

Reviving Z' portal dark matter with conversion mechanism

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In many new physics models with extended gauge symmetry, the new gauge boson Z' could mediate the interactions between the dark matter and standard model particles. For the conventional Z' portal dark matter, the collider and the direct detection constraints typically pose a significant challenge. To address this pressing issue, we present in this paper a new benchmark model based on the gauged $U(1)_{B-L}$ symmetry, which introduces a Dirac dark fermion $\tilde{\chi}_1$ and a heavier partner $\tilde{\chi}_2$ with zero and nonzero $U(1)_{B-L}$ charge, respectively. Including the mass term $\delta m \tilde{\chi}_1 \tilde{\chi}_2$ results in the dark fermions χ_1 and χ_2 in the mass eigenstate, where the lighter one χ_1 is regarded as the dark matter candidate. Various intriguing processes for the relic density arise with the compressed mass spectrum $m_{\chi_1} \simeq m_{\chi_2}$, such as the cospin scattering $\chi_2 f \rightarrow \chi_1 f$, the conversion $\chi_2 \chi_i \rightarrow \chi_1 \chi_j$, and the coannihilation $\chi_1 \chi_2 \rightarrow f \bar{f}$ processes. Suppressed by the small mixing angle θ between the dark fermions, the small effective gauge coupling of dark matter χ_1 to the gauge boson Z' is one distinct feature of this model, rendering phenomenology in many aspects more promising. In this paper, we investigate the production of dark matter through new mechanisms within the frameworks of resonance and secluded scenarios. The impacts of phenomenological constraints from colliders, dark matter, and cosmology are also taken into account. We report that the conversion mechanism is both favored by the resonance and secluded scenarios under current constraints.

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I. INTRODUCTION

The cosmological and astrophysical observations provide evidence for the feasibility of particle dark matter (DM) [1,2]. The prevailing hypothesis considers that the DM particle is in thermal equilibrium with the standard model (SM) bath during the early Universe, subsequently freezing out to yield the observed thermal relic density. This is known as the weakly interacting massive particle (WIMP) paradigm. Typically, the DM candidate is a singlet under the SM gauge symmetry, which then requires a portal to mediate the interactions between DM and SM [3,4]. Based on the spin of the mediator, it can be categorized into

the fermion portal [5–10] and the scalar portal [11–16], as well as the vector portal [17–21].

The gauged $U(1)_{B-L}$ symmetry is the simplest anomaly-free extension of the SM [22], in which three right-handed neutrinos N are introduced for anomaly cancellation. These right-handed neutrinos could produce tiny neutrino mass via the seesaw mechanism [23–26] and explain the baryon asymmetry via the leptogenesis mechanism [27–31]. Meanwhile, DM can be implemented in the $U(1)_{B-L}$ symmetry [32–36]. In this work, we focus on the phenomenon of DM. For simplicity, the influence of right-handed neutrinos on DM is disregarded by assuming m_N at the canonical seesaw scale $\mathcal{O}(10^{14})$ GeV.

In the traditional Z' portal Dirac DM scenario, the observed relic density of DM χ requires the gauge coupling $g' \sim \mathcal{O}(0.1)$ with new symmetry charge $Q_\chi = -1$ [37,38]. This results in the spin-independent DM-nucleon scattering cross section mediated by Z' being strictly constrained by current DM direct detection experiments [39–41]. Furthermore, the direct searches of Z' on various colliders favor exceedingly small g' for $m_{Z'}$ below the TeV scale [42–46]. To satisfy current constraints, one improvement approach is to treat Q_χ as a free parameter. In this way, the effective gauge coupling of DM $g_\chi = Q_\chi g'$ is expected to be much larger than g' , which can alleviate the pressure

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from direct detection in both Z' resonance [47] and secluded [48] schemes. An alternative pathway is the inelastic DM [49–51], in which the relic density is determined via coannihilation of the DM and the dark partner. With sufficiently large mass splitting, the inelastic DM-nucleon scattering is kinematically suppressed by the DM velocity [49].

Motivated by the above two methods, we investigate the Z' portal inelastic Dirac DM with the $U(1)_{B-L}$ symmetry [52]. This model includes one Z_2 odd Dirac fermion $\tilde{\chi}_1$ and its partner $\tilde{\chi}_2$, which carry zero $Q_{\tilde{\chi}_1} = 0$ and nonzero $Q_{\tilde{\chi}_2} \neq 0$ charge under the $U(1)_{B-L}$ symmetry, respectively. In this way, the interaction between the DM candidate χ_1 in the mass eigenstate and Z' occurs through the mixing θ between the two dark fermions, which arises from a mass term $\delta m \tilde{\chi}_1 \tilde{\chi}_2$. Considering that $g_\chi = Q_{\tilde{\chi}_2} g' \sim 1$, the significant annihilation rate of χ_2 pairs may lead to DM production governed by the traditional coannihilation mechanism [53] or through the novel cospattering or conversion mechanism [54,55]. The prominent characteristic of the latter is that the relic density of DM is determined by the inelastic conversions with its heavier dark partner [56–71]. This transformation process in our model primarily consists of cospattering $\chi_2 f \rightarrow \chi_1 f$ and conversion $\chi_2 \chi_i \rightarrow \chi_1 \chi_j$ processes with f the SM fermion and $\{i, j\} = \{1, 2\}$. This scenario imposes very lenient requirements on DM annihilation, i.e., allowing the relevant small values of mixing θ and gauge coupling g' . Therefore, it is possible to obtain the viable parameter space allowed by various experiments.

Different from the minimal inelastic Dirac DM case with $Q_{\tilde{\chi}_2} = -1$ [52], we study the extended scenario with $Q_{\tilde{\chi}_2}$ as a free parameter in this paper. The different value of $Q_{\tilde{\chi}_2}$ leads to distinct mechanisms governing dark matter production. For instance, the coannihilation mechanism is the only viable dominate channel in the minimal scenario [52]. In contrast, cospattering and conversion mechanisms could become the dominant channel in the extended case when $g_\chi = g' \times Q_{\tilde{\chi}_2} \sim \mathcal{O}(1)$. It is worth noting that previous study of inelastic Dirac DM mainly concentrated on the MeV mass scale [50,52], however, our focus lies on the GeV–TeV mass scale.

The structure of this paper is organized as follows. In Sec. II, we provide a brief introduction to the theoretical model employed in our study. The calculation of relic density as well as the associated phenomenological constraints in the resonance scenario are discussed in Sec. III. Next, we investigate the secluded scenario in Sec. IV. Finally, we summarize the results in Sec. V.

II. THE MODEL

We extend two vectorlike Dirac fermions $\tilde{\chi}_1$ and $\tilde{\chi}_2$ beyond the SM. Among them, only $\tilde{\chi}_2$ carries a nonzero charge $Q_{\tilde{\chi}_2}$ under the $U(1)_{B-L}$ symmetry, and $\tilde{\chi}_1$ is

considered to be neutral. Meanwhile, these two dark fermions are stipulated to be Z_2 odd to ensure the stability of DM. In this configuration, when a dark scalar ϕ with $U(1)_{B-L}$ charge $Q_\phi = -Q_{\tilde{\chi}_2}$ acquires a vacuum expectation value $\langle \phi \rangle = v_\phi$ [50,52], the Yukawa term $y \phi \tilde{\chi}_1 \tilde{\chi}_2$ will naturally induce mass mixing $\delta m \tilde{\chi}_1 \tilde{\chi}_2$ with $\delta m = y v_\phi$.

The mass eigenstates χ_1 and χ_2 could be formed through a transformation,

$$\begin{pmatrix} \chi_1 \\ \chi_2 \end{pmatrix} = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \tilde{\chi}_1 \\ \tilde{\chi}_2 \end{pmatrix}. \quad (1)$$

Here, we consider that m_{χ_1} is slightly smaller than m_{χ_2} , thus χ_1 is regarded as the DM candidate. In principle, the scalar portal interactions also contribute to the DM relic abundance. However, the corresponding Yukawa coupling y is subjected to dual suppression from small mixing θ and mass splitting $m_{\chi_2} - m_{\chi_1}$ [52], which is favored by the conversion mechanism in this study. For simplicity, we further assume that the mass of the new scalar m_ϕ is much larger than $m_{\chi_{1,2}}$, so the contribution resulting from ϕ portal interaction is significantly smaller than that of the Z' portal [52].

In the mass eigenstates, the Z' portal interactions could be expressed as

$$\begin{aligned} \mathcal{L} \supset & +Z'_\mu g_\chi \left(\cos^2 \theta \overline{\chi}_2 \gamma^\mu \chi_2 - \frac{\sin 2\theta}{2} \overline{\chi}_2 \gamma^\mu \chi_1 - \frac{\sin 2\theta}{2} \overline{\chi}_1 \gamma^\mu \chi_2 \right. \\ & \left. + \sin^2 \theta \overline{\chi}_1 \gamma^\mu \chi_1 \right) - Z'_\mu g' Q_f \bar{f} \gamma^\mu f \end{aligned} \quad (2)$$

with g' the $U(1)_{B-L}$ gauge coupling and Q_f the $U(1)_{B-L}$ charge of the SM fermion f .

In such an inelastic model, we investigate the DM phenomenology with mass above the GeV scale. The specific execution based on the magnitude of $m_{Z'}$ and m_{χ_2} will be categorized into (1) the Z' resonance [47] and (2) secluded scenarios [48]. The total free parameters involved in the subsequent calculations are as follows:

$$\{m_{\chi_1}, \Delta_\chi, r_{Z'}, g', g_\chi, \theta\}, \quad (3)$$

where $\Delta_\chi = (m_{\chi_2} - m_{\chi_1})/m_{\chi_1}$ and $r_{Z'} = m_{Z'}/m_{\chi_2}$.

III. RESONANCE SCENARIO

A. Relic density

As will shown in Sec. III B, to evade the stringent collider constraints on Z' , the mass ratio $r_{Z'}$ has to be near the resonance region, i.e., $r_{Z'} \simeq 2$. We set the mass ratio $r_{Z'} = 2$ in this resonance scenario, which maximizes the suppression of the gauge coupling g' under the requirement of reproducing the observed dark matter relic density. The reactions associated with the dark particles include two

types: (co)annihilation and conversion. The former primarily involves s -channel processes mediated by Z' , i.e., $\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}$, $\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}$, and $\chi_2 \chi_1 \rightarrow f \bar{f}$. The latter encompasses not only the inelastic scattering $\chi_2 f \rightarrow \chi_1 f$ and the three-body decay $\chi_2 \rightarrow \chi_1 f \bar{f}$ which determine the

coscattering mechanism, but also the self-interactions $\chi_2 \chi_i \rightarrow \chi_1 \chi_j$ within the dark sector, which corresponds to the conversion mechanism.

The abundances of χ_1 and χ_2 could be calculated by numerically solving the Boltzmann equations,

$$\begin{aligned} \frac{dY_{\chi_1}}{dz} = & -\frac{s}{\mathcal{H}z} \left[\langle \sigma v \rangle_{\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}} \left(Y_{\chi_1}^2 - (Y_{\chi_1}^{\text{eq}})^2 \right) + \langle \sigma v \rangle_{\chi_2 \chi_1 \rightarrow f \bar{f}} \left(Y_{\chi_2} Y_{\chi_1} - Y_{\chi_2}^{\text{eq}} Y_{\chi_1}^{\text{eq}} \right) - \langle \sigma v \rangle_{\chi_2 f \rightarrow \chi_1 f} \left(Y_{\chi_2} Y_f^{\text{eq}} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_f^{\text{eq}} \right) \right. \\ & - \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_2}^2 - \frac{(Y_{\chi_2}^{\text{eq}})^2}{(Y_{\chi_1}^{\text{eq}})^2} Y_{\chi_1}^2 \right) - \langle \sigma v \rangle_{\chi_1 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_1} Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1}^2 \right) - \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_2} \left(Y_{\chi_2}^2 - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_{\chi_2} \right) \\ & \left. - \frac{\Gamma_{\chi_2 \rightarrow \chi_1 f \bar{f}}}{s} \left(Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} \right) \right], \end{aligned} \quad (4)$$

$$\begin{aligned} \frac{dY_{\chi_2}}{dz} = & -\frac{s}{\mathcal{H}z} \left[\langle \sigma v \rangle_{\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}} \left(Y_{\chi_2}^2 - (Y_{\chi_2}^{\text{eq}})^2 \right) + \langle \sigma v \rangle_{\chi_2 \chi_1 \rightarrow f \bar{f}} \left(Y_{\chi_2} Y_{\chi_1} - Y_{\chi_2}^{\text{eq}} Y_{\chi_1}^{\text{eq}} \right) + \langle \sigma v \rangle_{\chi_2 f \rightarrow \chi_1 f} \left(Y_{\chi_2} Y_f^{\text{eq}} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_f^{\text{eq}} \right) \right. \\ & + \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_2}^2 - \frac{(Y_{\chi_2}^{\text{eq}})^2}{(Y_{\chi_1}^{\text{eq}})^2} Y_{\chi_1}^2 \right) + \langle \sigma v \rangle_{\chi_1 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_1} Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1}^2 \right) + \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_2} \left(Y_{\chi_2}^2 - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_{\chi_2} \right) \\ & \left. + \frac{\Gamma_{\chi_2 \rightarrow \chi_1 f \bar{f}}}{s} \left(Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} \right) \right], \end{aligned} \quad (5)$$

where $z = m_{\chi_1}/T$, the entropy density $s = 2\pi^2 g_s T^3/45$, and the Hubble expansion rate is defined as $\mathcal{H} = \sqrt{4\pi^3 g_*/45} T^2/m_{pl}$ with the Planck mass $m_{pl} = 1.22 \times 10^{19}$ GeV. g_s and g_* are the number of relativistic degrees of freedom for the entropy density and energy density, respectively. The thermal average cross sections $\langle \sigma v \rangle$ of various channels are calculated numerically by MICROMEGAS [72,73].

Because of the kinematic prohibition of two-body decay $\chi_2 \rightarrow \chi_1 Z'$, the predominant decay process of χ_2 is the three-body decay $\chi_2 \rightarrow \chi_1 Z'^* \rightarrow \chi_1 f \bar{f}$. The corresponding thermal decay width is written as

$$\Gamma_{\chi_2 \rightarrow \chi_1 f \bar{f}} = \frac{\mathcal{K}_1 \left(\frac{m_{\chi_2}}{m_{\chi_1}} x \right)}{\mathcal{K}_2 \left(\frac{m_{\chi_2}}{m_{\chi_1}} x \right)} \tilde{\Gamma}_{\chi_2 \rightarrow \chi_1 f \bar{f}}, \quad (6)$$

where $\mathcal{K}_{1,2}$ are modified Bessel functions of the second kind. In the limit of small mass splitting, the three-body decay width can be estimated as [74]

$$\tilde{\Gamma}_{\chi_2 \rightarrow \chi_1 f \bar{f}} \simeq \frac{N_c^f g_\chi^2 g'^2 Q_f^2 \Delta_\chi^5 m_{\chi_1}^5 \sin^2 2\theta}{120\pi^3 m_{Z'}^4} \times \theta'(\Delta_\chi - 2m_f), \quad (7)$$

where N_c^f is the color number of f , and $Q_f = 1/3$ and -1 for quarks and leptons. θ' is the Heaviside theta function. More precise results are obtained numerically through MICROMEGAS in this paper. The abundance of χ_1 and χ_2

at thermal equilibrium can be expressed as [72]

$$Y_{\chi_1}^{\text{eq}} = \frac{45z^2}{2\pi^4 g_s} \mathcal{K}_2(z), \quad Y_{\chi_2}^{\text{eq}} = \frac{45z^2}{2\pi^4 g_s} \left(\frac{m_{\chi_2}}{m_{\chi_1}} \right)^2 \mathcal{K}_2 \left(\frac{m_{\chi_2}}{m_{\chi_1}} z \right). \quad (8)$$

Y_f^{eq} takes the value of 0.238.

