and ionization induced by the Migdal effect. In this case of pseudo-Dirac dark matter coupled through a vector mediator, A' , the relative contributions of these channels depend on the dark-matter mass and mass splitting, as well as on the relative strengths of the dark-matter couplings to leptons and quarks. The authors of Ref. [60] map their constraints on the thermal target of the iDM mediated by hidden vector particle. Reference [61] extends the analysis of exothermic inelastic scalar dark matter coupled through a dark photon to the PandaX-4T data, accounting for both direct electron scattering and ionization induced by the Migdal effect. However,

Refs. [61] derive constraints on the coupling without focusing on the thermal target.

In Refs. [58,60,61], the main focus is the signal amplification in the XENON1T and PandaX-4T experiments for the down-scattering regime, that implies the vector mediated iDM. In the present paper, we discuss simplified benchmark scenario with leptophilic scalar mediator implying Majorana iDM. Since the thermal target curves differ between vector and scalar mediator scenarios, the resulting constraints on the parameter spaces of inelastic dark matter are distinct. In addition, compared with Ref. [44], which focused on fixed-target searches for $\Delta \simeq \mathcal{O}(0.1)$, the present analysis studies the complementary regime $|\Delta| < \mathcal{O}(10^{-5})$ and liquid-xenon direct-detection experiments. Our analysis employs modern relativistic ionization form factors [56] and the LUX-ZEPLIN (LZ) data.

A more systematic description of light pseudo-Dirac dark matter coupled through a vector mediator, including its thermal history, was developed in Ref. [41]. In particular, for a vector mediator, light thermal pseudo-Dirac dark matter with a mass splitting $|\Delta| < \mathcal{O}(10^{-6})$ is constrained by cosmic microwave background (CMB) observations [41]. The phenomenology of thermal pseudo-Dirac dark matter with a vector mediator was further developed in Ref. [43] to sufficiently large mass splittings, including interactions associated with baryons and leptons. Reference [62] employs a similar vector-mediator model but focuses primarily on the resonant regime, $m_{A'} \simeq 2m_\chi$. In the resonant regime, the observed relic abundance can be reproduced with a small coupling, thereby avoiding the CMB constraints.

The studies [41,43,62] focus on the parameter space of thermal iDM with a vector mediator. Our present analysis extends the existing concept of kinematic enhancement of the process $\chi_2 e \rightarrow \chi_1 e$ to the case of a leptophilic scalar mediator.

As we discussed above, the scalar [63] and vector [41] portals give rise to different thermal target curves and, as a result, different cosmological constraints. In particular, the p-wave suppression for the scalar-mediator Majorana iDM does not impose constraints from CMB. Therefore, for a scalar mediator, the region of parameter space with small mass splittings, $|\Delta| < \mathcal{O}(10^{-6})$, remains relevant for searches in direct-detection experiments. Moreover, the couplings of the scalar mediator to the heavier lepton generations are additionally enhanced compared to its electron coupling. Specifically, for $m_\mu < m_{\chi_1}$, the model with a scalar mediator and Majorana iDM is unconstrained and is a subject for direct-detection experiments for mass splitting $|\Delta| < \mathcal{O}(10^{-5})$.

This paper is organized as follows. In Sec. II, we discuss the simplified benchmark model, the parameter space for inelastic DM with a scalar mediator, and the direct-detection experiments. In Sec. III, we summarize the general expressions which are used to estimate direct-detection constraints

for exothermic and endothermic reactions. In Sec. IV, we discuss the resulting direct-detection constraints on light scalar-mediated inelastic fermion dark matter. Finally, conclusions are drawn in Sec. V.

II. BENCHMARK SCENARIOS AND EXPERIMENTS

In this section, we describe the dark-matter models employed in our analysis and the parameter region under consideration. We also briefly review the main characteristics of the direct-detection experiments considered in this work.

A. Simplified benchmark scenario

The dimension-5 effective operator that couples a leptophilic scalar dark-sector mediator ϕ to the standard model (SM) charged-lepton sector reads as [63]

$$\mathcal{L}_{\text{eff}}^{\phi} \supset \frac{1}{2}(\partial_{\mu}\phi)^2 - \frac{1}{2}m_{\phi}^2\phi^2 - \sum_{l=e,\mu,\tau} c_{ll}^{\phi} \bar{l} l \phi, \quad (1)$$

where we use the flavor-dependent ratio for the coupling constants [64]:

$$c_{ee}^{\phi} : c_{\mu\mu}^{\phi} : c_{\tau\tau}^{\phi} = m_e : m_{\mu} : m_{\tau}. \quad (2)$$

The dark-matter sector consists of two fermion states χ_1 and χ_2 , described by the Lagrangian [37,65],

$$\mathcal{L}_{\text{kin.term.}}^{\text{DM}} \supset \sum_{i=1,2} \left[\frac{1}{2} \bar{\chi}_i i \gamma^{\mu} \partial_{\mu} \chi_i - \frac{1}{2} m_{\chi_i} \bar{\chi}_i \chi_i \right], \quad (3)$$

where m_{χ_i} denotes the physical fermion masses, and we take $m_{\chi_1} \lesssim m_{\chi_2}$ such that χ_1 is the lightest stable DM candidate.

We focus on the effective benchmark Lagrangian involving a leptophilic scalar dark-sector mediator that couples to a pair of fermions as [17,66]

$$\mathcal{L}_{\text{eff}}^{(+)\text{DM}} \supset \text{Re}(\lambda_{\chi_1\chi_2}^{\phi}) \bar{\chi}_1 \chi_2 \phi. \quad (4)$$

In our simplified scenario, only off-diagonal terms contribute to the effective interaction [44,67,68], and we use the standard notation, $\alpha_{\text{IDM}} = (\text{Re}[\lambda_{\chi_1\chi_2}^{\phi}])^2/(4\pi)$ for the dark sector fine structure constant.

