

Heavy dark photon handbook: Cosmological and astrophysical bounds

Andrea Caputo^{1,2,3,*}, Jaeyoung Park^{4,1,†} and Seokhoon Yun^{5,6,‡}

¹Theoretical Physics Department, CERN, 1211 Geneva 23, Switzerland

²Dipartimento di Fisica, “Sapienza” Università di Roma and Sezione INFN Roma1,
Piazzale Aldo Moro 5, Roma 00185, Italy

³Department of Particle Physics and Astrophysics, Weizmann Institute of Science, Rehovot 7610001, Israel

⁴Department of Physics and Astronomy and Center for Theoretical Physics, Seoul National University,
Seoul 08826, Korea

⁵Department of Physics, Kyungpook National University, Daegu 41566, Korea

⁶Center for Theoretical Physics of the Universe, Institute for Basic Science (IBS), Daejeon 34126, Korea



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We investigate cosmological and astrophysical constraints on dark photons with masses $\sim 10^{-1}$ – 10^3 MeV. These dark photons can be copiously produced either in the early Universe or during core-collapse supernovae, potentially leaving distinct observational signatures. First, we derive updated constraints from cosmological and astrophysical observables that rely on the thermal relic abundance of dark photons, including the cosmic microwave background spectrum, primordial light element abundances, and Galactic/extragalactic γ -ray flux. We consider the minimal reheating temperature possible, $T_{\text{RH}} = 6$ MeV, such that our constraints are conservative, but unavoidable within the minimal dark photon model. Then, for supernova-sourced dark photons, we systematically examine all relevant observational bounds, revisit the standard cooling argument and derive limits from other arguments such as fireball formation, low energy supernovae and Galactic positron injection.

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I. INTRODUCTION

New light and weakly coupled degrees of freedom appear ubiquitously in many well-motivated extensions of the Standard Model (SM). These hypothetical particles offer a fertile ground for probing physics beyond the SM through cosmological and astrophysical observations.

In this work, we explore one of the simplest and most theoretically motivated extensions of the SM: a new spin-1 vector field identified as the gauge boson of an additional Abelian $U(1)'$ gauge symmetry. We denote the corresponding quantum field by A'_{μ} , and refer to its particle excitation as the dark photon (DP), denoted by γ' . More specifically, we investigate the constraints and prospects on the most minimal communication channel with the visible sector in the absence of other new particles at the low-energy scale probed by processes relevant to cosmology and astrophysics: the renormalizable Abelian kinetic mixing [1–5].

The DP can be probed through a wide variety of experimental and observational means [6]. Laboratory searches, including fixed-target [7–12] and beam dump experiments [13,14], rare meson decay measurements [8,15–25], and collider signatures [26–33], have excluded wide regions of the DP parameter space, particularly at higher kinetic mixing, where DPs are efficiently produced and decay promptly. However, at lower values of the kinetic mixing, where the DP becomes long-lived on experimental timescales, terrestrial experiments lose sensitivity, and the role of astrophysical and cosmological probes becomes crucial.

In these regimes, DPs can be efficiently produced in extreme environments, such as the hot and dense interiors of stars, supernovae (SNe), and neutron star mergers. Depending on its mass and mixing, DPs may escape the object, deposit energy outside via decay, or remain trapped and contribute to local energy transport, thereby modifying observable quantities such as neutrino emission, light curves, or postmerger signals.

In parallel, the cosmological implications of MeV-scale DPs offer complementary and powerful avenues of constraint. In the early Universe, DPs produced through thermal processes can contribute to the expansion history. Their radiative decays may inject entropy, alter the predicted light element abundances of big bang nucleosynthesis (BBN), or contribute to the effective number of

* Contact author: andrea.caputo@cern.ch

† Contact author: rs1958@snu.ac.kr

‡ Contact author: seokhoon.yun@knu.ac.kr

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relativistic species and distort the cosmic microwave background (CMB), all of which are precisely measured by current observations. At later times, decays of relic DPs into photons or electron-positron pairs can inject energy into the intergalactic medium, and be probed through cosmological observables such as the CMB anisotropy spectrum by affecting the ionization history. In Galactic and extragalactic environments, the decay products can lead to distinctive features in x-ray signals if DPs are long-lived and decay on astrophysical scales. Constraints from the Milky Way, clusters, and nearby galaxies provide stringent bounds in this regime.

In this work, we provide a comprehensive survey of the phenomenology of MeV-scale DPs, bringing together a wide range of observational and experimental constraints (see Ref. [34] instead for smaller DP masses). Our analysis includes cosmological limits from early-Universe physics, and astrophysical probes from SNe-associated phenomena. In particular, we highlight the role of recent developments in supernova probes and neutron star merger simulations for probing long-lived DPs. We also collect updated colliders constraints in the regime of interest.

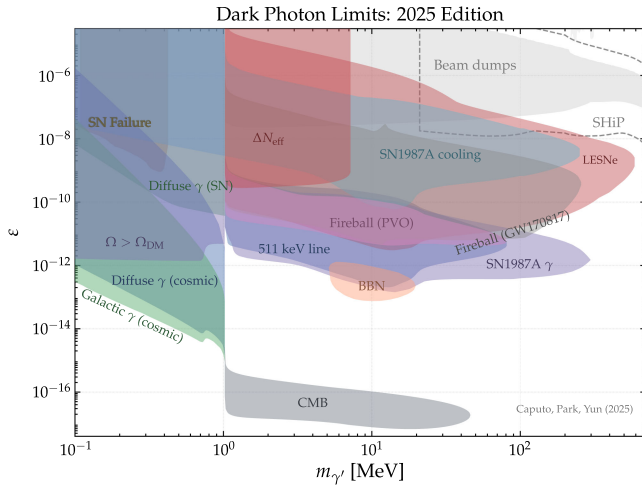


FIG. 1. Astrophysical, cosmological, and colliders constraints on the dark photon model in MeV mass range. In particular, we show constraints from SN1987A cooling (dark blue), failed supernovae (mustard), low-energy SNe (LESNe, light red), BBN (salmon), ΔN_{eff} (red), relic-density $\Omega_{\gamma'}$ (olive), 511 keV line (violet blue), diffuse supernova γ -rays, SN1987A γ -ray constraints (lilac), Galactic and diffuse γ -ray limits from cosmic relic densities (green), CMB anisotropies (dark gray), fireball formation both from both SN1987A (magenta) and GW170817 (gray). The light gray region labeled “Beam dumps” shows the combined exclusion from fixed-target and collider experiments, obtained using DarkCast [35–37] and Ref. [38]. We stress that the combination of low-energy supernovae, fireball formation, supernova γ -ray signals, and Galactic positron injection is independent of the supernova explosion mechanism and, taken together, supersedes the standard cooling argument. See the text for the details of each constraint.