According to the classification in Ref. [60], a benchmark that falls within the coscattering regime primarily satisfies three conditions at the freeze-out temperature: (1) χ_1 is in kinetic equilibrium with the SM. (2) There is no chemical potential for χ_1 . (3) The last reaction to decouple, which changes the number density of χ_1 , is exchange reactions between χ_1 and χ_2 . Regarding the first condition, since the elastic reaction $\Gamma_{\chi_1 f \rightarrow \chi_1 f}$ in our scenario is suppressed by the small mixing θ , we assume that χ_1 stays in the kinetic equilibrium with SM via the intense inelastic scattering process $\chi_2 f \rightarrow \chi_1 f$, which can be parametrized by $\Gamma_{\chi_2 f \rightarrow \chi_1 f} \gg \mathcal{H}$ [54,60]. Obtaining precise results requires solving the full unintegrated Boltzmann equations, which may introduce an $\mathcal{O}(10\%)$ distinction [55] compared to χ_1 not being in kinetic equilibrium. The requirement (2) is satisfied through $\Gamma_{\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}} > \Gamma_{\chi_2 f \rightarrow \chi_1 f}$. The strong reaction of $\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}$ rapidly consumes χ_2 , thereby disrupting the chemical equilibrium between χ_1 and χ_2 in the process $\chi_2 f \leftrightarrow \chi_1 f$. As for condition (3), it suffices that $\Gamma_{\chi_2 f \rightarrow \chi_1 f}$ is greater than $\Gamma_{\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}}$.

The process $\chi_2 \rightarrow \chi_1 f \bar{f}$ that is closely related to cospattering has an ignored contribution in this paper. The impact of $\chi_2 Z' \rightarrow \chi_1 Z'$ is also very small in the resonance scenario, because $m_{Z'} > m_{\chi_{1,2}}$ causes $Y_{Z'}^{\text{eq}}$ to undergo an exponential suppression earlier than $Y_{\chi_{1,2}}^{\text{eq}}$. In this way, $Y_{Z'}^{\text{eq}}$ is already diminished to a negligibly small value at the decoupling temperature of DM, naturally suppressing the contribution of process $\chi_2 Z' \rightarrow \chi_1 Z'$ in the Boltzmann equations. Therefore, we only need to focus on the inelastic scattering $\chi_2 f \rightarrow \chi_1 f$. Moreover, in the cospattering regime, both the WIMP-like pair annihilation $\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}$ and coannihilation $\chi_1 \chi_2 \rightarrow f \bar{f}$ processes have the reaction rates smaller than $\mathcal{H}$ to ensure that χ_1 has departed from thermal equilibrium at the freeze-out temperature. On the whole, the relevant processes approximately satisfy the relationship $\Gamma_{\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}} > \Gamma_{\chi_2 f \rightarrow \chi_1 f} \gg \mathcal{H} \gtrsim \Gamma_{\chi_1,2 \bar{\chi}_1 \rightarrow f \bar{f}}$ at the freeze-out temperature. Regarding the conversion process $\chi_i \chi_2 \rightarrow \chi_j \chi_1$, $\Gamma_{\chi_2 f \rightarrow \chi_1 f} > \Gamma_{\chi_i \chi_2 \rightarrow \chi_j \chi_1}$ suffices to guarantee that the contribution of cospattering is dominant, while its relative magnitude compared to $\mathcal{H}$ becomes irrelevant.

By employing a similar judgment method, when the dominant $\chi_2 f \rightarrow \chi_1 f$ is exceeded by $\chi_i \chi_2 \rightarrow \chi_j \chi_1$, it results in a conversion regime. At this point, the reaction rates of these processes satisfy the relationship $\Gamma_{\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}} > \Gamma_{\chi_i \chi_2 \rightarrow \chi_j \chi_1} \gg \mathcal{H} \gtrsim \Gamma_{\chi_1,2 \bar{\chi}_1 \rightarrow f \bar{f}}$. In other cases where either $\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}$ or $\chi_2 \chi_1 \rightarrow f \bar{f}$ makes the most significant contribution, this can be referred to as coannihilation, which requires $\Gamma_{\chi_1,2 \bar{\chi}_2 \rightarrow f \bar{f}} \sim \mathcal{H} > \Gamma_{\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}}$. As for the coannihilation phase, the respective reaction rates of the cospattering process $\chi_2 f \rightarrow \chi_1 f$ and the conversion process $\chi_i \chi_2 \rightarrow \chi_j \chi_1$ may be either greater than or less than $\Gamma_{\chi_1,2 \bar{\chi}_2 \rightarrow f \bar{f}}$ as long as the conditions for determining cospattering and conversion regimes are not satisfied.

We utilize three benchmark points corresponding to the cospattering, the conversion and the coannihilation phases to present the evolutions of dark fermion abundances in Fig. 1, which also aims to elucidate how to distinguish them. The thermal rates of $2 \rightarrow 2$ processes in Figs. 1(d)–1(f) are denoted as $\Gamma_i = n_a^{\text{eq}} \langle \sigma v \rangle_i$ with n_a^{eq} the number density of particle a at the thermal equilibrium. Figures 1(a) and 1(d), 1(b) and 1(e), and 1(c) and 1(f) satisfy the criteria for determining cospattering, conversion, and coannihilation, respectively. Among these phases, cospattering is the first to exhibit an obvious deviation from thermal equilibrium, occurring approximately at $\Gamma_{\chi_2 \chi_1 \rightarrow f \bar{f}} \sim \mathcal{H}$. Then Y_{χ_1} subsequently continues to show a rapid decline until the depletion of inelastic reaction $\chi_2 f \rightarrow \chi_1 f$ [72]. In contrast, the decoupling of conversion and coannihilation occurs at a later stage. And for the same m_{χ_1} , significant differences are observed in g' due to variations in the defined conditions.

Furthermore, the inability to achieve freeze-out raises concerns with tiny g' and θ . Specifically, DM χ_1 cannot

reach thermal equilibrium with the SM bath. This can be parametrized as $\Gamma_{\chi_2 \chi_1 \rightarrow f \bar{f}} / \mathcal{H} < 1$, as the contribution of pair annihilation $\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}$ is deemed negligible due to its minimal impact. Here, we use MICROMEGAS to perform calculations over a broad range of $g' \in [10^{-10}, 10^{-2}]$, $\Delta_\chi \in [10^{-3}, 10^{-1}]$, and $g_\chi \in [0.1, 1]$. We report that the lower limit for freeze-out of χ_1 is $\theta \gtrsim \mathcal{O}(10^{-6})$ when $m_{\chi_1} \sim \text{TeV}$. For the GeV scale, this result can drop to $\mathcal{O}(10^{-8})$. Therefore, in the subsequent phenomenological study, the selection of parameters adheres to the freeze-out condition.

B. Phenomenology of Z'

The phenomenology in the aspect of Z' mainly arises from searches at various colliders. At the scale above GeV that we are focusing on, constraints come from the probe of dark photons at current experiments BABAR [42,43] and LHCb [44,45], dilepton searches of large electron positron (LEP) at Z peak [76,77], as well as of CMS and ATLAS [46,78]. These experimental limits collectively provide stringent constraints on coupling g' as illustrated in the shaded gray area of Fig. 2, in which the parameter space with $g' \gtrsim 3 \times 10^{-4}$ is excluded. The future Belle II sensitivity for long-lived [79] and invisible [80] Z' are colored by an orange dashed line, which could probe $m_{Z'} \lesssim 8 \text{ GeV}$. The projection sensitivity of searching for dark photons at FCC-ee [81] with $\sqrt{s} = 90 \text{ GeV}$ is colored by a pink dashed line, which is likely to detect $m_{Z'} \in [10, 100] \text{ GeV}$ with $g' \sim 6 \times 10^{-5}$. For $m_{Z'} \gtrsim 150 \text{ GeV}$ and $g' \gtrsim 4.3 \times 10^{-4}$, the sensitivity of future CMS and ATLAS with a projected 3 ab^{-1} luminosity in the search for dileptons improves by approximately 1 order of magnitude compared to the current results [76], which is marked with a purple dashed curve.

Figure 2 illustrates the impact of constraints from Z' on various phases of benchmark cases. In Fig. 2(a), the parameters θ , g_χ , and Δ_χ are fixed as 5×10^{-3} , 1, and 10^{-2} , respectively, while the changing parameter is $r_{Z'}$. The conventional nonresonant case of $r_{Z'} = 3$ is dominated by coannihilation, and almost all portions except for $m_{Z'} \sim 100 \text{ GeV}$ are excluded by the current constraints. When $r_{Z'}$ decreases to the resonance region $2 \lesssim r_{Z'} \lesssim 2.07$, as shown by the red and blue curves, it is indeed possible to overcome these limitations. Furthermore, as one approaches the extreme resonance with $r_{Z'} = 2$, the required g' decreases, which facilitates the emergence of conversion and cospattering phases. In the extreme resonance case $r_{Z'} = 2$, coannihilation favors $g' \lesssim 10^{-3}$, while the corresponding $m_{Z'}$ does not exceed 2850 GeV. A larger g' naturally leads to a more intense annihilation reaction of χ_2 pairs. When $\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}$ exceeding the conversion process $\chi_i \chi_2 \rightarrow \chi_j \chi_1$, it belongs to the conversion phase based on the specified criteria, which distributes at $g' \in [10^{-3}, 5.8 \times 10^{-3}]$ with $m_{Z'} \in [1560, 2850] \text{ GeV}$.

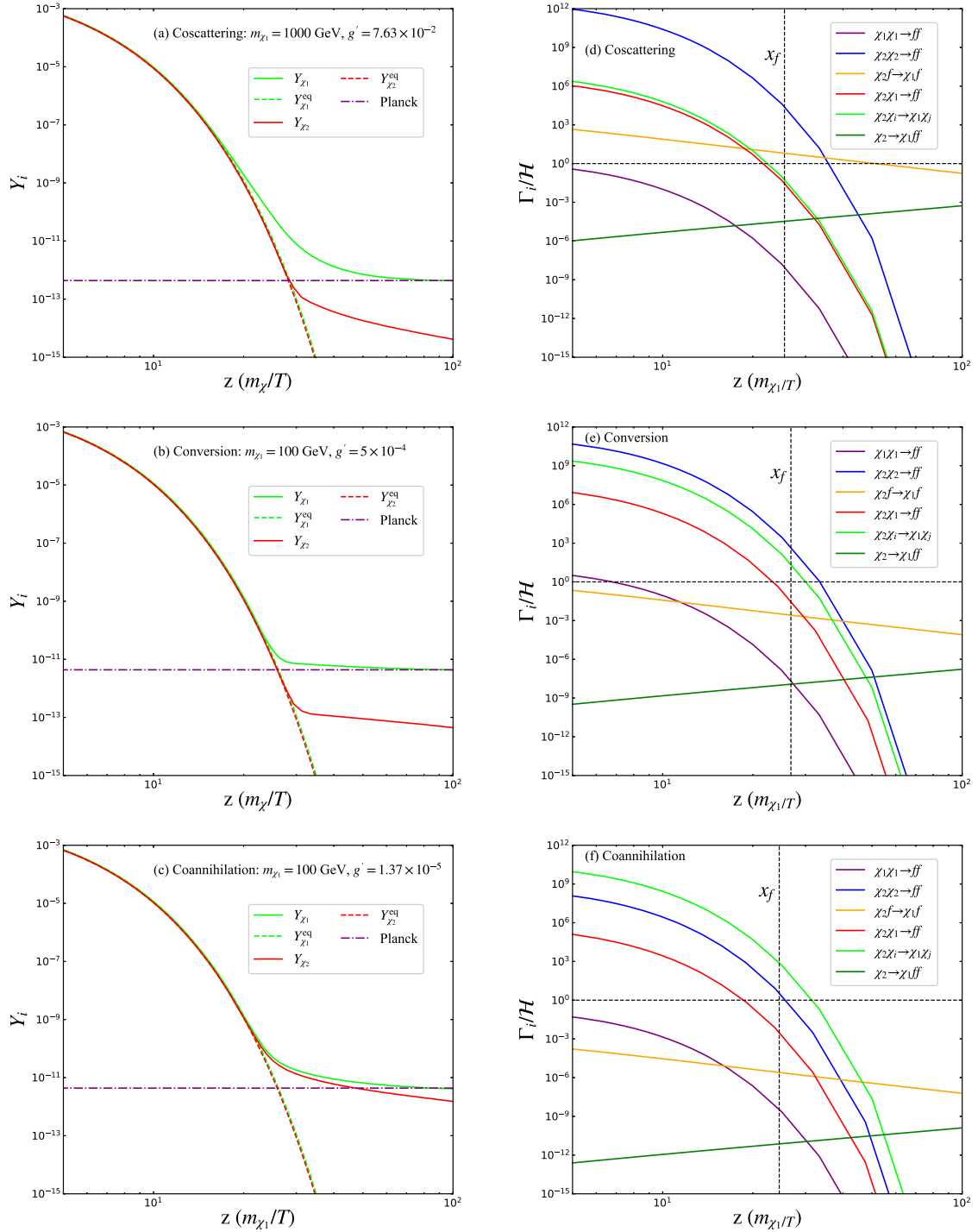


FIG. 1. The evolutions of various abundances Y_i for (a) coscattering, (b) conversion, and (c) coannihilation benchmarks in the resonance scenario. (d)–(f) The thermal rates of various reaction processes in three different phases. We fix $\Delta_\chi = 10^{-2}$, $r_{Z'} = 2$, and $\theta = 5 \times 10^{-4}$. The solid green and red lines in (a)–(c) represent the abundance of χ_1 and χ_2 , whereas the dashed lines are their thermal equilibrium. Purple dot-dashed curve is the observation of DM $\Omega_{\chi_1} h^2 = 0.12$ [75]. (d)–(f) The black vertical dashed line corresponds to the thermal decoupling temperature when $Y_{\chi_1}/Y_{\chi_1}^{\text{eq}} = 2.5$, and the horizontal black line is $\Gamma_i = \mathcal{H}$. Moreover, $\chi_i \chi_j \rightarrow \chi_k \chi_l$ is the sum of conversion channels $\chi_2 \chi_2 \rightarrow \chi_1 \chi_1$, $\chi_1 \chi_2 \rightarrow \chi_1 \chi_1$, and $\chi_2 \chi_2 \rightarrow \chi_1 \chi_2$.