B. Relic abundance of inelastic dark matter

In this subsection, we discuss the freeze-out mechanism, assuming a kinetic and chemical equilibrium between DM and the SM thermal bath in the early Universe [69,70]. As the Universe expands, dark matter departs from thermal equilibrium and becomes depleted from the thermal bath.

Furthermore, the depletion mechanism via portal interactions leads to observed density of dark matter [43,63,71].

The current value of the cold DM relic abundance obtained from the Planck 2018 combined analysis is [1,72]

$$\Omega_c h^2 = 0.1200 \pm 0.0012.$$

The relic density in the case of the coannihilation channel $\chi_1\chi_2 \rightarrow \ell^+\ell^-$ is estimated to be [73,74]

$$\Omega_c h^2 \propto \left(\int_{x_f}^{\infty} \frac{\langle \sigma_{\text{eff}} v \rangle}{x^2} dx \right)^{-1}, \quad (5)$$

where $\langle \sigma_{\text{eff}} v \rangle$ is a effective thermally averaged coannihilation cross section and $x = m_{\chi_1}/T$ is a ratio of DM mass to the temperature of the SM plasma. A more detailed discussion and explicit expression of the thermally averaged coannihilation cross section for the considered benchmarks can be found in the Ref. [44].

It should be emphasized that we focus on relative mass splittings $|\Delta| \ll 1/20$. In this regime, Boltzmann suppression of the heavier dark-matter state occurs after the chemical decoupling of the visible and dark sectors, at $T \simeq m_{\chi} |\Delta|$ [41]. Thus, for the relative mass splittings considered here, the thermal relic abundance curves for inelastic dark matter coincide with those of the elastic case.

The additional energy injection of DM annihilation in the early Universe can leave an imprint on the measured anisotropy and polarization spectra of CMB [75]. However, straightforward constraints from CMB for the considered dark-matter model with a scalar mediator can be weakened due to the p-wave annihilation. Thus, inelastic fermion dark matter with a scalar mediator can have an unconstrained region of parameter space similar to the case with a vector mediator [76].

After the process $\chi_1\chi_2 \rightarrow \ell^+\ell^-$ becomes suppressed, inelastic fermion dark matter chemically decouples from the visible sector at a temperature $T_f \simeq m_{\chi_1}/20$, which fixes the total dark-matter abundance. However, in the presence of a scalar leptophilic mediator, the processes $\chi_2\ell \rightarrow \chi_1\ell$ and $\chi_2\chi_2 \rightarrow \chi_1\chi_1$ can modify the fraction of the heavier dark-matter state. In particular, the kinetic decoupling temperature can be estimated at the order-of-magnitude level as $T_{\text{kd}}^{\chi} \simeq \mathcal{O}(m_e)$, due to the start of electron depletion in the early plasma [41]. For small mass splitting, the chemical decoupling between the heavy and light dark-matter states, occurring at temperature T_{ch}^{χ} , is controlled by the process $\chi_2\chi_2 \rightarrow \chi_1\chi_1$ [43].

It is also worth emphasizing that coannihilation is suppressed relative to the process $\chi_2\chi_2 \rightarrow \chi_1\chi_1$ by the coupling between the visible and hidden sectors, particularly for the benchmark value $\alpha_{\text{IDM}} = 0.5$. Moreover, one can estimate that the channel $\chi_1\chi_1 \rightarrow \chi_2\chi_2$ is effectively open for hidden-sector temperatures satisfying $T_{\chi} \gtrsim \mathcal{O}(4|\Delta|m_{\chi_1}/3)$, where

we have used $\sigma_{\chi_1\chi_1 \rightarrow \chi_2\chi_2} \sim \sqrt{v^2 - 8|\Delta|}$ and $\langle v^2 \rangle = 6T_\chi/m_{\chi_1}$. Nevertheless, this process can also proceed at lower temperatures through the high-velocity tail of the dark-matter distribution. We therefore adopt the standard assumption of Ref. [40] that $T_{\text{ch}}^\chi < T_f$. The fraction of the heavier state is given by $f_2 = e^{-|\Delta|m_{\chi_1}/T_\chi}/(1 + e^{-|\Delta|m_{\chi_1}/T_\chi})$, as long as the process $\chi_1\chi_1 \rightarrow \chi_2\chi_2$ remains efficient [41].

For $T \lesssim T_{\text{ch}}^\chi$, the ratio of number densities of the dark-matter particles in excited and grounded states can be treated as constant for the sufficiently small mass splittings [62]. In particular, for small mass splittings and a vector mediator, the number densities of the excited and ground states remain comparable at late times [43]. It should also be noted that, for small mass splitting $|\delta| < 2m_e$, decays of the heavy dark-matter state into visible-sector particles are suppressed [41,76]. In addition, this parameter region is expected to evade searches for the semivisible mode in accelerator based experiments, since the decay channel $\chi_2 \rightarrow \chi_1 e^+ e^-$ is suppressed.

One can assume that the Boltzmann suppression at temperature $T_\delta \simeq \delta \lesssim \mathcal{O}(10^{-5})m_{\chi_1}$ occurs after the internal dark-sector chemical decoupling, $T_\delta \lesssim T_{\text{ch}}^\chi$, for dark-matter masses in the range from 1 MeV to 1 GeV. Under these assumptions, the fraction of the heavy state is $f_{\chi_2} \simeq 1/2$. Otherwise, Boltzmann suppression can reduce the abundance of the heavier state to a negligibly small level.

We left for future work a more precise calculation of the relevant temperatures at which the two dark-matter states freeze out and evolve in number density—specifically for the inelastic fermion DM model with a scalar leptophilic mediator. The current analysis uses simplified assumptions.

As an approximate approach, one can treat the dark-matter state fractions f_i (the relative abundances of the two states) as free parameters. These fractions effectively rescale the constraints derived from direct-detection experiments: the electron-scattering cross section σ_e inferred from the data is scaled down by a factor of f_i to account for the fact that only a fraction of the total dark matter participates in the inelastic scattering process.