By combining all available constraints, we present an up-to-date map of the viable parameter space for MeV-scale DPs (see Fig. 1). This work serves as both a status report and a resource for future investigations, identifying open windows for discovery and clarifying the complementarity of different probes across cosmic and terrestrial frontiers.

II. DARK PHOTON THEORY, INTERACTIONS, AND PRODUCTION RATES

This section defines the theoretical framework that forms the basis of our phenomenological analysis. We begin by introducing the low-energy effective DP interactions and discussing plausible UV origins for the two key parameters appearing on the Cartesian axes of Fig. 1. We then examine how Abelian kinetic mixing induces interactions between the DP and electrically charged SM particles. This requires identifying the physical propagating states, which we carry out both in vacuum and in a medium. Finally, we use these results to derive the master formula in Eq. (5), which quantifies the number of DPs produced per unit time and volume. This expression will be repeatedly used in the following sections to study DP production in both the early Universe and supernova environments.

The low-energy Lagrangian for the DP that captures the relevant phenomenology reads as follows:

$$\mathcal{L}_{\text{DP}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - \frac{1}{4}F'_{\mu\nu}F'^{\mu\nu} + \frac{\epsilon}{2}F_{\mu\nu}F'^{\mu\nu} + \frac{m_{\gamma'}}{2}A'_\mu A'^\mu + eA_\mu J_{\text{EM}}^\mu. \quad (1)$$

Here, $F_{\mu\nu}^{(\prime)} = \partial_\mu A_\nu^{(\prime)} - \partial_\nu A_\mu^{(\prime)}$ is the photon (DP) field strength tensor, $m_{\gamma'}$ is the DP mass, and J_{EM}^μ is the electromagnetic (EM) current. We notice how this Lagrangian contains the two parameters that define our framework: the kinetic mixing ϵ and the DP mass $m_{\gamma'}$.

The first row of Eq. (1) contains the kinetic terms of the Lagrangian, including the operator that serves as the sole interaction channel between the visible and dark sectors: the kinetic mixing between the two Abelian field strengths.¹ This interaction is conventionally normalized by the dimensionless parameter ϵ . Because particles charged under both Abelian symmetries induce radiative corrections, the kinetic mixing parameter typically runs with the energy scale [1]. In any given UV complete theory, its value must be fixed at a specific energy scale, with the

¹If we were interested in a Lagrangian valid at arbitrarily high energies, the Abelian gauge group to which the DP couples would be the hypercharge $U(1)_Y$ [1]. However, since the processes relevant to our analysis take place at energies well below the Fermi scale, we work with the SM in its broken electroweak phase and accordingly couple the DP to the electromagnetic gauge group; coupling to the neutral weak gauge group is suppressed by the Fermi scale as $(m_{\gamma'}/m_Z)^2$ [39–41].

low-energy value determined accordingly. In this work, we adopt a practical approach and treat ε as the kinetic mixing parameter evaluated at the low energies relevant for the constraints used in our analysis.

The mass term for the DP in the second row of Eq. (1) explicitly breaks the $U(1)'$ gauge symmetry. However, there are several mechanisms that can generate such a mass at low energies while preserving gauge invariance at high energies. Two widely studied examples are the dark Higgs mechanism [42] and the Stueckelberg mechanism [43]. The former, analogous to the well-established mechanism for electroweak symmetry breaking, introduces a new scalar field that spontaneously breaks the Abelian symmetry, resulting in a massive gauge boson and an additional radial mode in the spectrum. In contrast, the Stueckelberg mechanism generates a gauge-invariant mass term without spontaneous symmetry breaking or additional physical scalars, by coupling the gauge field to a pseudoscalar degree of freedom. In this work, we remain agnostic about the origin of the DP mass. We assume that any additional states associated with the mass-generation mechanism are heavy enough to be phenomenologically irrelevant, and we treat $m_{\gamma'}$ as a free parameter throughout our analysis.

Having introduced the effective parameters that define our framework and outlined our low-energy, pragmatic approach, we now turn to the physical implications of a nonzero kinetic mixing. As we will show below, this term induces a small effective coupling between the DP and electrically charged SM particles.

We start our discussion from the simplest case where physical processes happen in the vacuum and the SM photon field is massless. The diagonalization of the photon and DP fields in Eq. (1) can be achieved via the gauge-invariant shift of $A_\mu \rightarrow A_\mu + \varepsilon A'_\mu$ at the leading order of ε (higher-order terms in ε arise from the canonical field normalization) [44]. Therefore, the interactions of DPs with the SM contents in the vacuum are simply written as [44]

$$\mathcal{L}_{\text{DP}|_{\text{vacuum}}} \supset \varepsilon e A'_\mu J_{\text{EM}}^\mu. \quad (2)$$

The assumption that we are in vacuum is not always justified and, in some of the cases considered in this work, it is not valid. For this reason, we are also forced to consider situations in which the system is embedded in a dense medium. The effective DP couplings in this case are affected by plasma effects, which modifies photon dispersion. Considering the thermally corrected photon propagator in the form of the polarization tensor, these couplings can be described by the replacement of the kinetic mixing as [45,46]

$$\varepsilon \rightarrow \varepsilon \frac{m_{\gamma'}^2}{m_{\gamma'}^2 - \pi_{\text{T,L}}}, \quad (3)$$

where $\pi_{\text{T,L}}$ quantifies the refractive properties of the intermediate photon propagator for transverse (T) and longitudinal (L) polarizations, respectively. Explicit expressions for the polarizations can be found in Appendix A. This expression indeed covers the case of the vacuum, where both polarizations $\pi_{\text{T,L}}$ vanish.

We conclude this section with a very useful expression that quantifies the DP production rate. The DP coupling structure, which is provided in Eq. (3), follows that of photons (i.e., EM interactions) up to an overall factor proportional to the kinetic mixing ε . Thus, we can express any interaction rates of DPs in terms of the total absorption rate of SM photons. To this purpose, it is useful to appreciate how the imaginary components of $\pi_{\text{T,L}}$ rely on processes that bring photons into thermal equilibrium through scatterings, such as absorption and/or decay. When particles participating in such processes are in thermal equilibrium, detailed balance yields [47–49]

$$\text{Im } \pi_{\text{T,L}} = -\omega(1 - e^{-\omega/T})\Gamma_{\text{T,L}}. \quad (4)$$

Here, $\Gamma_{\text{T,L}}$ represents the total absorption rate of photons with a modified dispersion relation that corresponds to an external DP with energy ω and momentum $\vec{k}$, or in other words satisfying the relation $\omega^2 - k^2 = m_{\gamma'}^2$. Therefore, all the functions appearing in Eq. (4) are determined by properties of the QED plasma populating the early Universe.