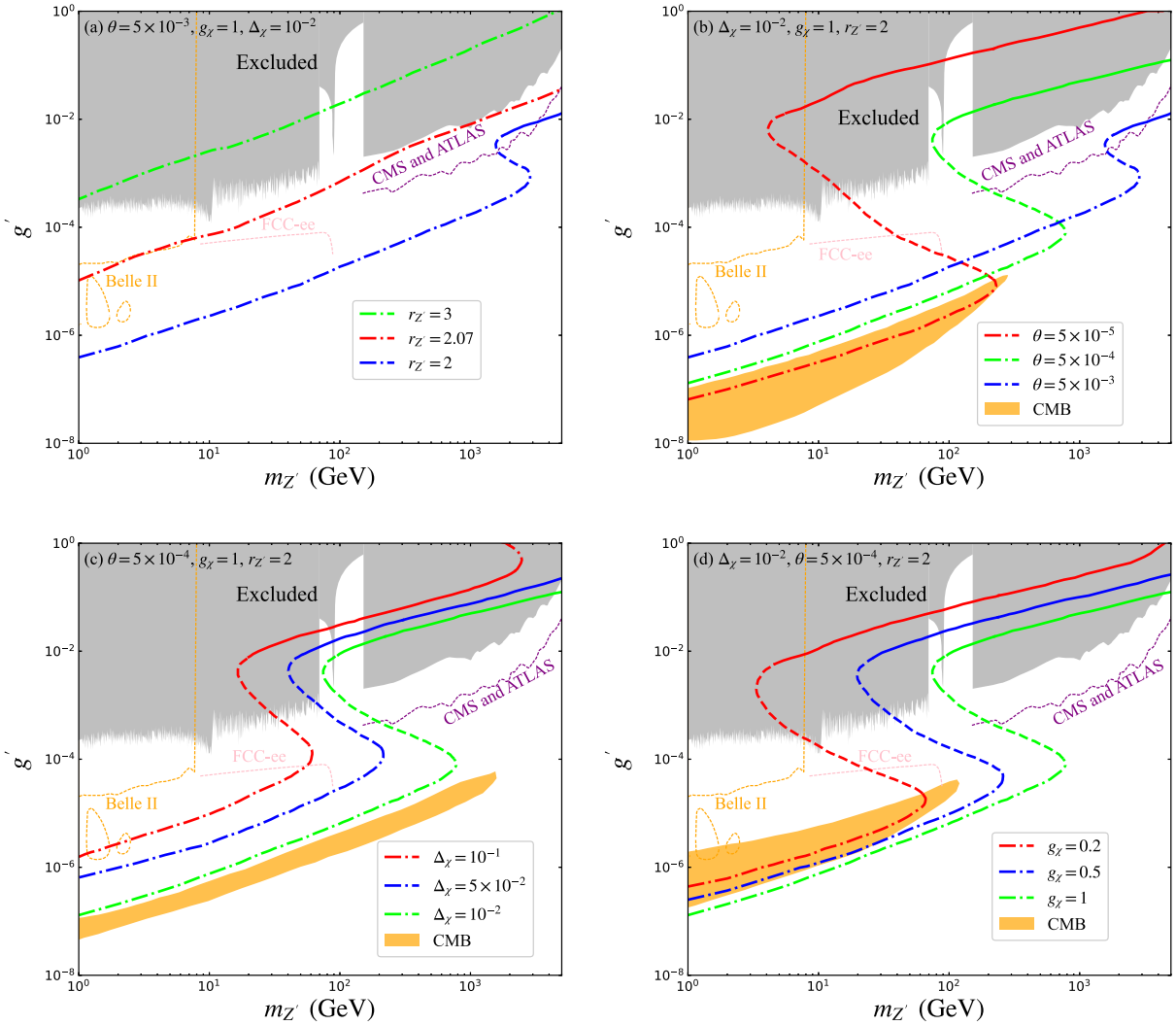


FIG. 2. Constraints on the $m_{Z'} - g'$ parameter space in the resonance scenario. (a)–(d) Different selections with fixed parameters. Each panel features red, green, and blue lines representing three distinct benchmarks, all of which are consistent with DM observation. The solid, dashed, and dot-dashed parts on each line correspond to coscattering, conversion, and coannihilation phases, respectively. The gray shaded area indicates the current exclusion on Z' by various colliders. The future sensitivities of Belle II, FCC-ee, CMS, and ATLAS are illustrated by the orange, pink, and purple dashed lines, respectively. The orange shaded areas appearing in (b)–(d) indicate the promising coannihilation region that can be probed by future CMB on long-lived χ_2 .

For $g' \gtrsim 5.8 \times 10^{-3}$, the contribution brought by inelastic process $\chi_2 f \rightarrow \chi_1 f$ is greater than that of $\chi_i \chi_2 \rightarrow \chi_j \chi_1$, so it becomes the coscattering phase. In this phase, $m_{Z'}$ continuously rises to 5000 GeV as g' increases to 0.013. Obviously, there is an approximate proportional relationship between g' and $m_{Z'}$ in the coannihilation and coscattering phases, whereas the situation in conversion is quite the opposite. In the future, the projected sensitivity of CMS and ATLAS is expected to capture conversion and coscattering within $1610 \lesssim m_{Z'} \lesssim 3300$ GeV when $g' \sim \mathcal{O}(10^{-3})$. Furthermore, by comparing the parameter $r_{Z'}$, it can be observed that the current stringent constraints compel us to select a more optimistic Z' resonance scenario, which serves as the focal point for the subsequent three panels.

In Fig. 2(b), we vary θ to obtain different benchmark cases, while keeping other parameters as $\Delta_\chi = 10^{-2}$, $g_\chi = 1$, and $r_{Z'} = 2$. As θ increases from 5×10^{-5} to 5×10^{-3} , coannihilation exhibits an upward trend on g' , where its maximum achievable value at $(m_{Z'}, g')$ rises from (225 GeV, 8.6×10^{-6}) to (2850 GeV, 10^{-3}). In contrast, both conversion and coscattering demonstrate a distinct rightward shift tendency, with the critical points of these two phases increasing from (6.5 GeV, 1.42×10^{-2}) to (1943 GeV, 5.8×10^{-3}). For coannihilation, the required g' is not subject to any exclusion at present, and the future Belle II is sensitive to it with $m_{Z'} \sim \mathcal{O}(1)$ GeV when $\theta \sim \mathcal{O}(10^{-2})$. The conversion of $\theta \lesssim 5 \times 10^{-4}$ is excluded within the range of $g' \gtrsim 3 \times 10^{-4}$ and $m_{Z'} \lesssim 100$ GeV.

However, under the exclusion limits, there is significant potential for conversion to be captured by future Belle II and FCC-ee. Furthermore, when $m_{Z'} \gtrsim 150$ GeV with $g' \sim \mathcal{O}(10^{-3})$, conversion is also promising with $\theta \in [5 \times 10^{-4}, 5 \times 10^{-3}]$ for CMS and ATLAS. Because of the large magnitude of g' , most cospattering regions are not permitted. In the small unrestricted area, we anticipate that there is hope at $m_{Z'} \gtrsim 1000$ GeV when $\theta \sim \mathcal{O}(10^{-3})$ and the narrow region slightly below 100 GeV.

Figure 2(c) researches the impact of Δ_χ , while $\theta = 5 \times 10^{-4}$, $g_\chi = 1$, and $r_{Z'} = 2$ are fixed. The three phases shift toward larger $m_{Z'}$ as Δ_χ decreases, among them, coannihilation, conversion, and cospattering are distributed at $g' \lesssim 10^{-4}$, $10^{-4} \lesssim g' \lesssim 10^{-2}$, and $g' \gtrsim 10^{-2}$, respectively. However, when $\Delta_\chi = 0.1$, the excessive mass splitting prevents cospattering from being valid within the range of $m_{Z'} \gtrsim 2450$ GeV. Therefore, returning to small $m_{Z'}$, $g' \gtrsim 0.5$ can only satisfy dark matter observations through coannihilation. Under current constraints, it is regrettable that a small portion of conversion with $\Delta_\chi \gtrsim 10^{-2}$ and the bulk of cospattering are restricted at larger g' . The remaining small g' that are likely to be detected by future experiments include coannihilation when $\Delta_\chi \gtrsim 10^{-1}$ and conversion with relatively free Δ_χ .

In Fig. 2(d), the increase of g_χ also leads to a rightward shift of the curve. Similar to that, in Fig. 2(c), a portion of the conversion will be excluded as long as $g_\chi \lesssim 1$, while nearly all cospatterings are destined to face exclusion. When $g_\chi \ll 0.2$, coannihilation with $g' \sim \mathcal{O}(10^{-6})$ has the potential to be identified by Belle II. Meanwhile, conversion is more encouraging, since future experiments hold greater promise for capturing it when $g_\chi \sim \mathcal{O}(0.1)$. Too large g_χ may lead to nonperturbative issues; therefore, such cases with $g \geq 1$ are not within the scope of our consideration.

In general, within the resonance scenario, various promising phases exhibit distinct parameter selections. For instance, coannihilation favors $\theta \sim \mathcal{O}(10^{-2})$, $\Delta_\chi \gtrsim 10^{-1}$, and $g_\chi \ll 0.2$. Conversion with $\theta \lesssim 5 \times 10^{-3}$, $g_\chi \sim \mathcal{O}(0.1)$, and relatively unlimited Δ_χ demonstrates significant potential for future experiments. The cospattering situation is not very optimistic owing to the limitations imposed by current constraints. In addition, when the gauge coupling g' and the mixing angle θ are small enough, the dark partner χ_2 becomes long-lived, which can be probed by future cosmic microwave background (CMB) experiments. In Fig. 2, we depict the CMB sensitive regions in orange, which will be discussed in detail in Sec. III E.

C. Phenomenology of χ_1

For DM above the GeV scale, the spin-independent scattering cross section with nucleons is strictly

constrained by current direct detection experiments, such as DarkSide-50 [39], PandaX-4T [40], and LZ [41]. The combined exclusion region is represented by the orange shaded region in Fig. 3, in which the strongest constraint is located at $m_{\chi_1} \sim 30$ GeV with $\sigma_{\text{SI}} \sim 2 \times 10^{-48} \text{ cm}^2$. The future sensitivity represented by the orange dashed curve is provided by DarkSide-LowMass [82], SuperCDMS [83], and LZ [84], which is nearly an order of magnitude lower than the current results.

In this model, the spin-independent cross section can be calculated by

$$\sigma_{\text{SI}} = \frac{m_{\chi_1}^2 m_n^2 \sin^4 \theta g_\chi^2 g'^2}{\pi m_{Z'}^4 (m_{\chi_1} + m_n)^2}, \quad (9)$$

where the mass of nucleons $m_n \simeq 0.939$ GeV. Compared with the traditional Z' portal DM, the cross section is further suppressed by the small mixing angle θ . The predictions are shown in Fig. 3 as red, blue, and green lines for the benchmark cases.

In Fig. 3(a), although the future outcomes are expected to capture $m_{\chi_1} \sim \mathcal{O}(10)$ GeV in the nonresonant case, the constraints on Z' from the collider already disfavor such a region. In the resonance cases, the σ_{SI} for three phases decreases as m_{χ_1} increases. Being suppressed by θ , σ_{SI} reaches its maximum $6.8 \times 10^{-52} \text{ cm}^2$ for $r_{Z'} = 2$ at 1 GeV, which is nearly 6 orders of magnitude lower than future sensitivity. Therefore, such a resonance case is challenging for the direct detection experiments.

In Figs. 3(b)–3(d), all benchmarks predict $\sigma_{\text{SI}} \lesssim 10^{-50} \text{ cm}^2$. None of them is within the reach of future experiments. However, it is evident that adjusting certain parameters can alter this situation. For instance, when $\theta \gg 5 \times 10^{-3}$ in Fig. 3(b), it can be inferred that coannihilation is completely dominant based on the trend in changing θ . As θ increases, m_{χ_1} may be detectable over a wide range.

When the mass splitting $\delta = m_{\chi_2} - m_{\chi_1}$ is less than $\mathcal{O}(100)$ keV, the inelastic scattering $\chi_1 n \rightarrow \chi_2 n$, in principle, can yield an observable signature [49]. However, such an inelastic scattering cross section is also suppressed by the small mixing angle θ , which makes it unpromising [50]. In this paper, we consider $\delta = m_{\chi_1} \Delta_\chi \gtrsim 10^{-3}$ GeV, so it is far above the experimental sensitive region [85].

Within the resonance scenario, DM pairs could annihilate into SM fermions $f\bar{f}$ via the Z' mediator. The present annihilation cross section $\langle \sigma v \rangle$ is constrained by outcomes from indirect detection experiments. As the leptonic final state $\ell^+ \ell^-$ is the dominant annihilation channel, we take the limits of the $e^+ e^-$ final state as an example, which is illustrated in Fig. 4. The existing constraints with $m_{\chi_1} \lesssim 5$ GeV come from experiments involving XMM-NEWTON x-rays [86] and CMB (s-wave) [87,88] observations, while the results for larger m_{χ_1} are taken from

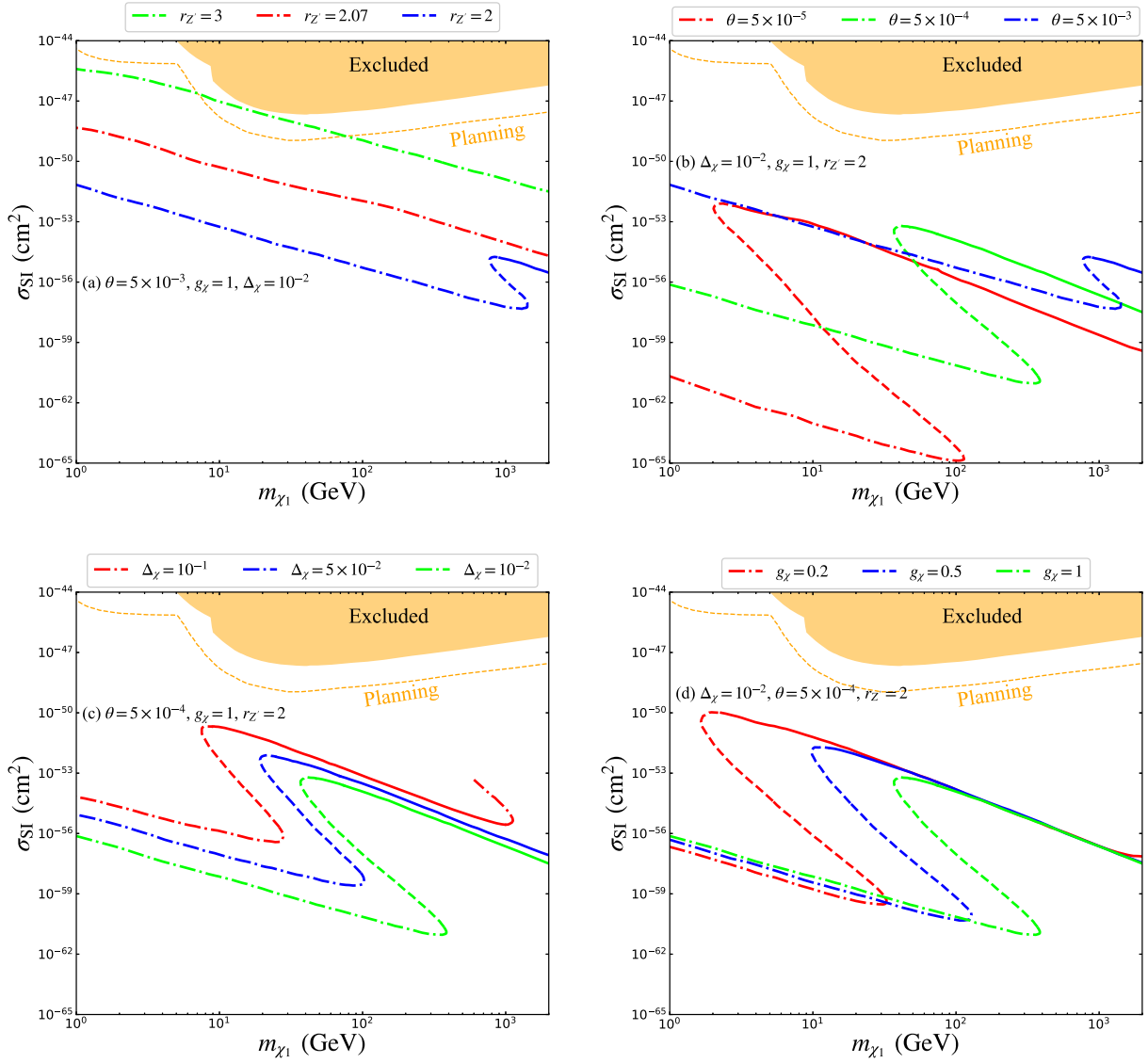


FIG. 3. The constraints of direct detection experiments in the resonance scenario. The fixed parameters in (a)–(d), as well as the selections of benchmarks in each panel, are consistent with those presented in Fig. 2. The solid, dashed, and dot-dashed components of each benchmark line still correspond to coscattering, conversion, and coannihilation phases. The orange region and dashed line represent the results of current and future direct detection experiments.