In this work, we focus on the parameter space of light thermal dark matter with scalar mediator defined by

$$\frac{m_{\chi_1}}{m_\phi} = \frac{1}{3}, \quad |\Delta| \lesssim 10^{-5}, \quad m_e \lesssim m_{\chi_1}, \quad \alpha_{\text{IDM}} = 0.5. \quad (6)$$

For this region, we introduce two scenarios:

- (i) *Endothermic (up-scattering) scenario*. The abundance of the heavier dark-matter state is negligibly small, implying $f_{\chi_1} \simeq 1$. Consequently, down scattering in direct-detection experiments is suppressed by the small fraction $f_{\chi_2} \simeq 0$. As a result, up-scattering process provides the dominant channel for the constraints from direct-detection experiments.
- (ii) *Exothermic (down-scattering) scenario*. The abundances of the heavier and lighter dark-matter states

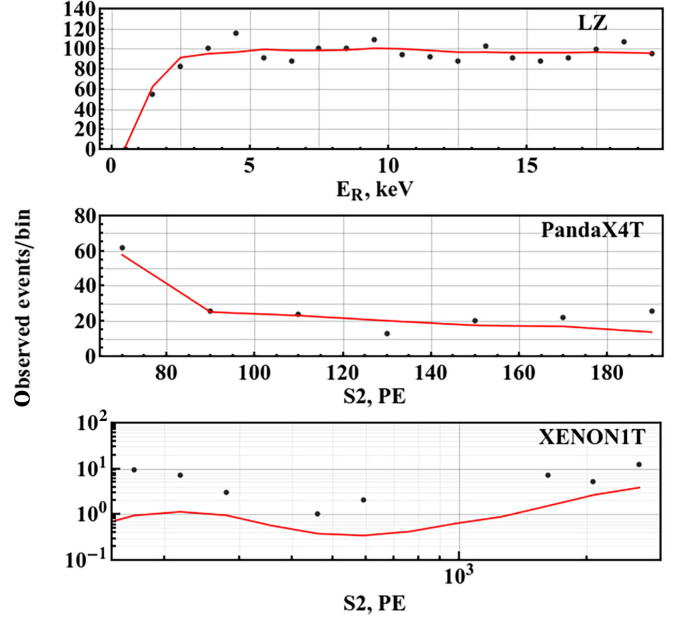


FIG. 1. Event and background public data for XENON1T [77] PandaX-4T [78], and LZ [79] experiments. The red line corresponds to case of simulated background, and black points are observed data after imposing all cuts in experiments.

are comparable, $f_{\chi_2} \simeq f_{\chi_1} \simeq 1/2$. In this case, down scattering provides the dominant channel for constraints from direct-detection experiments.

The results obtained within the proposed scenarios allow us to estimate the sensitivity of direct-detection experiments to each of the scattering channels.

C. Experiments

Direct-detection experiments search for dark matter via rare interactions in a terrestrial target, where such interactions can lead to small electron reconstructed energy. Strong background rejection is employed to isolate candidate recoil signals. The public event and background data used in our analysis for XENON1T, PandaX-4T, and LZ are shown in Fig. 1.

1. XENON1T

The XENON1T experiment is an underground direct-detection search for dark matter operated from 2015 to the end of 2018 with a integrated time of 258.2 days at the Laboratori Nazionali del Gran Sasso in Italy. The core of XENON1T is a dual-phase time projection chamber (TPC) containing of two tonnes of liquid xenon, bounded by a grounded electrode at the top and a cathode at the bottom. Energy deposited by charged recoils can generate both prompt scintillation (S1) and ionization electron signal (S2). To achieve an ultralow-background environment, the active detector is shielded by multiple layers, including an approximately 3600 m water-equivalent rock overburden,

an active water Cherenkov muon veto, and an additional 1.2 ton of liquid xenon (LXe) surrounding the TPC. The dominant background contributions arise from β -decays of radioactive impurities, mainly Pb^{214} , in the active detector volume; from neutrino backgrounds, primarily coherent elastic neutrino-nucleus scattering (CE ν NS); and from β -decays originating on the cathode. In particular, the nearly flat background from β -decays of impurities is reduced through the use of cleaner materials. Also, cathode-related backgrounds are suppressed by applying a selection on the S2 signal width. We use the data from the Ref. [77] which reports an S2-only data based on ionization electrons. The measured S2-signal energies for electronic recoils span the range from 0.186 keV (150 PE), with lower-energy events excluded because of poorly controlled backgrounds, up to $\simeq 3$ keV (3000 PE). Also, the full detector efficiency is already incorporated into the response matrix. It should also be emphasized that the publicly available XENON1T S2-only data do not include the analysis-specific cathode background.

2. PandaX-4T

The PandaX-4T experiment is an underground direct-detection search for dark matter located in B2 hall of the China Jinping Underground Laboratory (CJPL-II) in Sichuan, China. Its commissioning run started on November 28, 2020, and ended on April 16, 2021, comprising 95.0 calendar days of stable data taking. The detector is a multiton dual-phase xenon TPC with a sensitive target of 3.7 ton of LXe contained within a double-vessel cryostat holding 5.6 ton of LXe in total. The TPC is a cylindrical LXe volume with a cathode at the bottom and gate and anode grids near the surface, providing drift and extraction fields. Energy deposited by particle interactions generates both prompt scintillation photons (S1) in the LXe and a delayed ionization electron signal (S2). Both signals are collected by photomultiplier tube arrays at the top and bottom of the TPC. To achieve a low-background environment, PandaX-4T benefits from an overburden of ~ 2.4 km rock (corresponding to ~ 6720 m water equivalent) and is surrounded by an ultrapure-water shield in a stainless-steel tank with a diameter of 10 m and a height of 13 m. The dominant background contribution in the DM-electron ionization-only channel arises from β -decays of the internal radioactive contaminants. We use the public PandaX-4T data based on ionization electrons from Ref. [78]. The S2-signal energy range extends from 0.07 keV (60 PE) to 0.23 keV (200 PE) for electronic recoils. The lower and upper boundaries are determined by high background rate at very low S2 and by sensitivity to DM search, respectively.