In particular, the differential DP production rate per volume can be written as follows [45,49]:

$$\begin{aligned} \frac{d\mathcal{N}_{\gamma'}}{dV dt d^3\vec{k}} &= \frac{\varepsilon^2}{(2\pi)^3} \sum_{i=\text{T,L}} g_i e^{-\omega/T} \Gamma_i \left| \frac{m_{\gamma'}^2}{m_{\gamma'}^2 - \pi_i} \right|^2 \\ &= \frac{1}{(2\pi)^3} \frac{\varepsilon^2 m_{\gamma'}^4}{\omega} \frac{1}{e^{\omega/T} - 1} \\ &\quad \times \sum_{i=\text{T,L}} g_i \frac{|\text{Im } \pi_i|}{(m_{\gamma'}^2 - \text{Re } \pi_i)^2 + |\text{Im } \pi_i|^2}, \end{aligned} \quad (5)$$

with $g_i = 2(1)$ for the transverse (longitudinal) polarization. Under the typical condition of a narrow resonance width (i.e., $\text{Re } \pi_i \gg \text{Im } \pi_i$), the DP production during the resonance can be interpreted as the resonant conversion from background photons and becomes independent of the details of scattering processes, as we will discuss.

III. COSMOLOGICAL DARK PHOTONS: THERMAL PRODUCTION AND SIGNALS

The hot plasma in the early Universe serves as a natural source for the production of massive DPs. In this paper, we focus on the primordial DP abundance with masses around the MeV scale, generated through thermal processes such as particle scatterings within the thermal bath. Even if these thermally produced DPs contribute only a small fraction of

the total dark matter density, their decay into the electromagnetic sector can leave observable imprints on astrophysical and cosmological phenomena. In the following sections, we examine sensitivities to probe a signal from the relic DP abundance via their decay by exploiting some observables, which align closely with predictions from the SM.

The cosmological evolution of the DP number density $n_{\gamma'}$ is governed by the Boltzmann equation,

$$\frac{dn_{\gamma'}}{dt} + 3Hn_{\gamma'} = \gamma_{\gamma'} \left(1 - \frac{n_{\gamma'}}{n_{\gamma'}^{\text{eq}}}\right), \quad (6)$$

where H accounts for the Hubble rate, and $\gamma_{\gamma'}$ is the DP interaction rate. For the sake of convenience, we rewrite the Boltzmann equation in terms of the comoving number density $Y_{\gamma'} = n_{\gamma'}/s$, as follows:

$$\frac{dY_{\gamma'}}{d \log T^{-1}} = \left(1 - \frac{1}{3} \frac{d \log g_{*s}}{d \log T^{-1}}\right) \frac{\gamma_{\gamma'}}{sH}, \quad (7)$$

where $s = (2\pi^2/45)g_{*s}T^3$ denotes the entropy density, T is the temperature, and g_{*s} represents the effective entropic degrees of freedom. Since DPs with MeV-scale masses are produced during the radiation-dominated era, the Hubble rate reads $H = 1.66 \times g_*^{1/2} T^2 / m_{\text{pl}}$ with $m_{\text{pl}} = G^{-1/2} = 1.22 \times 10^{19}$ GeV the Planck mass, and g_* the total number of relativistic degrees of freedom.

The interaction rate $\gamma_{\gamma'}$ is proportional to $(n_{\gamma'}^{\text{eq}} - n_{\gamma'})$, where $n_{\gamma'}^{\text{eq}}$ denotes the equilibrium number density of DPs. If the DP coupling strength (i.e., the kinetic mixing ϵ) is sufficiently large, the DP number density follows thermal equilibrium and eventually freezes out. On the other hand, for small kinetic mixing, DPs never reach thermal equilibrium but are gradually produced from the thermal bath via the so-called “freeze-in” production mechanism. In this freeze-in regime, the DP production rate is predominantly contributed from collisions that generate DPs, approximately given by $\gamma_{\gamma'} \propto n_{\gamma'}^{\text{eq}}$.

Building upon the description in Sec. II, the cosmological DP production rate can be expressed as

$$\gamma_{\gamma'} = \int d^3\vec{k} \frac{dN_{\gamma'}}{dV dt d^3\vec{k}}. \quad (8)$$

The plasma frequency ω_{pl} , which determines $\text{Re } \pi_i$ as outlined in Eqs. (A1) and (A2), is depicted in Fig. 2, incorporating the baryon asymmetry $n_B/n_\gamma \simeq 6 \times 10^{-10}$ with the proton abundance $n_p/n_B \simeq 0.878$ after BBN, leading to $n_{e^-} - n_{e^+} \simeq n_p$ due to the electric-charge neutrality condition. The imaginary component $\text{Im } \pi_i$ encodes information on the absorption (and, correspondingly, production in the detailed balance relation) of photon

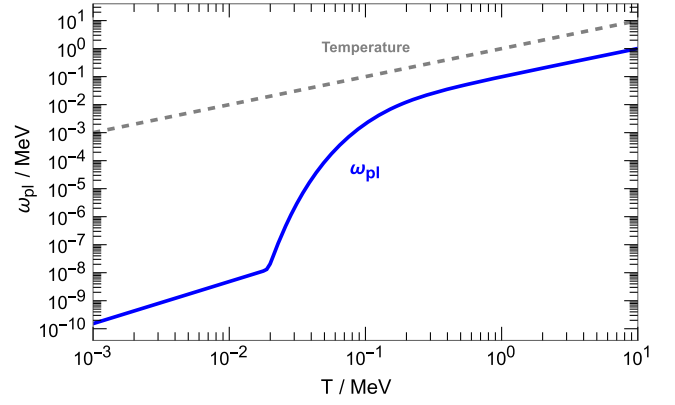


FIG. 2. Cosmic plasma frequency as a function of temperature.

states with the energy ω and the dispersion relation of $\omega^2 - k^2 = m_{\gamma'}^2$.

It is noteworthy that cosmological DP production is IR-dominated. This IR-sensitive feature stems from the renormalizable nature of DP couplings, governed by the dimensionless kinetic mixing ϵ . Moreover, at high temperatures of $\omega_{\text{pl}} > m_{\gamma'}$, plasma effects further suppress the effective DP couplings and, consequently, the interaction rate. Hence, the thermal production of cosmological DPs becomes efficient at lower temperatures when $\omega_{\text{pl}} \leq m_{\gamma'}$.

A notable contribution to the DP abundance can arise from the resonance point, $\text{Re } \pi_i = m_{\gamma'}^2$, similar to the case of supernovae discussed above. Since the resonance width is narrow (i.e., $\text{Im } \pi_i \ll \text{Re } \pi_i$), the resonance contribution to the DP abundance becomes independent of the imaginary components $\text{Im } \pi_i$, as discussed in Eq. (24), and can be understood as a nonadiabatic resonant conversion of photons into DPs. Details of the analytic derivation can be found in Appendix. B.