literature [89,90], which are the convolutions of the bounds from AMS positron [91,92], Fermi-LAT dwarfs [93], and H.E.S.S. Galactic Center (GC) observations [94,95]. These constraints collectively exclude the purple shaded area with $\langle\sigma v\rangle \gtrsim 10^{-28} \text{ cm}^3/\text{s}$. The purple dashed line represents the sensitivities of the future MeV telescopes AMEGO [96–98], E-ASTROGAM [99,100], and MAST [101] in probing weak-scale DM, which is derived from [102]. The future limit is roughly 2 orders of magnitude lower than the current one. The maximum detection capability is observed at the GeV scale with $\langle\sigma v\rangle \sim 10^{-31} \text{ cm}^3/\text{s}$.

In this model, $\langle\sigma v\rangle$ could be numerically calculated through [48]

$$\begin{aligned} \langle\sigma v\rangle_{\chi_1\bar{\chi}_1 \rightarrow f\bar{f}} &\simeq \frac{\sin^4\theta g_{\chi'}^2 g^2}{2\pi} \sum_f N_c^f Q_f^2 \frac{2m_{\chi_1}^2 + m_f^2}{(4m_{\chi_1}^2 - m_{Z'}^2)^2 + m_{Z'}^2 \tilde{\Gamma}_{Z'}^2} \\ &\times \sqrt{1 - \frac{m_f^2}{m_{\chi_1}^2}}, \end{aligned} \quad (10)$$

where $\tilde{\Gamma}_{Z'}$ represents the total decay width of Z' , which can be decomposed into

$$\tilde{\Gamma}_{Z' \rightarrow f\bar{f}} = \sum_f \frac{N_c^f g^2 Q_f^2 m_{Z'}}{12\pi} \left(1 + \frac{2m_f^2}{m_{Z'}^2}\right) \sqrt{1 - \frac{4m_f^2}{m_{Z'}^2}}, \quad (11)$$

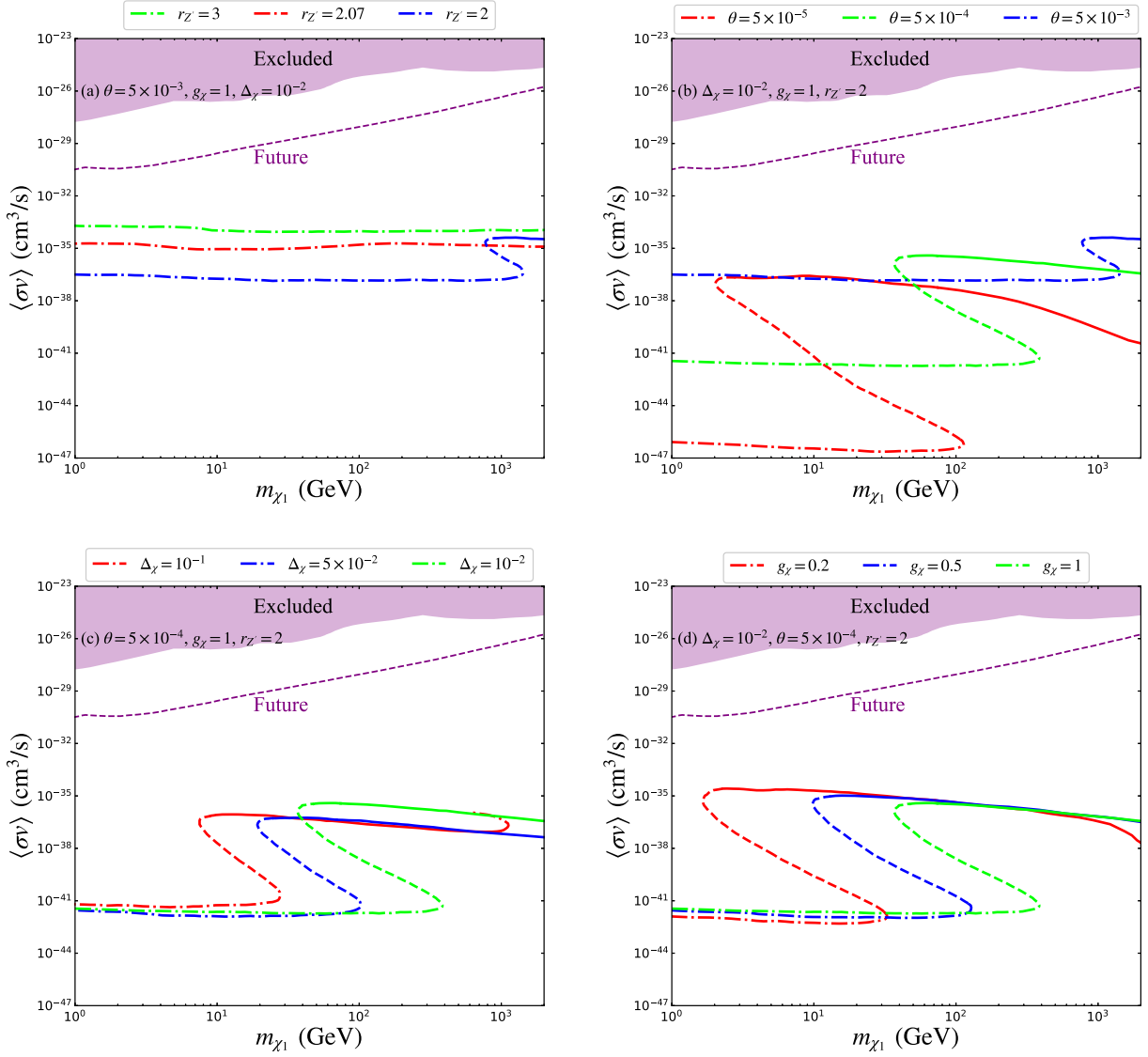


FIG. 4. The constraints of indirect detection experiments in the resonance scenario. The legends of (a)–(d) and markers of benchmark lines are consistent with those in Fig. 2. The purple region is not permitted by the existing constraints of indirect detection experiments, while the projected sensitivity of future experiments is represented as a purple dashed line.

$$\tilde{\Gamma}_{Z' \rightarrow \chi_1 \bar{\chi}_1} = \frac{g_\chi^2 \sin^4 \theta m_{Z'}}{12\pi} \left(1 + \frac{2m_{\chi_1}^2}{m_{Z'}^2} \right) \sqrt{1 - \frac{4m_{\chi_1}^2}{m_{Z'}^2}}, \quad (12)$$

$$\tilde{\Gamma}_{Z' \rightarrow \chi_2 \bar{\chi}_2} = \frac{g_\chi^2 \cos^4 \theta m_{Z'}}{12\pi} \left(1 + \frac{2m_{\chi_2}^2}{m_{Z'}^2} \right) \sqrt{1 - \frac{4m_{\chi_2}^2}{m_{Z'}^2}}, \quad (13)$$

$$\begin{aligned} \tilde{\Gamma}_{Z' \rightarrow \chi_1 \chi_2} &= \frac{g_\chi^2 \sin^2 2\theta}{48\pi m_{Z'}^5} ((m_{\chi_2}^2 - m_{\chi_1}^2)^2 \\ &\quad + m_{Z'}^2 (m_{\chi_1}^2 + m_{\chi_2}^2 - 6m_{\chi_1} m_{\chi_2} - 2m_{Z'}^2)) \\ &\quad \times \sqrt{m_{\chi_1}^2 (m_{\chi_1}^2 - 2m_{\chi_2}^2 - 2m_{Z'}^2) + (m_{Z'}^2 - m_{\chi_2}^2)^2}. \end{aligned} \quad (14)$$

The present $\langle\sigma v\rangle$ of DM annihilation is represented by the red, blue, and green curves in Fig. 4. Similar to the DM-nucleon scattering cross section in direct detection, the current $\langle\sigma v\rangle$ is also suppressed by the mixing angle θ . Consequently, all benchmarks exhibit $\langle\sigma v\rangle \lesssim 10^{-34} \text{ cm}^3/\text{s}$, which represents a detection range far beyond what future experiments can achieve. Then again, this relationship with θ implies that, as long as $\theta \gg 5 \times 10^{-3}$, coannihilation at the GeV scale first remains promising.

D. Phenomenology of χ_2

In this scenario, the phenomena induced by the three-body decay $\chi_2 \rightarrow \chi_1 f \bar{f}$ primarily arises from two aspects. First, it is the collider signature. As long as the decay length of χ_2 does not exceed the detection range of colliders [103],

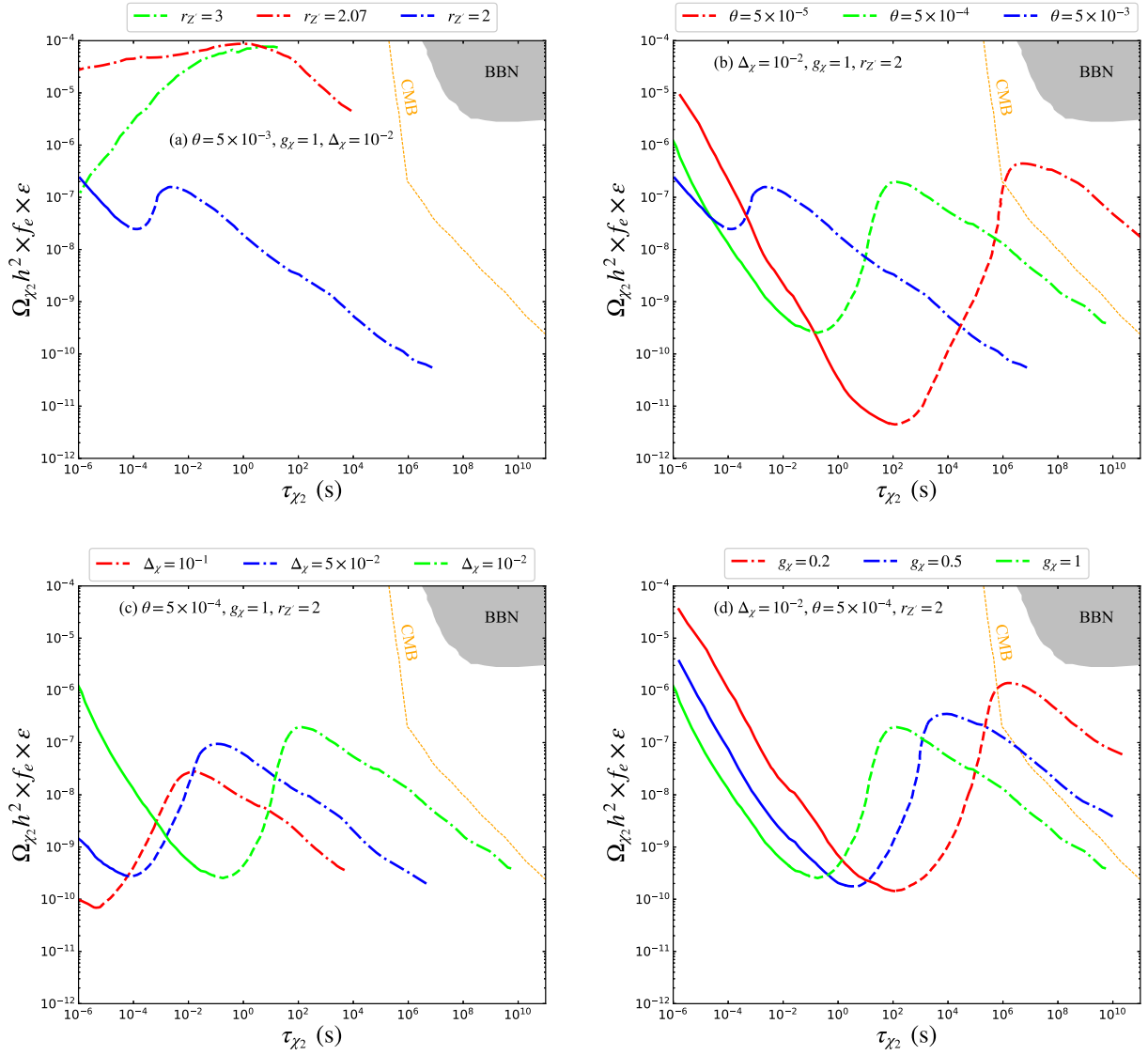


FIG. 5. Cosmological constraints on long-lived χ_2 in the resonance scenario. The horizontal axis represents the lifespan of χ_2 , while the vertical axis indicates the relative relic density. Here, f_e is the branching ratio of χ_2 decay to the e^+e^- final state, and $\epsilon = (m_{\chi_2}^2 - m_{\chi_1}^2)/2m_{\chi_2}^2$ is the fraction of the energy of χ_2 that has been transferred to electron. The legends of (a)–(d) and markers of benchmark lines are consistent with those in Fig. 2. The gray area is excluded by BBN, and the orange dashed line represents the future CMB results.

the prompt or displaced signals have the potential to be observed. For instance, $pp \rightarrow Z' \rightarrow \chi_1 + \chi_2 \rightarrow \chi_1 \chi_1 \ell^+ \ell^-$ at LHC [104] and $e^+e^- \rightarrow Z' + Z/\gamma$ and $Z' \rightarrow \chi_1 + \chi_2 \rightarrow \chi_1 \chi_1 \ell^+ \ell^- / \chi_1 \chi_1 jj$ at the lepton collider [105]. Certainly, since vertex $Z' \chi_1 \chi_2$ is suppressed in this model, multi-leptons or two displaced vertices induced by $Z' \rightarrow \chi_2 \bar{\chi}_2$ are more appropriate once kinematically allowed. Furthermore, when χ_2 decays outside the detector or the decay final states are too soft to be detected, χ_2 becomes invisible. Then, the expectation shifts to monophoton signal $e^+e^- \rightarrow Z' + \gamma \rightarrow \chi_2 \bar{\chi}_2 \gamma$ [43].