3. LZ

The LZ experiment is a direct-detection search for dark matter conducted at the Sanford Underground Research

Facility in Lead, South Dakota, at a depth of 4850 ft (4300 m water equivalent). For the 2024 weakly interacting massive particle search, LZ used data collected between 2021 and 2024, corresponding to an exposure 4.2 ton-years. To suppress backgrounds arising from the radioactivity of detector components, the TPC is enclosed within a system consisting of a 2 ton LXe gamma-tagging detector and an outer detector containing 17.3 tons of gadolinium-loaded liquid scintillator, optimized for neutron detection. These active components are further surrounded by 238 tons of ultrapure water, providing additional passive shielding and ensuring a low-background environment. Particle interactions in the TPC produce prompt scintillation light (S1) and ionization electrons, which generate secondary electroluminescence (S2). The dominant background contribution arises from the β -decay of Pb^{214} , for which a radon-tagging technique is employed. We use the publicly available LZ data based on ionization electrons from Ref. [79]. The electron-recoil signal energy range extends from 1 to 20 keV for electronic recoils.

III. SIGNATURES OF THE DIRECT DETECTION

In this section, we summarize the general expressions that can be used for the estimation of constraints on dark-matter models from direct-detection experiments. As mentioned above, we focus on small relative mass splittings, which imply $m_{\chi_1} \simeq m_{\chi_2} \equiv m_\chi$. We consider both exothermic and endothermic processes of dark-matter scattering on electrons.

A. Kinematics

Let us consider the inelastic scattering of dark matter off an atomic electron,

$$\chi_i(p_i) + e^-(p_2) \rightarrow \chi_f(p_f) + e^-(p_4),$$

where the momentum transfer is defined as $q^\mu \equiv p_i^\mu - p_f^\mu$ and the mass splitting is $\delta = m_f - m_i$. This process is endothermic ($\delta > 0$) in the case of up scattering, while down scattering leads to an exothermic process with $\delta < 0$. The momentum transfer depends on the dark-matter velocity v_χ and deposited energy E_d as [58]

$$q_\pm(v) = \left| m_\chi v_\chi \pm \sqrt{m_\chi^2 v_\chi^2 - 2m_\chi(E_d + \delta)} \right|, \quad (7)$$

which implies the following condition on the dark-matter velocity for exothermic processes with $|\delta| < E_d$ and endothermic processes as

$$v_\chi^2 > 2(E_d + \delta)/m_\chi. \quad (8)$$

In particular, for the up-scattering process, one finds the estimate $\Delta < (v_{\chi_1})^2/2$. Assuming that the dark-matter population in the Solar System is gravitationally bound

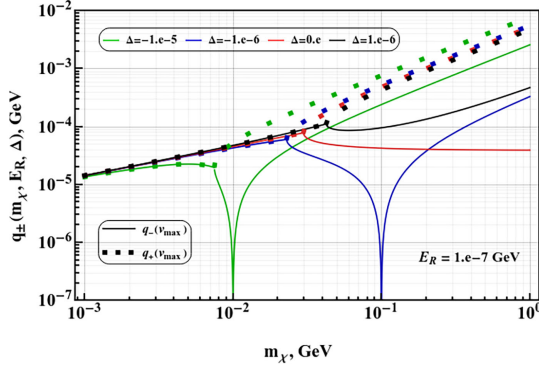


FIG. 2. Transferred momentum limits (11) as functions of the DM mass with the fixed deposited energy. Each color corresponds to a different value of the relative mass splitting. The dotted and solid lines denote the upper and lower limits, respectively.

to the Milky Way [80,81], the dark-matter velocity is bounded by

$$v_{\chi}^{\max} = v_{\text{esc}} + v_E \simeq 2.58 \times 10^{-3}. \quad (9)$$

Thus, in the case of up-scattering process, signal events in direct-detection experiments are only possible for a relative mass splitting $\Delta \lesssim \mathcal{O}(10^{-6})$.

The minimum dark-matter velocity as function of momentum transfer is

$$v_{\min}(q) \simeq \left| \frac{E_d + \delta}{q} + \frac{q}{2m_{\chi}} \right|. \quad (10)$$

Therefore, the lower limit of the dark-matter velocity reaches zero only in the exothermic case with $|\delta| > E_d$. Imposing the dark-matter velocity condition $v_{\min}(q) < v_{\chi}^{\max}$ leads to the following constraints on momentum transfer:

$$q_{-}(v_{\chi}^{\max}) < q < q_{+}(v_{\chi}^{\max}). \quad (11)$$

The corresponding momentum ranges for different parameter choices are shown in Fig. 2. It should also be noted that, for the considered processes and the mass splittings satisfying $E_d < m_{\chi}|\Delta|$, the limits on the momentum transfer tend to $m_{\chi}(v_{\chi}^{\max} \pm \sqrt{(v_{\chi}^{\max})^2 - 2\Delta})$. In the region $E_d > m_{\chi}(v_{\chi}^{\max})^2/2, m_{\chi}|\Delta|$, kinematic suppression occurs, such that $q_{+} \simeq q_{-} \simeq \sqrt{2m_{\chi}E_d}$. In addition, in the regime,

$$m_{\chi} \simeq E_d/|\Delta|, \quad (12)$$

the momentum limits become 0 and $2m_{\chi}v_{\chi}^{\max}$, which leads to an enhancement in this mass region compared with the elastic case. Also, increasing the deposited energy shifts the enhancement region toward larger masses.