We notice that the resonant production is dominated by the transverse components, which occur when $\omega_{\text{pl}} \simeq m_{\gamma'}$. This is because the resonance of the longitudinal polarization occurs at the specific low energy of $\tilde{\omega} = \mathcal{O}(\omega_{\text{pl}})$, considerably smaller than temperature. Therefore, the photon number density available for resonant conversion into longitudinal DPs is suppressed by factors of $(\omega_{\text{pl}}/T)^2$ and $(1/\pi)^2$ relative to the equilibrium number density n_γ ; the second factor comes from the kinematic space integral. There is a compensatory enhancement of (T/ω_{pl}) since the nonadiabatic resonant conversion probability is inversely proportional to ω . As a result, the resonance contribution to the longitudinal DP abundance is approximately 2 orders of magnitude lower than that of transverse ones.

The thermal production of DPs persists until $T \sim m_{\gamma'}$. The conventional method for evaluating production rates assumes that statistical factors for the final states are neglected, and it exploits the advantage of invariance under Lorentz transformations to the final states. This approach

allows the use of Mandelstam variables to describe scattering processes [50] in the c.m. frame. However, in principle, this approach is not directly applicable to calculating the DP production rate because the plasma effect on the DP couplings depends on polarization, which is defined in a specific frame characterized by thermodynamic properties such as temperature.

Nevertheless, once the temperature drops to the lower temperature regime of $\omega_{\text{pl}} < m_{\gamma'}$, the plasma effect turns out to be insignificant. At such low temperatures, the effective coupling of DPs becomes effectively independent of polarization and behaves equivalent to the vacuum case. In light of this and considering that calculations are relatively insensitive to the initial temperature, the conventional Lorentz-invariant method remains a reliable approximation. In Appendix B 2, we provide the cosmological DP production rate $\gamma_{\gamma'}$ for the relevant scattering processes (semi-Compton scatterings and e^-e^+ annihilation) and the inverse-decay processes (the so-called “coalescence”) in terms of the Lorentz invariant variables.

Figure 3 shows the freeze-in DP abundance $Y_{\gamma'} = n_{\gamma'}/s$ for the initial temperatures $T_{\text{RH}} = 6$ MeV (black solid line) and $T_{\text{RH}} = 100$ MeV (gray dashed line). To compute $Y_{\gamma'}$, we perform the Boltzmann equation in Eq. (7) with the production rate given in Eq. (8), incorporating the plasma effect described in Sec. II. We take the effective degrees of freedom for the entropy (g_{*s}) and energy (g_*) densities within the standard cosmology from Ref. [51].

For $T_{\text{RH}} = 6$ MeV (black solid line), the DP abundance at masses below $\sqrt{3/2}\omega_{\text{pl}}(6 \text{ MeV}) \simeq 0.74$ MeV is dominated by resonant production, while at higher masses scattering processes without the plasma effect control the abundance. A sizable contribution from the inverse electron-positron decay arises only when it becomes kinematically allowed, i.e., for $m_{\gamma'} > 2m_e$. For the higher temperature of $T_{\text{RH}} = 100$ MeV, resonance production can

occur even for $m_{\gamma'} > 2m_e$, but its contribution is subdominant compared to inverse-decay process.

When the DP mass is sufficiently small compared to the initial temperature, the thermal relic abundance becomes insensitive to the initial temperature due to the IR-dominated nature of the production, as noted above. On the other hand, for too-heavy DP masses, abundance is exponentially suppressed by the Boltzmann factor, as clearly shown in Fig. 3. These results do not account for the subsequent decay of DPs; thus, the actual density at a given time should include an additional exponential suppression factor from decay.

The scaling behavior of the freeze-in DP abundance, in regimes where it becomes independent of the initial temperature (i.e., for $m_{\gamma'} \ll T_{\text{RH}}$), can be understood analytically as follows: for masses above the electron-positron threshold ($2m_e \simeq 1.02$ MeV), the dominant production mechanism is inverse decay, the rate of which is $\gamma_{\gamma'} \sim m_{\gamma'}^2 T \Gamma_{e\bar{e}} K_1(m_{\gamma'}/T) \propto \epsilon^2 m_{\gamma'}^2 T^2$ at $T > m_{\gamma'}$; see Eq. (B13). With this expression, we find

$$\begin{aligned} Y_{\gamma'} &\simeq \frac{135}{1.668\pi^3} \frac{g_{\gamma'} M_{\text{Pl}} \Gamma_{\gamma' \rightarrow e^+ e^-}}{g_s g_*^{1/2} m_{\gamma'}^2} S(x) \\ &\simeq \frac{8.2 \times 10^{17}}{m_{\gamma'}/\text{MeV}} \left(\frac{10.75}{g_s} \right)^{3/2} \\ &\quad \times S(x) \left(1 + \frac{2m_e^2}{m_{\gamma'}^2} \right) \sqrt{1 - 4 \frac{m_e^2}{m_{\gamma'}^2}} \end{aligned} \quad (9)$$

where $S(x) \simeq [0.266 + 0.734e^{-x}]e^{-x}(1+x^2)^{5/4}$, and $x \equiv m_{\gamma'}/T_{\text{RH}}$. Notice that for $x \gg 1$ (equivalently, $m_{\gamma'} \gg T_{\text{RH}}$), then the abundance gets suppressed as $S(x) \sim x^{5/2}e^{-x}$, while for $x \ll 1$, then $S(x) \sim 1$. Derivation of these results can be found in Appendix B. For lower masses, production is governed by the resonant scattering, which relies on the resonant temperature T_{res} . From Eq. (B2), the resonant contribution scales as $[Y_{\gamma'}]_{\text{res}} \propto \epsilon^2 m_{\gamma'}^2 T_{\text{res}}^{-3}$. When $m_{\gamma'} > 10^{-2}$ MeV, electrons and positrons are still relativistic at T_{res} that the plasma frequency scales as $\omega_{\text{pl}} \propto T$. This results in $Y_{\gamma'} \propto \epsilon^2 m_{\gamma'}^{-1}$, consistent with the inverse decay regime. In the intermediate mass range of $10^{-8} \text{ MeV} < m_{\gamma'} < 10^{-2} \text{ MeV}$, resonance occurs during the epoch when positrons are being annihilated. As illustrated in Fig. 2, T_{res} depends weakly on $m_{\gamma'}$ across this range, which leads to $Y_{\gamma'} \propto \epsilon^2 m_{\gamma'}^2$. At last, for $m_{\gamma'} < 10^{-8} \text{ MeV}$, positrons are depleted and the remaining asymmetric electron population determines the plasma frequency that $\omega_{\text{pl}} \propto n_e^{1/2} \propto T^{3/2}$. This results in a mass-independent abundance $Y_{\gamma'} \propto \epsilon^2$.