According to the cosmological constraints on χ_2 in Fig. 5, the decay length $d_{\chi_2} = c\tau_{\chi_2}$ can be roughly

estimated via the horizontal axis τ_{χ_2} . It is evident that searching for displaced vertex signals is more suitable within the coscattering regime; the other two phases can study the monophoton signals because of too long-lived χ_2 . In the extreme resonance scenario with $r_{Z'} = 2$, the decay $Z' \rightarrow \chi_2 \bar{\chi}_2$ is forbidden, thus the missing energy comes from the decay $Z' \rightarrow \chi_1 \chi_2$. However, the decay width of $Z' \rightarrow \chi_1 \chi_2$ is also suppressed by the small mixing θ and phase space, so the collider signature for the case with $r_{Z'} = 2$ is not promising.

For the long-lived dark partner, the additional energetic injection of $\chi_2 \rightarrow \chi_1 f \bar{f}$ will affect the big bang nucleosynthesis (BBN) predictions and the CMB anisotropy

power spectra. As illustrated in Fig. 5, since the area with larger τ_{χ_2} where BBN and CMB take effect is primarily distributed by coannihilation, the minor mass splitting results in $\chi_2 \rightarrow \chi_1 e^+ e^-$ being the dominant process. Thus, we consider the limits of the electron final state as an example. The gray area represents the current constraints imposed by BBN on the $e^+ e^-$ final state [106], which excludes regions for $\tau_{\chi_2} \gtrsim 3 \times 10^6$ s with relative relic density $\Omega_{\chi_2} h^2 \times f_e \times \varepsilon \gtrsim 3 \times 10^{-6}$. The orange dashed line stands for the upcoming CMB results of purely electromagnetic decay [107], which is capable of detecting regions for $\tau_{\chi_2} \gtrsim 2 \times 10^5$ s with $\Omega_{\chi_2} h^2 \times f_e \times \varepsilon \gtrsim 2 \times 10^{-10}$.

In Fig. 5, the maximum τ_{χ_2} of all benchmarks is cut off at $m_{\chi_1} = 1$ GeV. The detectability for $m_{\chi_1} \leq 1$ GeV can be seen in the previous study of [52]. It is evident that the current BBN constraints [106] do not pose any threat to benchmarks due to the small τ_{χ_2} and $\Omega_{\chi_2} h^2 \times f_e \times \varepsilon$. Coannihilation with $\theta \lesssim 5 \times 10^{-4}$, $\Delta_\chi \lesssim 10^{-2}$, and $g_\chi \lesssim 1$ is likely to be detected by future CMB [107], which corresponds to m_{χ_1} typically below TeV. Furthermore, it is noteworthy that the evolutionary trends of coannihilation in the resonance and non-resonance conditions depicted in Fig. 5(a) are completely opposite. This phenomenon arises because the conversion reaction of $\chi_i \chi_2 \rightarrow \chi_j \chi_1$ at the resonance is significantly more intense, resulting in a continued decline in the abundance of χ_2 after freeze-out, until this reaction is depleted. In more extreme cases of $\theta \lesssim 5 \times 10^{-5}$ and $g_\chi \lesssim 0.2$, conversion with $m_{\chi_1} \lesssim 100$ GeV also becomes increasingly optimistic for future CMB.

Additionally, $\chi_2 \rightarrow \chi_1 \nu \bar{\nu}$ also exhibits a considerable branching ratio. For $\tau_{\chi_2} \gtrsim \mathcal{O}(10^5)$ s with m_{χ_1} below 100 GeV, the CMB observation set no limit due to too small fraction of energy injection into SM plasma from neutrinos [108]. The neutrinos from delayed decay also contribute to the effective number of relativistic neutrino species N_{eff} [109]. The benchmarks predict $(f_\nu \varepsilon \Omega_{\chi_2} / \Omega_{\chi_1})^2 \tau_{\chi_2} \lesssim \mathcal{O}(0.1)$ s, which is far below the current Planck limit $(f_\nu \varepsilon \Omega_{\chi_2} / \Omega_{\chi_1})^2 \tau_{\chi_2} \lesssim 5 \times 10^9$ s [108].

E. Combination and discussion

In the resonance scenario, the direct and indirect detection experiments can hardly detect the benchmarks. A relatively large gauge coupling g' could induce a detectable collider signature of Z' . On the other hand, a sufficiently small g' leads to long-lived χ_2 , which could affect CMB observables due to delayed decay $\chi_2 \rightarrow \chi_1 f \bar{f}$. As shown in Figs. 2 and 5, the future CMB experiment is

primarily sensitive to coannihilation with $g' \lesssim 10^{-4}$. Therefore, the collider search for Z' and the CMB observables provides complementary pathways to probe such a scenario.

In Fig. 2, we depict the combined results. For Fig. 2(a) with $\theta = 5 \times 10^{-3}$, $g_\chi = 1$, and $\Delta_\chi = 10^{-2}$, the resonance benchmark line with $r_{Z'} = 2$ is already the lowest. We report that the insufficient τ_{χ_2} at coannihilation fails to meet the sensitivity of CMB as shown in Fig. 5(a). In contrast, coannihilation located at $m_{Z'} \lesssim \mathcal{O}(100)$ GeV and $g' \lesssim \mathcal{O}(10^{-6})$ often corresponds to smaller θ in Fig. 2(b). Consequently, very large τ_{χ_2} can easily fall within the detection range of CMB. However, excessively long τ_{χ_2} may result in χ_2 behaving as a decaying DM. To avoid this issue, we set a lower limit of $\theta \gtrsim \mathcal{O}(10^{-6})$, which also aligns with the effective range of freeze-out. Similar in Fig. 2(c), reducing Δ_χ leads to an increase of τ_{χ_2} . An overly small Δ_χ not only results in χ_2 having a lifetime longer than that of the current Universe but also restricts electromagnetic decay processes for χ_2 . Considering these factors comprehensively, we adopt a lower limit of $\Delta_\chi \gtrsim \mathcal{O}(10^{-3})$. The CMB detectable range can reach TeV scale with corresponding g' between $\mathcal{O}(10^{-7})$ and $\mathcal{O}(10^{-5})$. Finally, examining changes of g_χ illustrated in Fig. 2(d) reveals that coannihilation is distributed over larger g' as g_χ decreases. Although the detectable range of $m_{Z'}$ remains similar to that in Fig. 2(b), both being below $\mathcal{O}(100)$ GeV, the overall distribution of g' is elevated by an order of magnitude compared to that observed in Fig. 2(b).

In summary, the current constraints from colliders have excluded most of the cospin scattering. Future colliders with sensitivities at $\mathcal{O}(10^{-5}) \lesssim g' \lesssim \mathcal{O}(10^{-3})$ hold promise for conversion. However, for a lower coannihilation dominated region $\mathcal{O}(10^{-8}) \lesssim g' \lesssim \mathcal{O}(10^{-5})$ beyond future colliders' reach, the CMB observation can play a significant role.

IV. SECLUDED SCENARIO

A. Relic density

We now consider the secluded scenario with $m_{Z'} < m_{\chi_{1,2}}$ [110]. In comparison with the resonance scenario, the most notable difference of the secluded scenario is the emergence of $\chi_{1,2} \chi_{1,2} \rightarrow Z' Z'$, which is the dominant process with $g_\chi \gg g'$. For simplicity, we denote the thermal bath particles f and Z' as ζ in the following discussion. The Boltzmann equations related to the dark fermions are as follows:

$$\begin{aligned}
\frac{dY_{\chi_1}}{dz} = & -\frac{s}{\mathcal{H}z} \left[\langle \sigma v \rangle_{\chi_1 \bar{\chi}_1 \rightarrow \zeta \bar{\zeta}} \left(Y_{\chi_1}^2 - (Y_{\chi_1}^{\text{eq}})^2 \right) + \langle \sigma v \rangle_{\chi_2 \chi_1 \rightarrow \zeta \bar{\zeta}} \left(Y_{\chi_2} Y_{\chi_1} - Y_{\chi_2}^{\text{eq}} Y_{\chi_1}^{\text{eq}} \right) - \langle \sigma v \rangle_{\chi_2 \bar{\zeta} \rightarrow \chi_1 \zeta} \left(Y_{\chi_2} Y_{\zeta}^{\text{eq}} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_{\zeta}^{\text{eq}} \right) \right. \\
& - \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_2}^2 - \frac{(Y_{\chi_2}^{\text{eq}})^2}{(Y_{\chi_1}^{\text{eq}})^2} Y_{\chi_1}^2 \right) - \langle \sigma v \rangle_{\chi_1 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_1} Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1}^2 \right) - \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_2} \left(Y_{\chi_2}^2 - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_{\chi_2} \right) \\
& \left. - \frac{\Gamma_{\chi_2 \rightarrow \chi_1 f \bar{f}}}{s} \left(Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} \right) \right], \quad (15)
\end{aligned}$$

$$\begin{aligned}
\frac{dY_{\chi_2}}{dz} = & -\frac{s}{\mathcal{H}z} \left[\langle \sigma v \rangle_{\chi_2 \bar{\chi}_2 \rightarrow \zeta \bar{\zeta}} \left(Y_{\chi_2}^2 - (Y_{\chi_2}^{\text{eq}})^2 \right) + \langle \sigma v \rangle_{\chi_2 \chi_1 \rightarrow \zeta \bar{\zeta}} \left(Y_{\chi_2} Y_{\chi_1} - Y_{\chi_2}^{\text{eq}} Y_{\chi_1}^{\text{eq}} \right) + \langle \sigma v \rangle_{\chi_2 \bar{\zeta} \rightarrow \chi_1 \zeta} \left(Y_{\chi_2} Y_{\zeta}^{\text{eq}} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_{\zeta}^{\text{eq}} \right) \right. \\
& + \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_2}^2 - \frac{(Y_{\chi_2}^{\text{eq}})^2}{(Y_{\chi_1}^{\text{eq}})^2} Y_{\chi_1}^2 \right) + \langle \sigma v \rangle_{\chi_1 \chi_2 \rightarrow \chi_1 \chi_1} \left(Y_{\chi_1} Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1}^2 \right) + \langle \sigma v \rangle_{\chi_2 \chi_2 \rightarrow \chi_1 \chi_2} \left(Y_{\chi_2}^2 - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} Y_{\chi_2} \right) \\
& \left. + \frac{\Gamma_{\chi_2 \rightarrow \chi_1 f \bar{f}}}{s} \left(Y_{\chi_2} - \frac{Y_{\chi_2}^{\text{eq}}}{Y_{\chi_1}^{\text{eq}}} Y_{\chi_1} \right) \right], \quad (16)
\end{aligned}$$

where the various parameters are consistent with those in Eq. (4). The thermal average cross section $\langle \sigma v \rangle$ is also calculated using MICROMEGAS [72,73]. Furthermore, the newly identified Y_{ζ}^{eq} can be categorized into $Y_{\zeta}^{\text{eq}} = 0.238$ for $\zeta = f$ and

$$Y_{\zeta}^{\text{eq}} = \frac{45z^2}{4\pi^4 g_s} \left(\frac{m_{Z'}}{m_{\chi_1}} \right)^2 \mathcal{K}_2 \left(\frac{m_{Z'}}{m_{\chi_1}} z \right), \quad (17)$$

when ζ represents the massive Z' .

In Fig. 6, we introduce all relevant dark matter production mechanisms through the benchmarks, specifically cospattering, conversion, and coannihilation. The criteria for evaluating each mechanism are consistent with those outlined in the resonance scenario. It can be observed from Fig. 6 that conversion is located at light m_{χ_1} and small g' . As these two parameters gradually increase, the dominant phase transits toward cospattering and ultimately becomes entirely governed by coannihilation at $m_{\chi_1} \sim \text{TeV}$ and $g' \sim \mathcal{O}(1)$. The overall distribution of these three phases on g' exhibits a monotonic relationship with m_{χ_1} . This is significantly different from the resonance scenario, where the overall monotonic relationship no longer exists. Such correlation is mainly due to the processes that dominate the generation of dark fermions as $\chi_{1,2} \chi_{1,2} \rightarrow Z' Z'$, which are almost unaffected by g' in the secluded scheme. However, the situation is quite the opposite in the resonance scenario.

Additionally, it is noteworthy that we incorporate the cospattering process $\chi_2 Z' \rightarrow \chi_1 Z'$ into Eq. (15), whose contribution is comparable to that of $\chi_2 f \rightarrow \chi_1 f$ when $m_{Z'} \ll m_{\chi_{1,2}}$. We explore the demand for freeze-out within a broad range of $r_{Z'} \in [0.1, 1]$, $g' \in [10^{-10}, 10^{-2}]$, $\Delta_{\chi} \in [10^{-3}, 10^{-1}]$, and $g_{\chi} \in [0.1, 1]$. We report that it is sufficient for θ to be greater than $\mathcal{O}(10^{-8})$ with m_{χ_1} from GeV to TeV scale.

B. Phenomenology of Z'

In the secluded scenario, the constraints applied on Z' are consistent with those in the resonance scenario, encompassing the excluded region by the current experiments BABAR [42,43], LHCb [44,45], LEP [77], CMS, and ATLAS [46,78], and the sensitivities of future Belle II [79,80], FCC-ee [81], high luminosity CMS, and ATLAS [76]. They are depicted in Fig. 7 as shaded gray areas for the current excluded limits as well as orange, pink, and purple dashed lines for the future reaches.

In Fig. 7, we illustrate the impact of Z' induced constraints on various parameters. The mass ratio $r_{Z'}$ varies from 0.3 to 0.95 in Fig. 7(a) with fixed $\theta = 2 \times 10^{-2}$, $g_{\chi} = 1$, and $\Delta_{\chi} = 10^{-2}$. From the perspective of g' , all conversion regimes are distributed at $g' \lesssim 3 \times 10^{-5}$. This is understandable, since among all relevant reactions, the conversion process $\chi_i \chi_2 \rightarrow \chi_j \chi_1$ is not sensitive to the gauge coupling g' . For such small g' , the thermal rate of the cospattering process $\chi_2 f \rightarrow \chi_1 f$ is suppressed, resulting in a contribution that is less than that of the conversion process $\chi_i \chi_2 \rightarrow \chi_j \chi_1$, which is illustrated in Figs. 6(b) and 6(e) for specific details. At larger $g' \in [3 \times 10^{-5}, 1.8 \times 10^{-3}]$, the contribution from $\chi_2 f \rightarrow \chi_1 f$ surpasses that of $\chi_i \chi_2 \rightarrow \chi_j \chi_1$ and becomes dominant, corresponding to the cospattering phase. As g' continues to increase, the thermal rate of $\chi_2 f \rightarrow \chi_1 f$ exceeds those of the other production processes, thus violating the criteria for cospattering. Concurrently, contributions from $\chi_2 f \rightarrow \chi_1 f$ gradually become insufficient as $m_{Z'}$ increases, hence entering into the domain of coannihilation. Moreover, it is evident that when $g' \gtrsim 0.1$, the production process of the $f \bar{f}$ final state mediated by Z' begins to exert significant influence, causing $m_{Z'}$ to once again push toward larger values.