In the light-dark-matter mass regime, the typical deposited energies in direct-detection experiments are of order

$$E_d \simeq \mathcal{O}(10^{-2}) - \mathcal{O}(1) \text{ keV},$$

which corresponds to momentum transfers of order $q \lesssim \mathcal{O}(10^{-3}) - \mathcal{O}(1) \text{ MeV}$. Therefore, when light dark matter scatters off a target material, the deposited energy can be sufficient to induce inelastic atomic processes, and one must account for the bound-state nature of the initial electron [51].

The impact of the dark-matter mass splitting on the constraints from direct-detection experiments in the endothermic case becomes important when the splitting is comparable to the typical deposited energy, $|\delta| \simeq \mathcal{O}(1) \text{ keV}$. In this case, the up scattering of inelastic dark matter leads to a smaller deposited energy than in the elastic case, resulting in weaker constraints.

B. Experimental reach

In the general framework of nonrelativistic effective field theory [82], several atomic response functions must be taken into account for scattering off atomic electrons [83,84]. Let us consider first the effective spin-independent interaction $(\bar{e}e)(\bar{\chi}\chi)$ that reduces at leading nonrelativistic order to the operator $\mathcal{O}_1 = \mathbf{1}_{\chi}\mathbf{1}_e$. The matrix element for dark-matter scattering off a bound electron admits a simple factorization [83,84],

$$\mathcal{M}_{i \rightarrow f}^{\text{bound}}(\mathbf{q}) = \mathcal{M}_{\chi e}^{\text{free}}(q) \langle f | e^{i\mathbf{q} \cdot \hat{\mathbf{r}}} | i \rangle,$$

where $|i\rangle$ and $|f\rangle$ denote the initial and final electron states, respectively, and $\mathcal{M}_{\chi e}^{\text{free}}(q)$ is the matrix element for dark-matter scattering off a free electron. The transition amplitude $\langle f | e^{i\mathbf{q} \cdot \hat{\mathbf{r}}} | i \rangle$ encodes the structure of the target material, and its explicit form depends on the normalization convention adopted. Moreover, at momentum transfers of order $|\mathbf{q}|^2 \simeq (am_e)^2$, the bound-state nature of the initial electron becomes important. Following the standard approach, we introduce a reference cross section $\bar{\sigma}_e$ and the dark-matter form factor $F_{\text{DM}}^2(q)$ as follows [51]:

$$\bar{\sigma}_e \equiv \frac{\mu_{\chi e}^2}{16\pi m_{\chi}^2 m_e^2} |\mathcal{M}_{\chi e}^{\text{free}}(q)|^2_{|q|^2=(am_e)^2},$$

$$F_{\text{DM}}^2(q) \equiv \frac{|\mathcal{M}_{\chi e}^{\text{free}}(q)|^2}{|\mathcal{M}_{\chi e}^{\text{free}}(q)|^2_{|q|^2=(am_e)^2}}.$$

This representation factorizes the entire transferred-momentum dependence of the free-electron matrix element into the dark-matter form factor.

The differential event rate for dark matter with fraction f_i is given by [56]

$$\frac{dR}{dE_d} = \frac{N_T \rho_\chi f_i d\langle\sigma v\rangle}{m_T m_\chi dE_d}, \quad (13)$$

where N_T is the number of target atoms of mass m_T , E_d is the energy deposited in the target, $\mathbf{v}$ is the dark-matter velocity in the Earth frame, and $\rho_{\text{DM}} \simeq 0.4 \text{ GeV/cm}^3$ is the local dark-matter density near the Earth [81,85]. The thermally averaged differential cross section for dark-matter scattering off a bound electron is [51,56]

$$\frac{d\langle\sigma v\rangle}{dE_d} = \frac{\bar{\sigma}_e}{2\mu_{\chi e}^2} \int_{q_-(v_{\chi 1}^{\text{max}})}^{q_+(v_{\chi 1}^{\text{max}})} q \frac{K(E_d, q)}{E_H} \cdot \eta(v_{\min}) |F_{\text{DM}}(q)|^2 dq, \quad (14)$$

where $E_H = m_e \alpha^2$ is the Hartree energy, $K(E_d, q)$ is the corresponding total atomic ionization factor, and $\eta(v_{\min}) = \langle \frac{1}{v} \theta(v - v_{\min}) \rangle$ is the averaged inverse velocity of dark matter over the distribution function $f_{\text{TMB}}(\vec{v} + \vec{v}_E)$. In the case of the Standard Halo Model, one can use the local truncated Maxwell-Boltzmann distribution as [86]

$$f_{\text{TMB}}(\mathbf{v} + \mathbf{v}_E) = \begin{cases} \frac{1}{N_{\text{esc}}} e^{(-|\mathbf{v} + \mathbf{v}_E|^2/v_0^2)}, & |\mathbf{v} + \mathbf{v}_E| < v_{\text{esc}} \\ 0, & |\mathbf{v} + \mathbf{v}_E| > v_{\text{esc}}, \end{cases}$$

$$N_{\text{esc}} = \pi^{3/2} v_0^3 \left(\text{erf}(v_{\text{esc}}/v_0) - \frac{2}{\sqrt{\pi}} \frac{v_{\text{esc}}}{v_0} e^{-v_{\text{esc}}^2/v_0^2} \right).$$

where $v_0 \simeq 220 \text{ km/s}$ is the characteristic velocity, $v_{\text{esc}} \simeq 544 \text{ km/s}$ is the escape velocity, and $\mathbf{v}_E$ is the velocity of the Earth in the Galaxy with $v_E \simeq 232 \text{ km/s}$. An explicit expression for $\eta(v_{\min})$ is provided in the Ref. [87] that is nonzero within the momentum-transfer range defined in (11). Tabulated values of the total atomic ionization factor for different target materials can be found in Ref. [56], where this quantity is computed using a relativistic Hartree-Fock method. We explicitly take into account the momentum-transfer limits in order to improve the robustness of the numerical calculations. The total ionization factor and the thermally averaged inverse velocity for representative deposited energies are shown in Fig. 3.