In the main text, we focus on the irreducible DP relic abundance for the initial temperature of 6 MeV,

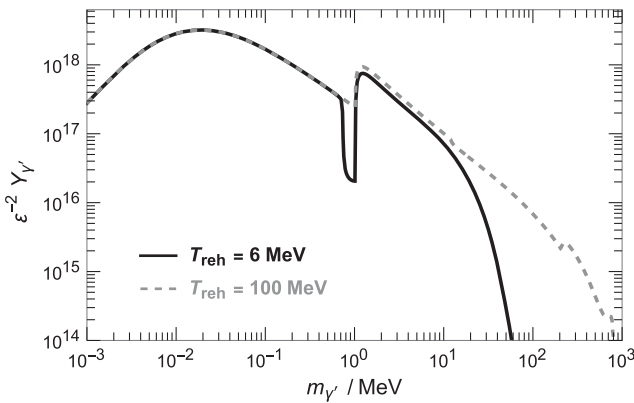


FIG. 3. Comoving dark photon density $Y_{\gamma'} = n_{\gamma'}/s$ generated by thermal processes with the initial temperatures $T_{\text{res}} = 6$ MeV (black) and $T_{\text{res}} = 100$ MeV (gray dashed), in the freeze-in scenario where thermal equilibrium is never achieved.

corresponding to the minimal onset temperature consistent with the standard cosmology (5.96 MeV at the 95% confidence level [52,53]), and explore the associated cosmological and astrophysical sensitivities to probe their signals. We also examine a higher temperature case as a more general case, with a detailed discussion provided in Appendix C.

The presence of a cosmological DP relic and its subsequent decay into the visible EM sector could leave an imprint on cosmological observations. In particular, cosmological probes are appropriate to searching for DPs with lifetimes shorter than the age of the Universe. Such short lifetimes are naturally realized for DP masses above the electron-positron threshold, where tree-level decay into electron-positron pairs becomes kinematically allowed with the rate described in Eq. (36). Hence, the phenomenological impacts of these decays during the early cosmological history become crucial for probing relatively heavy DPs. We examine the effects of decaying DPs with an irreducible abundance on various cosmological observations, including the CMB spectrum and the primordial abundance of light elements from BBN.

A. Overproduction

Excessive production of DPs can alter the expansion history of the Universe by increasing the total energy density beyond what is expected in the standard Λ cold dark matter cosmology. Once DPs are generated, they become nonrelativistic around $T \lesssim m_{\gamma'}$ and begin contributing to the matter density, which redshifts more slowly than radiation. An excess in matter abundance leads to a faster expansion rate in the early Universe and, in particular, brings forward the matter-radiation equality epoch. This shift in the expansion history affects various cosmological observables, including the growth of perturbations and the shape of the CMB power spectrum (e.g., through modifications to the early integrated Sachs-Wolfe effect and the acoustic peak structure).

In Fig. 4, the gray-shaded region labeled “ $\Omega_{\gamma'} > \Omega_{\text{rad}}$ ” (corresponding to the olive-shaded region in Fig. 1) is excluded based on the requirement that DPs must not dominate the energy density prior to matter-radiation equality. This overproduction bound appears only for masses below the electron-positron threshold, since at higher masses the DP lifetime is too short compared to the matter-radiation equality time $t_{\text{eq}} \simeq 1.2 \times 10^{12}$ sec. Even for lower masses, DPs can decay into three photons rapidly when the coupling ε is large, which prevents them from ever dominating the total energy density. However, as we discuss in the next subsection, there is an additional constraint arising from changes to the effective number of neutrino species on top of the SM contribution due to photon entropy injection by DP decay (red-shaded region in Fig. 1 and purple-shaded region in Fig. 4).

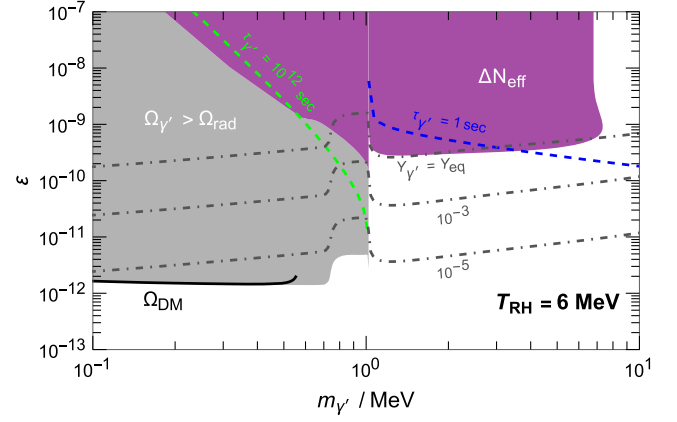


FIG. 4. Gray-shaded region is excluded by the overproduction of dark photons. The black solid line indicates the parameter values that reproduce the present dark matter abundance. The purple-shaded region is excluded by the ΔN_{eff} constraint (2σ confidence level). The green and blue dashed lines correspond to the DP lifetime of 10^{12} sec (around the CMB epoch) and 1 sec (around the neutrino decoupling), respectively. The gray dot-dashed lines indicate contours of the dark photon yield, $Y_{\gamma'} = 10^{-5}$, 10^{-3} , and the equilibrium value from bottom to top.

B. ΔN_{eff}

The decay of DP relics before the recombination era can leave observable imprints on the early Universe, such as the power spectrum of CMB anisotropies. If DPs decay after the neutrino decoupling epoch ($T \sim 2$ MeV), their energy is injected exclusively into the electron-photon plasma, while the neutrino sector remains unaffected. This selective heating increases the photon temperature relative to that of neutrinos, effectively reducing the fractional contribution of neutrinos to the total energy density. This effect is conventionally quantified by the effective number of neutrino species, $N_{\text{eff}} = (8/7)(11/4)^{4/3}(\rho_{\nu_e} + \rho_{\nu_\mu} + \rho_{\nu_\tau})/\rho_\gamma$, where ρ_{ν_i} and ρ_γ are the energy densities of ν_i flavor and photons, respectively. The SM prediction is $N_{\text{eff}} = 3.0440$ [54,55], while the most stringent observational constraint from the Planck Collaboration yields $N_{\text{eff}}^{\text{SM}} = 2.99 \pm 0.17$ at the time of recombination ($t_{\text{rec}} \simeq 1.2 \times 10^{13}$ sec). In the DP mass range of interest, where $m_{\gamma'} \gg T_{\text{CMB}}$, DPs cannot contribute to additional (invisible) radiation degrees of freedom. Instead, their decay reduces N_{eff} by injecting entropy solely into the photon bath.