From the viewpoint of $m_{Z'}$, conversion, cospattering, and coannihilation correspond, respectively, at the orders of magnitude around $\mathcal{O}(10)$, $\mathcal{O}(100)$, and $\mathcal{O}(1000)$ GeV.

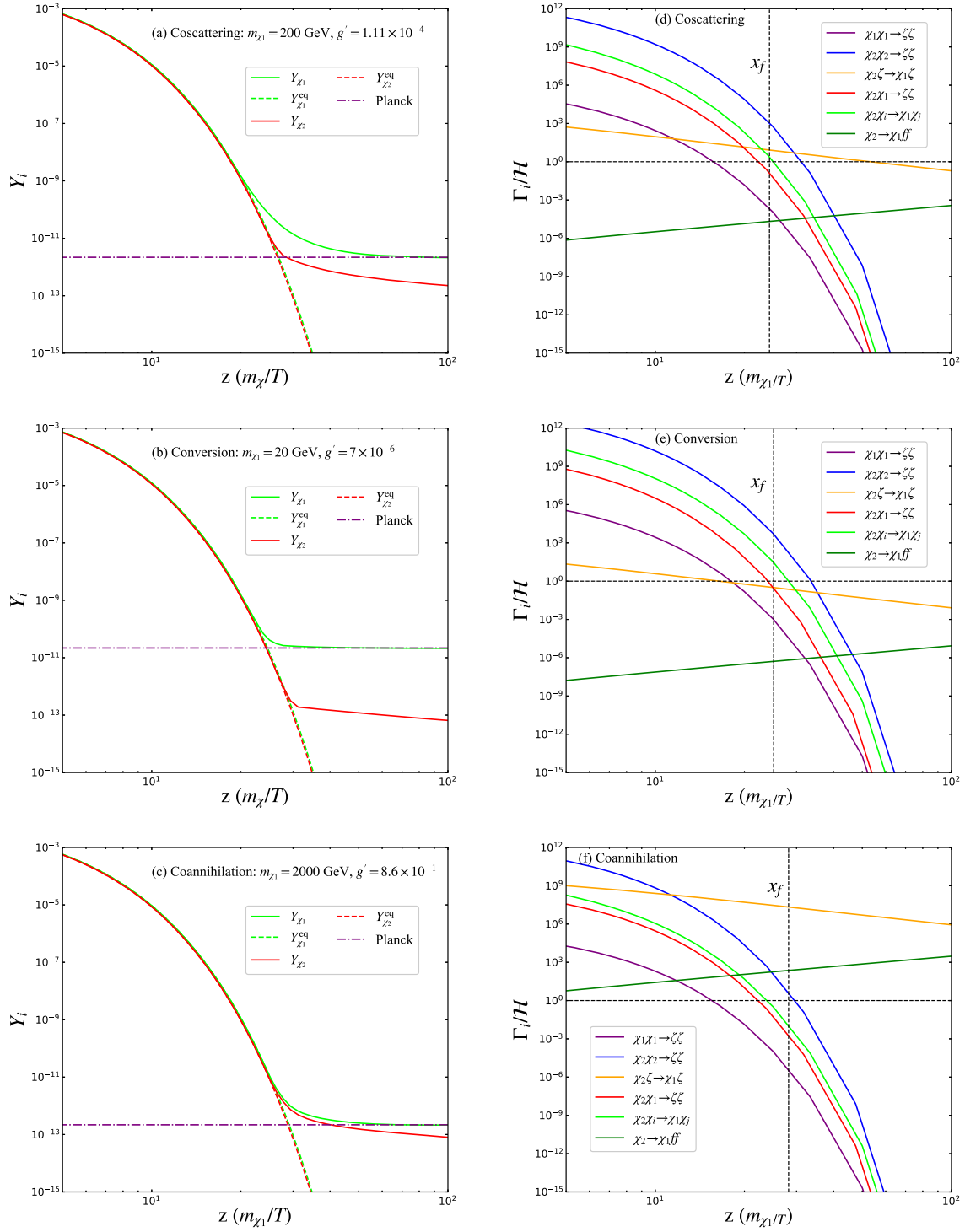


FIG. 6. The evolutions of various abundances Y_i for (a) coscattering, (b) conversion, and (c) coannihilation benchmarks in the secluded scenario. (d)–(f) The thermal rates of various processes in mechanisms described by (a)–(c), respectively. Here we fix $\Delta_\chi = 10^{-2}$, $r_{Z'} = 0.75$, and $\theta = 2 \times 10^{-2}$. The meanings of different colored curves can be referenced in Fig. 1.

Additionally, for smaller $r_{Z'}$, the dominance within the coscattering channel shifts toward $\chi_2 Z' \rightarrow \chi_1 Z'$. This particular pathway is solely influenced by θ and only becomes effective when θ approaches $\mathcal{O}(10^{-6})$. However, such a low

magnitude of θ does not attract the interest of DM-related phenomenology, so this aspect is not considered in the following discussion. Overall, any benchmark starts from the conversion with minimum g' and $m_{Z'}$. The distribution

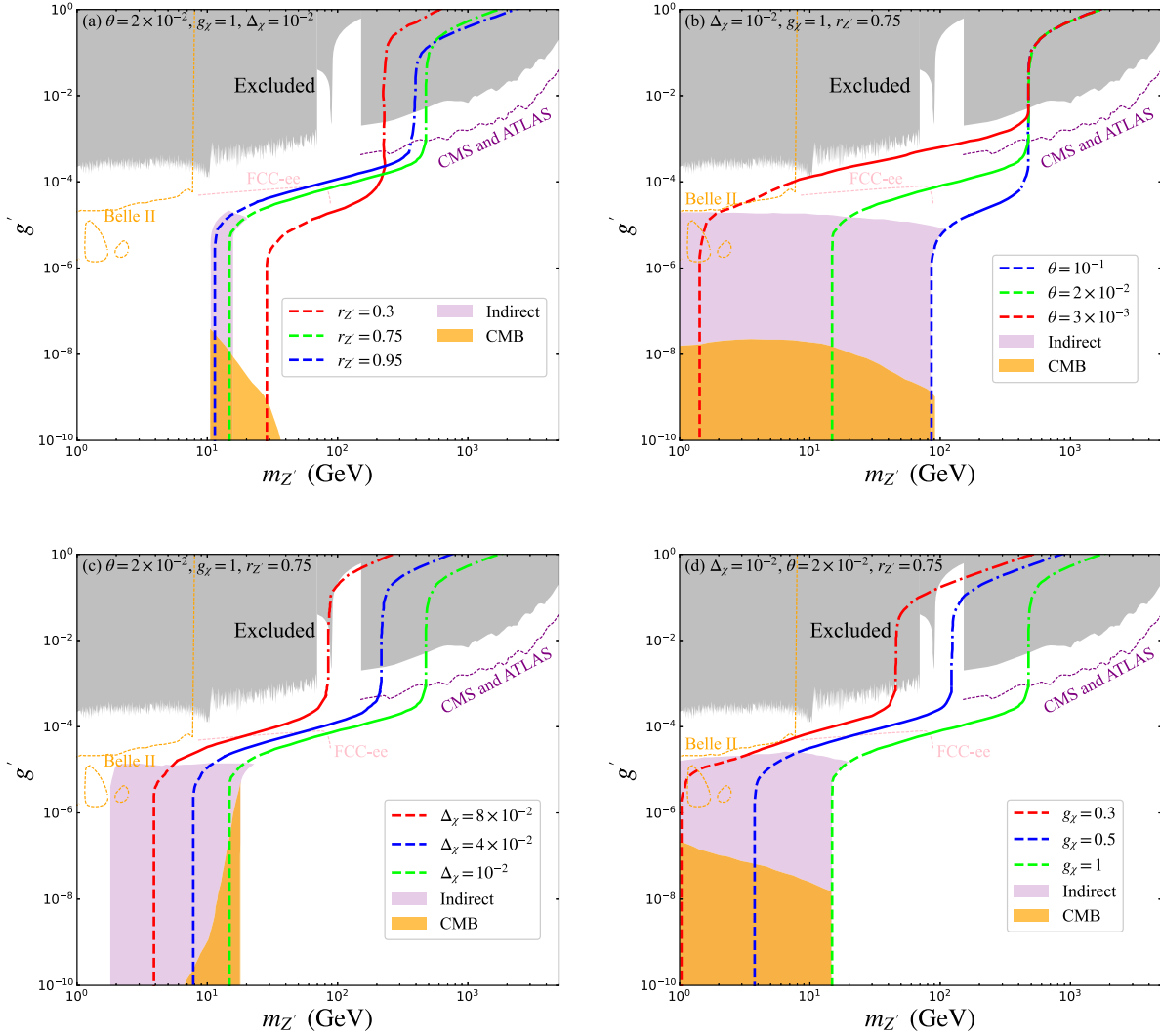


FIG. 7. Constraints on the $m_{Z'}$ – g' parameter space in the secluded scenario. The three fixed parameters and one varying benchmark marked by red, green, and blue lines in (a)–(d) are consistent with those in the resonance scenario of Fig. 2, although some parameters have different numerical values. Similarly, cospattering, conversion, and coannihilation phases are designated as solid, dashed, and dot-dashed lines, respectively. The gray shaded regions are excluded by the current limits on Z' at colliders, while the future sensitivities are indicated by orange, pink, and purple dashed lines.

of conversion is minimally affected by g' , because the key factors $\Gamma_{\chi_2 \bar{\chi}_2 \rightarrow Z' Z'}$, $\Gamma_{\chi_1 \bar{\chi}_2 \rightarrow \chi_j \chi_1}$, and $\Gamma_{\chi_1 \bar{\chi}_1 \rightarrow Z' Z'}$ in the conversion determination criteria are all independent of g' . When g' increases to $\mathcal{O}(10^{-5})$, it triggers a transition from the conversion phase to the cospattering phase because the reaction rate of the cospattering process is proportional to g' . Subsequently, $m_{Z'}$ must increase in synchrony with g' to ensure the judgment condition of cospattering $\Gamma_{\chi_2 \bar{\chi}_2 \rightarrow Z' Z'} > \Gamma_{\chi_2 f \rightarrow \chi_1 f}$. However, since $\Gamma_{\chi_2 f \rightarrow \chi_1 f}$ increases at an exceptionally rapid rate with increasing g' , the criterion governing cospattering breaks down when $g' \sim \mathcal{O}(10^{-3})$. Then it enters the coannihilation phase dominated by $\chi_2 \bar{\chi}_2 \rightarrow Z' Z'$ that is insensitive to g' , so a vertical line like the conversion phase exhibits until the emergence

of the coannihilation phase dominated by $\chi_2 \bar{\chi}_2 \rightarrow f \bar{f}$. Subsequently, similar to the resonance scenario, g' is proportional to $m_{Z'}$. The benchmarks corresponding to different parameter choices discussed below exhibit nearly identical evolutionary trends, differing only in their initial and final $m_{Z'}$. From a phenomenological perspective, under the constraints from Z' , the results in the secluded scenario are markedly different from those obtained in the resonance scenario. The coannihilation distributed in large g' is almost entirely excluded by the current Z' constraints. In the future, the collider experiments might probe the secluded cospattering phase with Z' around the electroweak scale and $r_{Z'} \sim \mathcal{O}(0.1)$.

In Fig. 7(b), we choose to change θ . For the conversion phase, the corresponding $m_{Z'}$ rises from $\mathcal{O}(1)$ GeV to 100 GeV as θ increases, with the upper limit on g'

decreasing from 10^{-4} to 10^{-5} . Meanwhile, the distribution of cospattering continuously shrinks from an initial $m_{Z'} \in [6.9, 480]$ GeV with $g' \in [8.4 \times 10^{-5}, 10^{-2}]$ for $\theta = 3 \times 10^{-3}$ to the final $m_{Z'} \in [172.1, 438.2]$ GeV with $g' \in [1.8 \times 10^{-5}, 9.2 \times 10^{-5}]$ for $\theta = 10^{-1}$. Ultimately, all benchmark lines converge at $m_{Z'} \sim 480$ GeV with the coannihilation phase, primarily due to the fact that the dominant process occurring at $\chi_2 \chi_2 \rightarrow \zeta \bar{\zeta}$ within coannihilation shows little dependence on the mixing θ . The Belle II, which is sensitive to light $m_{Z'}$, shows a preference for conversion and cospattering when $\theta \lesssim 3 \times 10^{-3}$. FCC-ee, CMS, and ATLAS further have the potential to capture cospattering as well as coannihilation that has not yet been excluded when $\theta \gtrsim 3 \times 10^{-3}$.

In Figs. 7(c) and 7(d), it is obvious that as either Δ_χ decreases or g_χ increases, the benchmark curve shifts toward higher $m_{Z'}$. For the three phases of conversion, cospattering, and coannihilation, the corresponding $m_{Z'}$ increase from $\mathcal{O}(1)$ GeV to $\mathcal{O}(10)$ GeV, from $\mathcal{O}(10)$ to $\mathcal{O}(100)$ GeV, and from $\mathcal{O}(100)$ GeV to TeV, respectively. In contrast, the distribution of certain phases on g' is not significantly influenced by parameters Δ_χ or g_χ , such as $g' \lesssim 3 \times 10^{-5}$ for conversion, $3 \times 10^{-5} \lesssim g' \lesssim 1.8 \times 10^{-3}$ for cospattering, and $g' \gtrsim 1.8 \times 10^{-3}$ for coannihilation. In the future, Belle II will be sensitive to conversion and cospattering with $\Delta_\chi \gtrsim 8 \times 10^{-2}$ and $g_\chi \lesssim 0.5$, while FCC-ee, CMS, and ATLAS will favor cospattering with $g_\chi \sim \mathcal{O}(0.1)$ and relatively free Δ_χ , and coannihilation with $\Delta_\chi \lesssim 4 \times 10^{-2}$ and $g_\chi \gtrsim 0.5$. It should be mentioned that the current constraints on Z' are lacking when $m_{Z'} \sim 100$ GeV. This allows a portion of coannihilation to persist at $\Delta_\chi \sim 8 \times 10^{-2}$ and $g_\chi \sim 0.5$, which is highly appealing for the subsequent discussion of DM phenomenology.

In short, the secluded and resonance schemes yield different results for the correct relic density under constraints of Z' . For example, promising conversions appear at the parameter space with $\theta \lesssim 3 \times 10^{-3}$, $\Delta_\chi \gtrsim 8 \times 10^{-2}$, and $g_\chi \lesssim 0.5$ in the secluded scenario; while hopeful cospattering requires $\theta \gtrsim 3 \times 10^{-3}$ and $g_\chi \sim \mathcal{O}(0.1)$, as long as the value of Δ_χ is not excessively large. The vast majority of coannihilation is excluded, but a small portion with $\theta \gtrsim 3 \times 10^{-3}$, $\Delta_\chi \lesssim 4 \times 10^{-2}$, and $g_\chi \gtrsim 0.5$ may potentially be detected by future CMS and ATLAS. In the conversion region with $g' \lesssim 10^{-5}$, it is promising at future indirect and CMB experiments, which will be considered in the following discussion.