Since the relevant particle masses are larger than $\mathcal{O}(1) \text{ MeV}$, the momentum transfer is much smaller than the particle masses. Therefore, we work in the regime $t \ll m_e^2, m_\phi^2, m_{\chi_1}^2$. After summing over final and averaging over initial internal degrees of freedom, the squared matrix element becomes

$$\begin{aligned} |\overline{\mathcal{M}}|^2 &= \frac{4\pi\alpha_{\text{DM}}(c_{ee}^\phi)^2}{(t - m_\phi^2)^2} (4m_e^2 - t)((m_{\chi_1} + m_{\chi_2})^2 - t) \\ &\simeq \frac{4\pi\alpha_{\text{DM}}(c_{ee}^\phi)^2}{(t - m_\phi^2)^2} (4m_e^2)(m_{\chi_1} + m_{\chi_2})^2. \end{aligned} \quad (15)$$

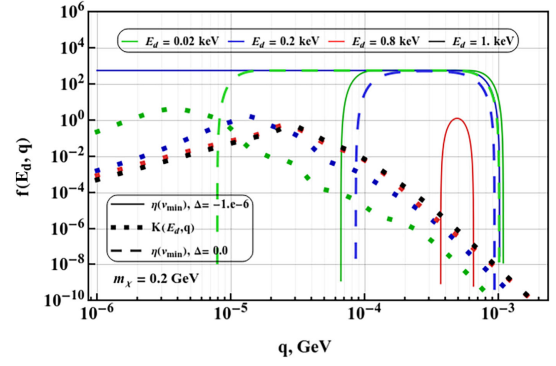


FIG. 3. Total ionization factor and thermally averaged inverse velocity as functions of the transferred momentum for different deposited energies. Different colors correspond to different deposited energies. The quantity $\eta(v_{\min})$ is shown for relative mass splittings $\Delta = 0.0$ and $\Delta = -10^{-6}$ by dashed and solid lines, respectively. The total ionization factor is shown by the dotted line.

The dark-matter form factor and the effective cross section read, respectively,

$$F_{\text{DM}}^2(t) = \frac{((\alpha m_e)^2 + m_\phi^2)^2}{(|\mathbf{q}|^2 + m_\phi^2)^2}, \quad (16)$$

$$\bar{\sigma}_e = \frac{4\pi\alpha_{\text{DM}}(c_{ee}^\phi)^2}{\pi} \frac{\mu_{\chi_1}^2}{(m_\phi^2 + (\alpha m_e)^2)^2}. \quad (17)$$

In the parameter region of interest, $m_\chi/m_\phi = 1/3$ and $m_e < m_\chi$, the dark-matter form factor can be set to $F_{\text{DM}}^2(t) \simeq 1$. However, in deriving the constraints, we explicitly included the form factor (16) in the calculation, and we also checked that $F_{\text{DM}}^2(t) \simeq 1$ is a reasonable approximation.

The Migdal effect [88] induced by DM-nucleus scattering can provide an additional contribution to low-energy electron signals [60,89]. However, in the present work, we focus on a leptophilic scalar mediator, for which the dominant interaction is the tree-level coupling to electrons. For this reason, we restrict our direct-detection analysis to the electron-scattering channel and leave a dedicated study of Migdal contributions for future work, specifically for hadron-specific mediator scenario.

In our analysis of direct-detection data, we constrain the signal strength using a one-dimensional profile-likelihood procedure. The expected number of events in each bin is given by

$$\mu_i(\bar{\sigma}_e) = n_{\text{Bkg},i} + \bar{\sigma}_e / (1 \text{ cm}^2) \cdot n_{\text{thr},i}(\bar{\sigma}_e = 1 \text{ cm}^2).$$

For the full set of bins, we construct the Poisson log likelihood as

$$\log L(\bar{\sigma}_e) = \sum_i [n_{\text{Obs},i} \ln(\mu_i(\bar{\sigma}_e)) - \mu_i(\bar{\sigma}_e) - \ln(n_{\text{Obs},i}!)].$$

We obtain the maximum-likelihood estimate of the upper limit by numerically maximization. The profile-likelihood-ratio test statistic is [90]

$$q(\bar{\sigma}_e) = -2[\log L(\bar{\sigma}_e) - \log L(\hat{\bar{\sigma}}_e)]. \quad (18)$$

The upper limit on the parameter is obtained by solving $q(\bar{\sigma}_e) = 1.642$, which corresponds to a one-sided 90% confidence level in the asymptotic Wilks approximation for a single parameter.

One can calculate number of theoretical signal events by the expression,

$$N_{\text{sign}} = \epsilon \cdot \sum_j \sum_i R_{ij} \frac{N_T \rho_\chi}{m_T m_\chi} \frac{d\langle\sigma v\rangle}{dE_d}(E_{di}) dE_{di}, \quad (19)$$

where $\sum_j$ and $\sum_i$ are sum over event bins and energy discretizations, respectively, and R_{ij} is the response matrix.

For XENON1T [77] and PandaX-4T [78], we employ the published response matrix, which already incorporates the full detector efficiency. LZ [79] publicly provides the total energy-dependent efficiency and the reconstructed-energy distributions, but without a response matrix that accounts for transitions between different energy bins. We therefore adopt the diagonal-response approximation $R_{ij}^{\text{LZ}} = \varepsilon_{\text{LZ}} \delta_{ij}$. Accordingly, the LZ result should be understood as a public-data recast within this approximation.

For an order-of-magnitude estimate of the constraints from direct-detection experiments, one can use the Bayesian approach and derive limits based on [91–93]

$$\Gamma(n_{\text{obs}} + 1, s_{\text{upper}} + N_{\text{bkg}}) = \alpha \Gamma(n_{\text{obs}} + 1, N_{\text{bkg}}). \quad (20)$$

This unbinned treatment, used to derive the signal bound, yields a conservative estimate of the cross section limit (shown by dashed lines in Figs. 4 and 5) and therefore provides weaker constraints.