The Planck constraint on N_{eff} implies limits on DPs with large kinetic mixing ε , where DPs thermalize with the SM plasma and subsequently decay, resulting in a suppressed N_{eff} . In this regime, the freeze-in approximation is no longer valid, and we compute the relic abundance and entropy injection by solving the full Boltzmann equations. We follow the procedure outlined in Ref. [56], while properly incorporating the plasma effect as discussed earlier. The purple-shaded region in Fig. 4 shows the

constraint from the Planck Collaboration on $|\Delta N_{\text{eff}} = N_{\text{eff}} - N_{\text{eff}}^{\text{SM}}| < 0.394$ at 2σ .

For DP masses above $2m_e$, the ΔN_{eff} constraint remains relevant up to about 7 MeV. Above this mass scale, the DP lifetime easily becomes much shorter than the neutrino decoupling time (~ 1 sec), so that the entropy injected by the decay of earlier-produced DP relics is efficiently distributed between the electron-photon plasma and the neutrino sector, which leaves negligible impact on N_{eff} . The small bump observed around $m_{\gamma'} \simeq 7$ MeV arises from the interplay between the DP lifetime and abundance, both of which control the efficiency of entropy injection.

We also derive the ΔN_{eff} constraint for DP masses below $2m_e$, where the relic abundance is substantially enhanced by resonant photon-DP conversion in the plasma, an effect that was not fully captured in earlier analyses [56]. In this low-mass regime, however, the loop-induced decay $\gamma' \rightarrow 3\gamma$ results in considerably long lifetimes. When the lifetime becomes comparable to or exceeds the recombination time, the corresponding entropy injection occurs too late to significantly modify the photon-to-neutrino temperature ratio. This interplay leads to the sharp, lifetime-dependent exclusion for $m_{\gamma'} < 2m_e$ seen in Fig. 4. Nevertheless, in this parameter region, DPs are so efficiently produced and long-lived that they already conflict with standard cosmology and its associated observables. In other words, the ΔN_{eff} constraint for $m_{\gamma'} < 2m_e$ lies entirely within the parameter space excluded by overproduction, as discussed in Sec. III A. For this reason, we do not include it in the main summary plot, Fig. 1.

Additional constraints related to ΔN_{eff} arise from BBN, which proceeds at temperatures of $T \sim 0.01$ –1 MeV. Due to higher temperatures, dark radiation from light DPs could, in principle, contribute to an increase of N_{eff} (i.e., $\Delta N_{\text{eff}} > 0$) during this BBN epoch, which impacts the expansion rate of the Universe as the cosmological parameter to govern dynamics in BBN, such as the freeze-out of the neutron-to-proton ratio. However, the production of lighter DPs is typically delayed relative to the freeze-out of weak interactions, because plasma effects strongly suppress DP–photon mixing at high temperatures. As a result, the direct contribution of DPs themselves to N_{eff} during BBN is negligible.

Furthermore, entropy injection from decaying DPs (particularly for $m_{\gamma'} > 2m_e$) can imprint observable effects on BBN. Such preferential entropy injection into the EM sector accounts for a reduction of the relative neutrino energy density (i.e., $\Delta N_{\text{eff}} < 0$), which alters not only the expansion rate but also neutrino-induced weak processes relevant for the proton-neutron interconversion [57–59]. In addition, entropy injection dilutes comoving baryon densities, which subtly shifts the predicted primordial element abundances [57,59]. We find that the resulting BBN constraints [57] are comparable to our Planck ΔN_{eff} bounds, as both depend on similar requirements for the DP

relic abundance and lifetime. Therefore, while BBN provides complementary information in terms of ΔN_{eff} , in our scenario it does not impose additional constraints beyond those derived from the CMB. Beyond the ΔN_{eff} considerations, energetic cascade photons produced by DP decay after BBN can dissociate light nuclei; these photodissociation bounds are discussed in detail in Sec. III D.

C. CMB anisotropy

An EM energy deposition through DP decay into electron-positron pairs, occurring after recombination, would heat up and ionize the neutral gas in the Universe. Such modifications to the cosmological thermal history during the dark ages leave observable imprints on the CMB, particularly in the temperature and polarization anisotropy spectra. As extensively discussed in literature [60–66], precise CMB measurements place stringent constraints on the properties of decaying (or annihilated) dark matters associated with energy injection.

To derive the CMB constraints on DPs, we follow the prescription outlined in Ref. [61], which provides a consistent fit to the results in the literature, using *Planck* 2018 data [67]. The CMB anisotropy constraints on DPs reads

$$\frac{m_{\gamma'} Y_{\gamma'}}{\rho_{\text{DM}}/s_0} \leq \frac{\tau_{\gamma'}}{\tau_{\text{min}}} \exp \left[\left(\frac{t_{\text{CMB}}}{\tau_{\gamma'}} \right)^{2/3} \right], \quad (10)$$

where $\tau_{\gamma'} = \Gamma_{e\bar{e}}^{-1}$ is the lifetime of nonrelativistic DPs, with the decay rate $\Gamma_{e\bar{e}}$ given by

$$\Gamma_{e\bar{e}} = \varepsilon^2 \frac{\alpha m_{\gamma'}^2}{3 \omega} \left(1 + \frac{2m_e^2}{m_{\gamma'}^2} \right) \sqrt{1 - 4 \frac{m_e^2}{m_{\gamma'}^2}}, \quad (11)$$

where at rest ($\omega = m_{\gamma'}$). Here, τ_{min} refers to the corresponding lifetime constraints for dark matters decaying into electron-positron pairs with a mass of $m_{\gamma'}$, as derived in Ref. [62], and we take $t_{\text{CMB}} \simeq 8.7 \times 10^{13}$ sec. The blue-shaded region in Figs. 1 and 5 describes the CMB constraint.

Notice that the lower part of the curve can be derived analytically. Ignoring the exponential in Eq. (10), and considering that the bound from Ref. [62] is almost flat between 1–10 MeV (around $\sim 3 \times 10^{24}$ s), one finds

$$\varepsilon \gtrsim 1.5 \times 10^{-17} \left(\frac{\text{MeV}}{m_{\gamma'}} \right)^{1/4}, \quad (12)$$

which is in good numerical agreement with the plot. The upper part of the curve instead coincides with having $\tau_{\gamma'} \simeq t_{\text{CMB}}$.