C. Phenomenology of χ_1

We first consider the direct detection of dark matter in the secluded scenario, where the constraints on the spin-independent scattering cross section of DM nucleons are identical to those presented in the resonance scenario depicted in Fig. 3. Meanwhile, the numerical result of

the scattering cross section can also be calculated through Eq. (9).

The constraints from direct detection experiments on the secluded scenario are presented in Fig. 8. Here, the results of the benchmarks are obtained by substituting the corresponding ones from Fig. 7 into Eq. (9). Consequently, all curves exhibit nearly identical shapes. However, unlike the shape shown in Fig. 7, the σ_{SI} of the cospattering and final coannihilation components are inversely proportional to m_{χ_1} due to the suppression of $m_{Z'}$. In Fig. 8(a) with varying $r_{Z'}$, conversion exhibits the highest $\sigma_{\text{SI}} \sim 2.2 \times 10^{-49}$ cm² at $m_{\chi_1} \sim 17$ GeV when $r_{Z'} = 0.95$, however, this value remains slightly lower than the result $\sigma_{\text{SI}} \sim 8.3 \times 10^{-50}$ cm² from future direct detection experiments. Cospattering usually predicts $\sigma_{\text{SI}} \lesssim 10^{-50}$ cm², thus it is also beyond the reach of direct detection. Therefore, only coannihilation holds promise for validation by future direct detection. However, such a coannihilation phase is not permitted by the constraints of Z' . In Fig. 8(b), the variations in θ give rise to conversion and cospattering having $\sigma_{\text{SI}} \lesssim 1.2 \times 10^{-49}$ cm². Meanwhile, the coannihilations that fall within the detection range are excluded by Z' constraints. Therefore, similar conclusions can be drawn as those presented in Fig. 8(a). The situations have been improved in Figs. 8(c) and 8(d), where cospattering below 60 GeV can be easily detected by the projected experiments when $\Delta_\chi \gtrsim 4 \times 10^{-2}$ and $g_\chi \lesssim 0.5$. Furthermore, due to the less stringent constraints imposed by Z' on $\Delta_\chi \sim 8 \times 10^{-2}$ and $g_\chi \sim 0.5$, this small portion of coannihilation near 100 GeV also shows considerable promise.

One particularly bright aspect of traditional secluded DM arises from the indirect detection experiments, as the typical secluded annihilation cross section $\langle \sigma v \rangle$ is still around the benchmark WIMP value 2×10^{-26} cm³/s [110]. In the secluded scenario of this model, the primary annihilation process of χ_1 is denoted as $\chi_1 \bar{\chi}_1 \rightarrow Z' Z'$ followed by the decay $Z' \rightarrow f \bar{f}$. Because the direct $\chi_1 \chi_1 Z'$ coupling is heavily suppressed by $\sin^2 \theta$, while the $\chi_1 \chi_2 Z'$ coupling is less suppressed by $\sin 2\theta$, the dominant contribution of secluded annihilation $\chi_1 \bar{\chi}_1 \rightarrow Z' Z'$ is mediated by the dark partner χ_2 through the t/u channel. The current thermal average cross section with nearly degenerate masses of dark fermions for a small mixing angle θ can be approximated as [48]

$$\langle \sigma v \rangle_{\chi_1 \bar{\chi}_1 \rightarrow Z' Z'} \simeq \frac{\sin^4 2\theta g_\chi^4}{256\pi m_{\chi_1}^2} \left(1 - \frac{m_{Z'}^2}{m_{\chi_1}^2}\right)^{3/2} \left(1 - \frac{m_{Z'}^2}{2m_{\chi_1}^2}\right)^{-2}. \quad (18)$$

Obviously, compared to the traditional secluded DM, the secluded annihilation cross section of this model is suppressed by the mixing angle θ . Therefore, the predicted value is expected to be much smaller than the benchmark WIMP value 2×10^{-26} cm³/s. The experimental limits from indirect detection on secluded DM also depend on the

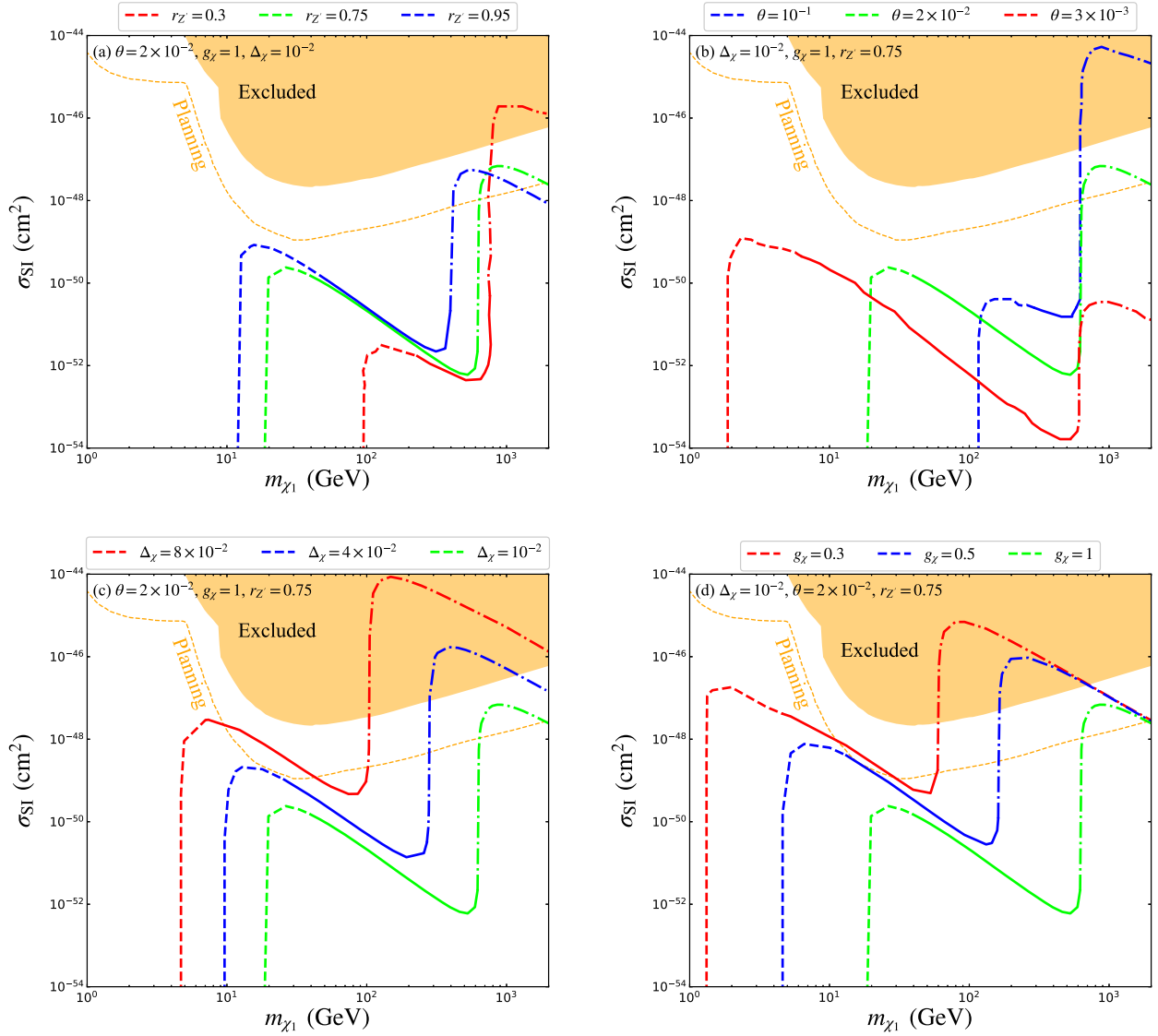


FIG. 8. The constraints of direct detection experiments in the secluded scenario. Legends and markers occurring (a)–(d) correspond directly to those presented in Fig. 7. The current and future constraints are represented by the orange shaded region and dashed line, respectively.

final states [111]. In this paper, we consider the electron final state $\chi_1 \bar{\chi}_1 \rightarrow Z' Z' \rightarrow 4e$ for illustration, which is presented in Fig. 9 accompanied by the predictions of $\langle \sigma v \rangle$ for various benchmarks. Here, the current limits are the joined results from Fermi-LAT dwarfs, CMB, and H.E.S.S. GC observations [111].

Limited by relic density, we report that lighter DM can be obtained with larger mass ratio $r_{Z'}$, smaller mixing angle θ , smaller mass splitting ratio Δ_χ , and smaller dark coupling g_χ . In Fig. 9, all benchmark lines exhibit nearly identical shapes. In the conversion and cospattering regions, the coupling g' is relatively small and the total $\langle \sigma v \rangle$ is well approximated by Eq. (18) and is inversely proportional to m_{χ_1} . However, in the final coannihilation segment, the influence of $\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}$ must be taken into account, as $\langle \sigma v \rangle_{\chi_1 \bar{\chi}_1 \rightarrow f \bar{f}}$ could be larger than $\langle \sigma v \rangle_{\chi_1 \bar{\chi}_1 \rightarrow Z' Z'}$.

This part asymptotically approaches a constant value due to the constraint imposed by observed relic density. In Fig. 9(a), the benchmark lines are cut off at the minimum m_{χ_1} , i.e., conversion regime, since the secluded process $\chi_1 \bar{\chi}_1 \rightarrow Z' Z'$ remains unaffected by g' . When $r_{Z'} \gtrsim 0.75$, conversion exhibits a relatively large annihilation cross section $\langle \sigma v \rangle \sim \mathcal{O}(10^{-29}) \text{ cm}^3/\text{s}$ below 30 GeV, naturally falling within the detection range of future experiments. However, the $\langle \sigma v \rangle$ for cospattering and coannihilation is suppressed by heavier m_{χ_1} , with most values dropping to a pessimistic level below $\mathcal{O}(10^{-31}) \text{ cm}^3/\text{s}$. In Figs. 9(b)–9(d), all benchmark lines possess small enough $\langle \sigma v \rangle$ that are not excluded by the current constraints, and the most promising one is still conversion. Depending on different parameters, the corresponding m_{χ_1} ranges from $\mathcal{O}(1)$ to $\mathcal{O}(100)$ GeV. Certainly, as observed in Figs. 9(c) and 9(d), when the Δ_χ is

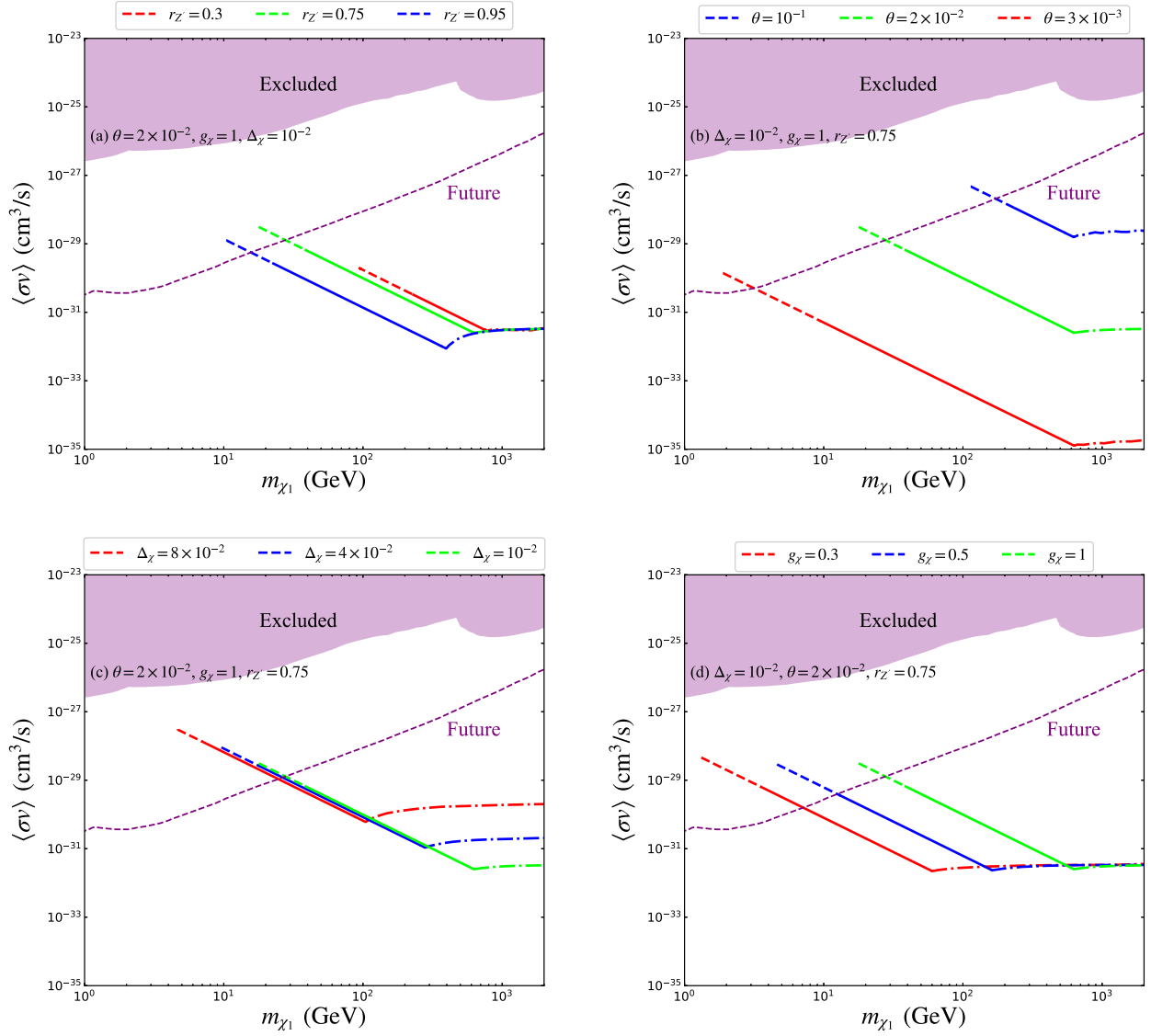


FIG. 9. The constraints of indirect detection experiments in the secluded scenario. The benchmarks selected in panels (a)–(d) are consistent with those presented in Fig. 7. The purple shadow region and dashed line represent the existing constraints [111] as well as the sensitivities of future experiments [102], respectively.

relatively large or g_χ is relatively smaller, e.g., $\Delta_\chi \gtrsim 0.04$ or $g_\chi \lesssim 0.5$, the cospattering below about 20 GeV will also be subjected to examination.

On the whole, for $m_{\chi_1} \lesssim \mathcal{O}(100)$ GeV, future direct detection experiments are more inclined toward cospattering and heavily rely on Δ_χ and g_χ , requiring them to be greater than 4×10^{-2} and less than 0.5, respectively. In terms of indirect detection, it is clear that conversion demonstrates superior performance, with the necessary condition $m_{\chi_{1,2}} \sim m_{Z'}$.

D. Phenomenology of χ_2

In the secluded scenario $m_{Z'} < m_{\chi_2}$, the dark fermion χ_2 is always generated via off-shell Z' at colliders, which naturally

suppresses the production cross section. So the collider signature will not be considered in this scenario. Consequently, the more promising phenomenology induced by the dark partner χ_2 still arises from cosmological aspects, specifically the constraints imposed by BBN [106] and CMB [107] on the delayed decay of $\chi_2 \rightarrow \chi_1 e^+ e^-$. Same as the resonance scenario, they are represented in Fig. 10 by the gray region and orange dashed lines, respectively.