We use the background-only best-fit spectrum, $n_{\text{Bkg},i}$, published by each collaboration as the background model in our analysis. PandaX-4T and LZ report relative total background uncertainties of about $\mathcal{O}(10)\%$ [78] and $\mathcal{O}(3)\%$ [79], respectively. For XENON1T, a comparable fractional uncertainty is difficult to estimate from the public data owing to the analysis-specific cathode background [77].

In general, background uncertainties broaden the range of signal strengths compatible with the data and therefore weaken the expected constraints [94]. For instance, using a background-only dataset at $m_{\text{DM}} = 4.5$ GeV, the DarkSide-50 Collaboration demonstrated that uncertainties in the theoretical background spectra and normalizations weakened the expected limit by 15%–31% [95]. By

analogy, we expect background uncertainties to affect the limits for the experiments considered here at a similar order of magnitude.

For $\Delta = 0$, our XENON1T limits [77] agree with the corresponding collaboration results within a factor of $\mathcal{O}(1)$, while the elastic bounds for PandaX-4T agree with those of Ref. [78] to within $\mathcal{O}(10\%)$. A full treatment of background uncertainties requires the experiment-level likelihood and is beyond the scope of this work.

IV. RESULTS AND DISCUSSION

In this section, we discuss the direct-detection constraints on light inelastic DM [see Eq. (4)] with a scalar leptophilic mediator, within the parameter space defined in Eq. (6). Specifically, the PandaX-4T, XENON1T, and LZ experiments considered here are sensitive to deposited energies in the signal region of order $\mathcal{O}(10^{-8})$, $\mathcal{O}(10^{-7})$, and $\mathcal{O}(10^{-6})$ GeV, respectively. We employ both binned Eq. (18) and unbinned Eq. (20) approaches to estimate the direct-detection constraints. It is also important to note that the limit obtained using the unbinned approach (20) allows one to derive direct-detection constraints up to an additional $\mathcal{O}(1)$ factor compared to the likelihood-based estimate, for the direct-detection experiments considered here. The corresponding direct-detection constraints for the endothermic and exothermic scenarios are shown in Figs. 5 and 4, respectively.

To map the constraints from the BABAR and NA62 experiments into parameter space of interest, we employed the reference cross section (17) where the corresponding dependence $c_{ee}^\phi(m_\phi)$ was taken from the Refs. [96] and [97,98], respectively. Similarly, we use the constraints on scalar-mediator radiation from fixed-target experiments [44]. We also compare our computed results in the limit of zero mass splitting with known results of XENON1T and PandaX-4T experiments, and the resulting curves agree at the level of an $\mathcal{O}(1)$ factor. We also use the relic-density curves obtained in our previous work [44], where the mass splitting has no significant impact on the freeze-out mechanism in the considered parameter region (6). Note that both NA64e and NA62 experiments rule out the typical masses $m_{\chi_1} \lesssim 100$ MeV within the adopted thermal target benchmarks.

Endothermic scenario (up scattering, Fig. 5). In the case of an endothermic reaction, increasing the mass splitting leads to weaker constraints, up to the kinematically forbidden region $\Delta > 10^{-6}$. However, for the relative mass splittings $\Delta = 10^{-7}$, the constraints for the up-scattering process differ from the elastic case only at the level of an $\mathcal{O}(1)$ factor. Moreover, for smaller mass splittings, $\Delta \ll 10^{-7}$, the constraints on inelastic DM are close to the elastic case. Indeed, as the mass splitting between the states decreases, its impact on the minimum dark-matter

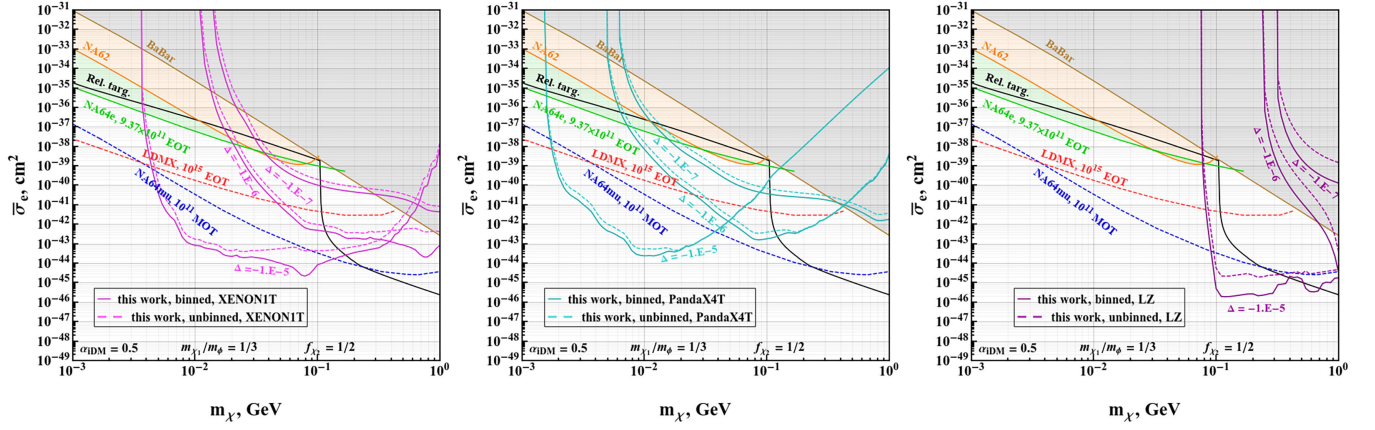
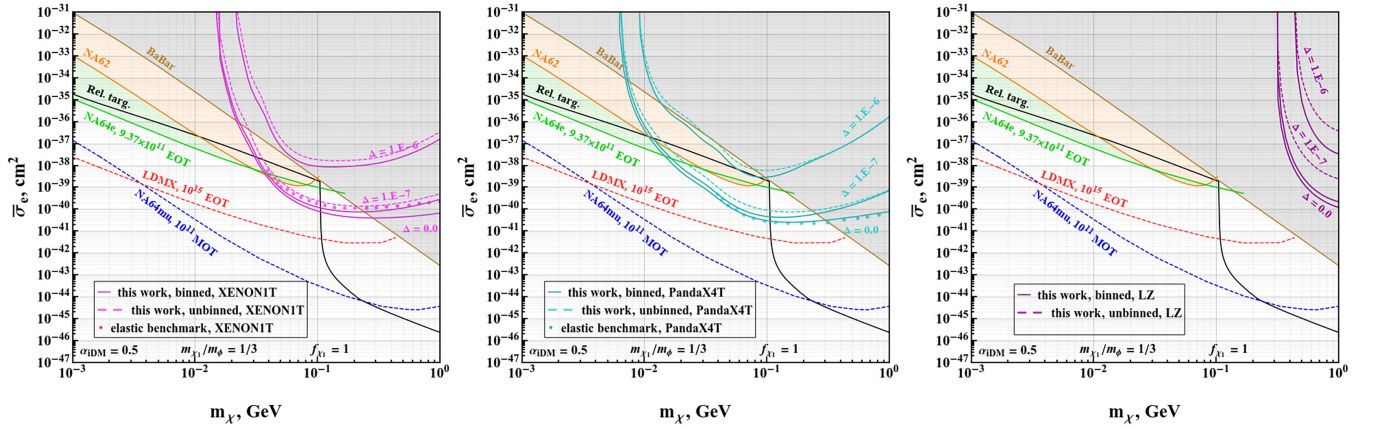
FIG. 4. The same as Fig. 5, but for down-scattering ($\Delta < 0$) scenario.