One may also consider constraints from distortions of the CMB energy spectrum, which is precisely measured and tightly constrained by cosmic background explorer–far infrared absolute spectrophotometer [68]. There are two

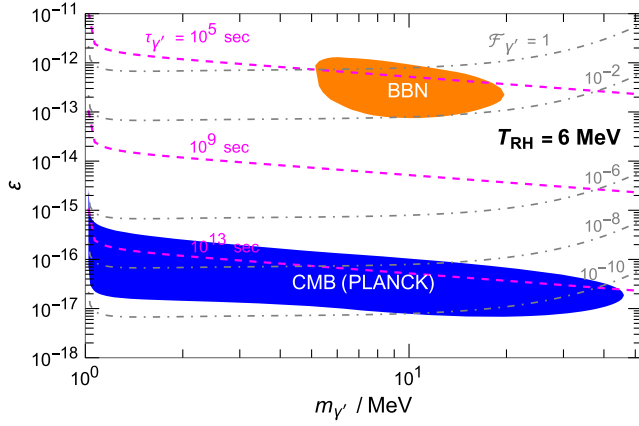


FIG. 5. Blue-shaded region corresponds to the CMB constraint from the latest *Planck* results. The orange-shaded region is excluded by BBN.

relevant types of spectral distortions, μ -type and y -type, which rely on photon-scattering processes that redistribute the injected energy (e.g., double Compton and Compton); see Refs. [65,66] and references therein. However, within the DP parameter space of interest, the fractional energy injection required to generate a detectable spectral distortion is substantially larger than that needed to imprint measurable changes on the CMB anisotropy power spectra. Consequently, current FIRAS sensitivities may not yield competitive constraints on DPs compared to the anisotropy bounds discussed above.

D. BBN

BBN offers a powerful probe of physics beyond the Standard Model, as it took place only a few seconds to minutes after the big bang, when light nuclei first formed. The observed primordial abundances of light elements are in remarkable agreement with the predictions of standard cosmology, and even small deviations can thus impose stringent constraints on new particles that could alter nuclear abundances.

The decay of DPs into the EM sector can modify the final outcome of BBN. Once the Universe cools below $T \simeq 0.1$ MeV, photons energetic enough to dissociate nuclei (especially deuterium) become too rare to counteract nuclear synthesis. At this stage, roughly three minutes after the big bang, the formation of light elements ceases and their primordial abundances become fixed.

The decay of DPs into electron-positron pairs triggers an EM cascade, as the injected electrons, positrons, and photons interact with the background plasma. Assuming that the cascade develops instantaneously, the resulting zero-generation photon spectrum takes the form of a broken power law with a sharp cutoff at the minimum value of $\{m_e^2/22T, m_{\gamma'}/2\}$, reflecting photon-photon pair creation processes [69]. That is to say, high-energy photons are efficiently dissipated before they can interact with nuclei. For details of this analytic spectrum, see Refs. [70,71].

Because nuclear binding energies exceed ~ 1 MeV, photo-dissociation by the EM cascade becomes relevant only after BBN is complete. We therefore use the standard primordial abundances as initial conditions, adopting the light-element yields computed with PArthENoPEedstar 3.0 [72–74].

We solve the Boltzmann equation, including nuclear destruction terms from energetic cascade photons produced by DP decays. The stationary distribution of these cascade photons depends on the DP number density and decay rate, as well as on the zero-generation photon spectrum discussed above. It results from the balance between photon injection and energy degradation through scatterings with the background medium [75,76]. Analytic cross sections for all relevant photo-dissociation channels, validated against experimental data, are taken from Ref. [77].

For the observationally inferred light-element abundances, we adopt the recommended Particle Data Group values [78]:

$$D/H = (25.47 \pm 0.29) \times 10^{-6}, \quad (13)$$

$$Y_p = 0.245 \pm 0.003. \quad (14)$$

Accounting for both observational and theoretical uncertainties—which for D/H amount to $\sigma_D^{\text{th}} = 0.6 \times 10^{-6}$ (arising from nuclear reaction rates), and are negligible for Y_p compared with the observational one—we derive the corresponding BBN constraints at the 2σ level. For DP masses around $m_{\gamma'} \sim 10$ MeV—relevant given the low initial temperature $T_i \simeq 6$ MeV—the dominant constraint arises from the photodissociation of primordial deuterium.

We note that recent Extremely Metal-Poor Representatives Explored by the Subaru Survey [79] observations report a low $Y_p = 0.2370^{+0.0034}_{-0.0033}$. Since our BBN bound is driven primarily by the deuterium abundance (with the impact on Y_p being subdominant for the DP energy-injection mechanism considered here), adopting this alternative helium value would shift/broaden the helium likelihood but does not qualitatively change the BBN constraint.

E. X-rays

The decay of cosmological DPs into the EM sector can contribute to various astrophysical observations, such as x-ray flux from the Galactic Center or diffuse x-ray background. For the decay of DPs to leave a significant observational imprint, their lifetime is required to be comparable to or longer than the age of the Universe. This long-lifetime condition can be naturally achieved for DPs with masses below $2m_e$, where the leading decay channel, $\gamma' \rightarrow \gamma\gamma\gamma$, is induced at the one-loop level. The corresponding rest-frame decay rate $\Gamma_{3\gamma}$ is given by [80,81]

$$\Gamma_{3\gamma} = \epsilon^2 \frac{17\alpha^4}{2^7 3^6 5^3 \pi^3} \frac{m_{\gamma'}^9}{m_e^8} \mathcal{G}(m_{\gamma'}), \quad (15)$$

where $\mathcal{G}(m_{\gamma'})$ accounts for deviations from the Euler-Heisenberg approximation at higher masses; see Ref. [81]. The inclusive photon spectrum per DP decay, defined as $dN_{\gamma}/d\hat{x} \equiv 3(d\Gamma_{3\gamma}/d\hat{x})/\Gamma_{3\gamma}$ with $\hat{x} = E_{\gamma}/(m_{\gamma'}/2)$ in the range of $[0, 1]$, is described by

$$\frac{dN_{\gamma}}{d\hat{x}} = 3 \times \frac{\hat{x}^3}{51} \left(1715 - 3105\hat{x} + \frac{2919}{2}\hat{x}^2 \right), \quad (16)$$

where the factor of 3 accounts for the three photons produced in each decay.

Photons stemming from the DP radiative decay can result in an excess over the expected photon background flux. Given that current x-ray observations are consistent with SM predictions within instrumental sensitivities, any such excess provides a sensitive probe of DP relics. The expected differential Galactic x-ray flux from DP decay reads

$$\frac{d\Phi_{\gamma}}{dE_{\gamma}d\Omega} = c \left[\frac{2}{m_{\gamma'}} \frac{dN_{\gamma}}{d\hat{x}} \right] \times \left[\frac{\Gamma_{3\gamma}}{4\pi m_{\gamma'}} \int ds \rho_{\gamma'0} \right], \quad (17)$$

where c is the speed of light and $\rho_{\gamma'0} = m_{\gamma'} Y_{\gamma'} s_0 e^{-\Gamma_{3\gamma} t_0}$ is the current dark photon relic density, accounting for exponential depletion via decay. Here, $Y_{\gamma'}$ is the yield in the absence of subsequent decays, s_0 is the current entropy density, and s denotes the line of sight. Integrating over the area of interest introduces the parameter $D_{\gamma'} = \int ds d\Omega \rho_{\gamma'0}$, which indicates the dark photon column density for the corresponding observed region. In our analysis, we assume that the local dark photon density profile follows the morphology of the total dark matter density; i.e., that the local and global ratios of dark photons with respect to the whole dark matter (DM) abundance remain consistent as $D_{\gamma'}/D = \rho_{\gamma'0}/\hat{\rho}_{\text{DM}}$ with the global dark matter density $\hat{\rho}_{\text{DM}}$.