In this scenario, we note that the conversion becomes independent of g' at the minimum m_{χ_1} , which exceeds 1 GeV in all benchmarks. Hence, there is no need for truncation as required in the resonance scenarios. With g' decreasing, the lifetime τ_{χ_2} can become significantly long. As illustrated in Fig. 10, since τ_{χ_2} is primarily governed by

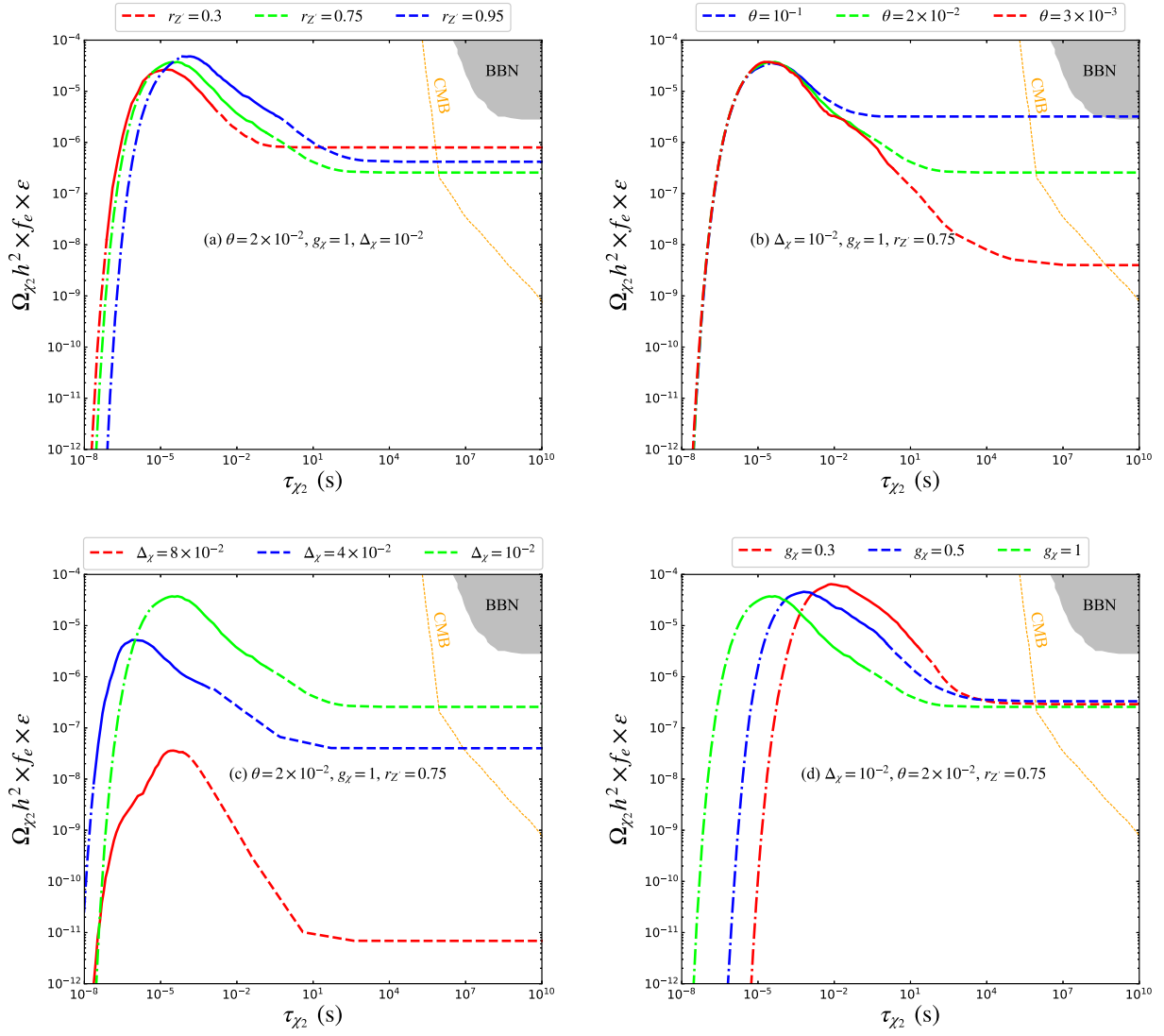


FIG. 10. Cosmological constraints on long-lived χ_2 in the secluded scenario. The parameters appearing in the coordinate axes are defined identically to those in Fig. 5. The benchmarks present in (a)–(d) correspond one-to-one with those shown in Fig. 7. The gray shadow and orange dashed line represent the region excluded by the current BBN constraints and the future CMB sensitivity, respectively.

g' , the increasing τ_{χ_2} in any benchmark line corresponds sequentially to the coannihilation, cospin, and conversion phase. The main factor inducing variation in the longitudinal axis is $\Omega_{\chi_2} h^2$. In the coannihilation phase, $\Omega_{\chi_2} h^2$ is predominantly suppressed by the $\chi_2 f \rightarrow \chi_1 f$ process. The reduction in $\Gamma_{\chi_2 f \rightarrow \chi_1 f}$ alleviates its suppression on $\Omega_{\chi_2} h^2$ as g' decreases (τ_{χ_2} increases). Consequently, $\Omega_{\chi_2} h^2$ begins to increase numerically. When the benchmark enters the cospin phase, the reduction in $m_{Z'}$ enhances $\Gamma_{\chi_1 \chi_2 \rightarrow \chi_1 \chi_1}$, leading to progressively stronger suppression of $\Omega_{\chi_2} h^2$, hence $\Omega_{\chi_2} h^2$ starts to decrease. Finally, in the conversion phase, $\Omega_{\chi_2} h^2$ remains a constant as the dominant $\Gamma_{\chi_1 \chi_2 \rightarrow \chi_1 \chi_1}$ stabilizes and no longer varies. In the phenomenological aspect, the maximum τ_{χ_2} of all benchmarks can

reach the range where the constraints of BBN and CMB are effective. However, different parameter choices result in varying relative relic densities on the vertical axis. Basically speaking, as long as $\theta \lesssim 0.1$, there will be no conversion excluded by BBN. For the conversion that has the potential to be captured by future CMB at $\tau_{\chi_2} \gtrsim \mathcal{O}(10^6)$ s, it is necessary for $\theta \gtrsim 3 \times 10^{-3}$ and $\Delta_\chi < 0.08$.

Furthermore, regarding the CMB constraint on neutrinos produced by the delayed decay of $\chi_2 \rightarrow \chi_1 \nu \bar{\nu}$, we consider the maximum τ_{χ_2} during the CMB epoch, i.e., the calculated result $(f_\nu e \Omega_{\chi_2} / \Omega_{\chi_1})^2 \tau_{\chi_2} \lesssim \mathcal{O}(10)$ s, which is 8 orders of magnitude lower than the current limit [108]. Therefore, constraints from neutrino final states do not pose any threat to our benchmarks.

E. Combination and discussion

Compared to the resonance scenario, the secluded scenario is more promising to be tested at future experiments. Especially in the conversion phase, the relatively large secluded annihilation cross section and long lifetime of the dark partner make this regime detectable. In Fig. 7, we also present the combined results of future indirect detection and CMB that exceed the detection range of colliders. Meanwhile, the sensitive cospattering of future direct detection experiments typically requires $g' \sim \mathcal{O}(10^{-4})$ for $m_{Z'} \sim \mathcal{O}(10)$ GeV, which may also be captured by upcoming colliders.

In Fig. 7(a), we consider $r_{Z'}$ to be greater than 0.2 to ensure the effectiveness of conversion. The future CMB is sensitive to conversion with $10 \lesssim m_{Z'} \lesssim 40$ GeV and $g' \lesssim \mathcal{O}(10^{-8})$. The indirect detection experiments can capture conversion of g' up to $\mathcal{O}(10^{-5})$, however, the sensitive region concentrates on $m_{Z'}$ being a dozen GeV with $r_{Z'} \gtrsim 0.75$ of the benchmarks. It should be mentioned that, in the secluded scenario, the relic density, in principle, does not depend on the coupling g' , which means it can be arbitrarily small. However, for a too small $g' < 10^{-11}$, the lifetime of new gauge boson $\tau_{Z'}$ becomes too large when $m_{Z'} > 1$ GeV, which may alert BBN. To avoid such an issue, we simply require $g' > 10^{-10}$ in this paper.

In Fig. 7(b), an excessively large θ will result in χ_1 being entirely produced through the coannihilation mechanism and the current BBN constraint imposes stringent restrictions on $\theta \gtrsim 0.1$ as shown in Fig. 10. Therefore, we consider θ to be less than 0.1. Correspondingly, conversion will be detected by CMB when $m_{Z'} \lesssim 90$ GeV and $g' \lesssim 2 \times 10^{-8}$. The sensitive region of indirect detection experiments will expand upward by 3 orders of magnitude of g' .

In the case of Fig. 7(c), we also assume that Δ_χ is within the range of $[10^{-3}, 10^{-1}]$. For varying Δ_χ , the future CMB can only capture conversion in the vicinity of 10 GeV. It is clear that, when Δ_χ decreases, the promising region of CMB quickly increases up to $\mathcal{O}(10^{-6})$. The capability for indirect detection is stronger, and it is desirable to detect conversion within the range of 1.8–18 GeV. Sub-GeV DM is not favored by the current constraints from indirect detection, and the corresponding Δ_χ should be less than $\mathcal{O}(0.1)$ of the benchmarks.

In Fig. 7(d) where $g_\chi \lesssim 1$, the future CMB shows sensitivity of $m_{Z'}$ below 15 GeV and g' smaller than 10^{-7} . We also report that the sensitive limit of g' from CMB decreases as g_χ increases. Similarly, the conversion that is expected to be probed by indirect detection experiments could still be expanded to $g' \sim \mathcal{O}(10^{-5})$.

As a whole for the secluded scenario, the majority of coannihilation is excluded by the current collider constraints, and certain cospattering within the range of $\mathcal{O}(10^{-5}) \lesssim g' \lesssim \mathcal{O}(10^{-3})$ could be tested by both future

colliders and direct detection experiments. Meanwhile, portions of conversion in the range of $\mathcal{O}(10^{-8}) \lesssim g' \lesssim \mathcal{O}(10^{-5})$ and $m_\chi \lesssim \mathcal{O}(10)$ GeV may be captured through indirect detection experiments. Furthermore, those conversions below $\mathcal{O}(10^{-8})$ hold promise for dual verification via both CMB and indirect detection experiments. Such a distinct feature makes the secluded scenario of this model quite different from the canonical secluded DM because the canonical GeV-scale secluded DM is tightly constrained by indirect detection and does not predict an observable CMB signature.

V. CONCLUSION

In addition to the pair annihilation of DM, alternative mechanisms, such as cospattering, conversion, and coannihilation, could also significantly modify the evolution of DM relic density. In order to study these new mechanisms within the framework of $U(1)_{B-L}$ symmetry, we introduce two Dirac fermions $\tilde{\chi}_1$ and $\tilde{\chi}_2$ with nearly degenerate masses. Among them, only the nondark matter fermion $\tilde{\chi}_2$ possesses a nonzero $U(1)_{B-L}$ charge, while the stability of DM is ensured by a Z_2 symmetry. The mass term $\delta m_{\tilde{\chi}_1 \tilde{\chi}_2}$ induces mixing between the dark fermions, which is crucial for activating the conversion mechanism. Based on the magnitude of the mass of mediator Z' and dark fermions, we categorize our study into the resonance and secluded scenarios.

In the resonance scenario, the gauge coupling g' in the cospattering phase is typically greater than $\mathcal{O}(10^{-2})$, which is easily excluded by the current Z' constraints. Therefore, we focus on the coannihilation and conversion with $g' \lesssim \mathcal{O}(10^{-3})$. For the conversion at the range of $\mathcal{O}(10^{-5}) \lesssim g' \lesssim \mathcal{O}(10^{-3})$, which is most likely to be probed by future colliders, the mass can vary from GeV to TeV scale, and require $\theta \lesssim 5 \times 10^{-3}$ and $g_\chi \sim \mathcal{O}(0.1)$. On the other hand, the magnitude of Δ_χ should not be excessively large to ensure the effectiveness of conversion. Regarding the coannihilation dominated region with g' between $\mathcal{O}(10^{-8})$ and $\mathcal{O}(10^{-5})$, although detection of Z' at the GeV scale may be achievable with future colliders, more promising signatures are anticipated from forthcoming CMB cosmological observables derived from the long-lived χ_2 . The CMB could capture the dark partner up to TeV, but it is necessary to satisfy $\theta \lesssim 5 \times 10^{-4}$, $\Delta_\chi \lesssim 10^{-2}$, and $g_\chi \lesssim 1$. Because of the suppression of the small mixing angle θ , future direct and indirect detection experiments for DM will face significant challenges in capturing any mechanisms. Unless θ far exceeds 5×10^{-3} , only coannihilation has hope at this moment.

In the secluded scenario with $m_{Z'} \lesssim m_{\chi_{1,2}}$, the situation is quite different from the resonance scenario. The coannihilation with the relatively large g' is naturally excluded by Z' constraints, while a relatively small g' is required by the cospattering and conversion that we are interested in.

For cospattering with $\mathcal{O}(10^{-5}) \lesssim g' \lesssim \mathcal{O}(10^{-3})$, future collider searches for Z' will be sensitive to $\theta \gtrsim 3 \times 10^{-3}$ and $g_\chi \sim \mathcal{O}(0.1)$, with corresponding new gauge boson mass not exceeding the TeV scale. When the DM mass is around dozens of GeV, the future direct detection experiments hold greater promise to probe the region with $\Delta_\chi \gtrsim 4 \times 10^{-2}$ and $g_\chi \lesssim 0.5$. Although the conversion below $\mathcal{O}(100)$ GeV with $g' \lesssim \mathcal{O}(10^{-5})$ is beyond the reach of future colliders, the advantages in indirect detection are even more pronounced. With minimal specific requirements for parameters, namely, $\Delta_\chi \lesssim \mathcal{O}(0.1)$, the conversion can easily satisfy current constraints. Moreover, for the lower regions of $g' \lesssim \mathcal{O}(10^{-8})$, CMB will also participate in the verification of conversion, which is sensitive to $\mathcal{O}(10^{-3}) \lesssim \theta \lesssim \mathcal{O}(0.1)$ and $\mathcal{O}(10^{-3}) \lesssim \Delta_\chi < \mathcal{O}(10^{-2})$.

In summary, although both scenarios can generate DM through the cospattering, conversion, and coannihilation mechanisms, there are significant differences in their phenomenological constraints. In the resonance scenario, cospattering is notably less favored, whereas in the secluded scenario, it changes into coannihilation. Furthermore, the magnitude of θ in the secluded scenario is 2

orders of magnitude larger than that in the resonance scenario, which makes it easier for the secluded case to be tested by future direct and indirect detection experiments. Finally, CMB can test coannihilation in the resonance scheme, but only conversion holds promise in the secluded scenario.

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DATA AVAILABILITY

The data that support the findings of this article are not openly available because no data was created or analyzed in this study.

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