FIG. 5. Constraints of effective cross section as function of dark-matter mass in cases of benchmark (4) and up-scattering ($\Delta > 0$) scenario for different direct-detection experiments. Magenta, cyan, and purple colors are related by for XENON1T (left panel), PandaX-4T (center panel), and LZ (right panel) experiments where solid and dashed lines show the binned and unbinned calculations, respectively. Magenta and cyan dots indicate the benchmark constraints in elastic cases from Ref. [77] for XENON1T and from Ref. [78] for PandaX-4T, respectively.

velocity, Eq. (10), becomes smaller for an endothermic reaction. For the $\Delta \simeq 10^{-7}$, significant deviations from the elastic case $\Delta = 0$ appear at masses $m_{\chi_1} > 1$ GeV for LZ and a $m_{\chi_1} > 100$ MeV for XENON1T and PandaX-4T. Thus, sufficiently small mass splittings $\Delta \ll 10^{-7}$ do not lead to any additional weakening of the direct-detection constraints.

Exothermic scenario (down scattering, Fig. 4). In the case of exothermic scattering, the typical masses about $m_{\chi_1} \simeq E_d/|\Delta|$ provide better sensitivity (this corresponds to the lowest upper limit on $\bar{\sigma}_e$ shown in Fig. 4), for which the lower bound of dark-matter velocity is minimal; see Eq. (10). Given Δ , the corresponding sensitivity enhancement arises from the effect of the mass splitting on the kinematic quantities, as can be seen directly from Figs. 2 and 3 (see Sec. III A for a more detailed discussion). In particular, for a relative mass splitting of $|\Delta| = 10^{-5}$, the sensitivity peaks arise near DM masses of $\mathcal{O}(10^{-2})$,

$\mathcal{O}(10^{-1})$, and $\mathcal{O}(1)$ GeV for PandaX-4T, XENON1T, and LZ, respectively.

The constraints at relatively large masses, $m_{\chi_1} \gtrsim E_d/|\Delta|$, are weakened compared to the peak sensitivity region $m_{\chi_1} \simeq E_d/|\Delta|$ due to kinematic suppression, which implies shrinking of the integration limits Eq. (11). At relatively small masses, $m_{\chi_1} \lesssim E_d/|\Delta|$, the sensitivities shift toward larger cross sections. As the relative mass splitting increases, direct-detection experiments rule out sufficiently large portion of the parameter space, within small DM mass region, i.e., the larger $|\Delta|$, the better the sensitivity. However, sufficiently large $|\Delta| \gtrsim 10^{-5}$ are forbidden due to kinematics.

It is worth noticing that the bounds for exothermic scenario, shown in Fig. 4, are sensitive to the energy binned distribution; this leads to irregularities in the sensitivity curves. Remarkably, the extended limits from considered direct-detection experiments are comparable to

expected from the NA64 μ experiment with a statistics of 10^{11} muons on target, which is the most sensitive probe of this model.

V. CONCLUSION

In this work, we derived direct-detection constraints on thermal inelastic fermion dark matter coupled to a leptophilic scalar mediator using public data from XENON1T, PandaX-4T, and LZ. The mass splitting $|\Delta| \simeq \mathcal{O}(10^{-5}) - \mathcal{O}(10^{-6})$ for the down-scattering setup can lead to an enhancement of the sensitivity for the experimental facilities of interest. For a heavy-state fraction $f_{\chi_2} = 1/2$ and the exothermic scenario, the signal can be sufficiently enhanced in the characteristic mass regions, $m_{\chi_1} \simeq E_d/|\Delta|$, up to the level $\sigma_e \simeq \mathcal{O}(10^{-45}) \text{ cm}^2$. In this case, direct-detection experiments can additionally exclude the parameter space in the mass range, $100 \text{ MeV} < m_{\chi_1} < 500 \text{ MeV}$.

However, following Eq. (13), we note that a smaller fraction $f_{\chi_2} \ll 1/2$ would result in a weaker constraint.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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