There are several x-ray observations of various galaxies where the observed flux is consistent with the SM prediction within the instrumental sensitivity limits: X-ray Multi-Mirror Mission-Newton [82–85], Nuclear Spectroscopic Telescope Array [86–88], and International Gamma-Ray Astrophysics Laboratory [89]. These studies provide constraints on the lifetime of single-component dark matter decaying into two photons, which imposes a monochromatic photon energy spectrum at the energy of half the dark matter mass for each decay.

For DPs with sub-MeV masses, however, the trident decay channel $\gamma' \rightarrow 3\gamma$ produces a continuum photon spectrum rather than a line, resulting in a broader x-ray feature. We adopt the analysis of Ref. [90], which derived Galactic x-ray constraints on DP dark matter using INTEGRAL data. Since Ref. [90] assumes that DPs constitute the entire dark matter abundance, we rescale their lifetime limit as

$$\frac{m_{\gamma'} Y_{\gamma'}}{\rho_{\text{DM}}/s_0} \leq \frac{\tau_{\gamma'}}{\tau_{\text{DM}}} \exp \left[\frac{t_{\text{U}}}{\tau_{\gamma'}} \right], \quad (18)$$

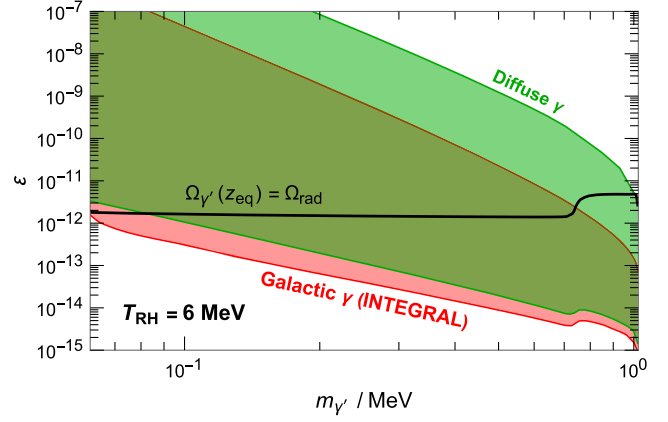


FIG. 6. Red- and green-shaded regions are excluded by Galactic and extragalactic x-ray observations, respectively, which constrain the x-ray flux from dark photon decays not to exceed the observed background within the corresponding instrumental sensitivities. The black solid line indicates the upper limit from the overproduction argument.

where t_{U} is the age of the Universe and τ_{DM} denotes the corresponding lifetime constraint reported in Ref. [90]. The red-shaded region in Fig. 6 is excluded by the equation above.

Besides their contribution to the present Galactic x-ray flux, the radiative decay of DPs during earlier cosmic epochs can also impact the diffuse extragalactic x-ray background. For masses around the MeV scale of interest, the production of DPs occurs well before the recombination era. Moreover, DPs become nonrelativistic at the time of decay, which happens much later than the recombination era to contribute to the cosmic radiation background. Under these conditions, the diffuse photon background flux from the decay of cosmic DPs reads

$$\begin{aligned} \frac{d\Phi_{\gamma}}{dE_{\gamma}d\Omega} &= \frac{Y_{\gamma'} s_0 \Gamma_{3\gamma}}{4\pi E_{\gamma}} \int_0^1 d\hat{x} \frac{c}{H(\hat{z})} e^{-\Gamma_{3\gamma} t(\hat{x})} \frac{dN_{\gamma}}{d\hat{x}} \\ &\times e^{-\kappa(\hat{z}, E_{\gamma})} \Theta \left(\hat{x} - \frac{E_{\gamma}}{m_{\gamma'}/2} \right), \end{aligned} \quad (19)$$

with

$$t(\hat{x}) = \frac{2}{3H_0\sqrt{\Omega_{\text{m}}}} \sinh^{-1} \sqrt{\frac{\Omega_{\Lambda}}{\Omega_{\text{m}}(\hat{x}(m_{\gamma'}/2)/E_{\gamma})^3}}. \quad (20)$$

Here, $\kappa(\hat{z}, E_{\gamma})$ accounts for the optical depth, which factors in the attenuation of photons during propagation with the present energy E_{γ} emitted at a redshift $\hat{z} = \hat{x}(m_{\gamma'}/2)/E_{\gamma}$; we exploit the result from Ref. [91] to describe this attenuation, in particular using their public *Mathematica* notebooks. Based on the given upper limits on the cosmic background radiation in Ref. [92], the green-shaded region

in Fig. 6 illustrates the constraint from the diffuse x-ray background.

IV. SUPERNOVA DARK PHOTONS

During a core-collapse supernova (CCSN), the iron core of a massive red supergiant becomes gravitationally unstable and collapses under its own weight [93–96]. This collapse halts when nuclear densities are reached, leading to the formation of a stiff, neutron-rich object known as a protoneutron star (PNS). In the first few seconds after core bounce, the temperature of the PNS can reach approximately ~ 30 – 50 MeV. Hence, massive DPs with masses in the MeV to several hundred MeV range can be copiously produced within the PNS. Based on Eq. (5), the DP luminosity per volume is given by

$$\begin{aligned} \frac{d\mathcal{P}_{\gamma'}}{dV} &= \int d^3\vec{k} \frac{d\mathcal{N}_{\gamma'}}{dV dt d^3\vec{k}} \omega \\ &= \frac{\varepsilon^2 m_{\gamma'}^4}{2\pi^2} \int \frac{d\Omega_{\gamma'}}{4\pi} d\omega \frac{\omega^2 v_{\gamma'}}{e^{\omega/T} - 1} \\ &\quad \times \sum_{i=\text{T,L}} g_i \frac{|\text{Im } \pi_i|}{(m_{\gamma'}^2 - \text{Re } \pi_i)^2 + |\text{Im } \pi_i|^2}, \quad (21) \end{aligned}$$

with $v' = \sqrt{1 - m_{\gamma'}^2/\omega^2}$. Here, we retain the angular integral of $d\Omega_{\gamma'}/4\pi$. As we will see, although produced DPs radiate isotropically, their absorptive width along the propagation trajectory, i.e., optical depth, exhibits angular dependence.

In a PNS, electrons are highly degenerate and relativistic, $(3\pi n_e)^{1/3} \gg T, m_e^{\text{eff}}$, where n_e is the electron number density and m_e^{eff} the electron effective mass in the medium. The plasma frequency, $\omega_{\text{pl}}^2 \simeq (64\pi/3)^{1/3} a n_e^{2/3}$, is of order $\mathcal{O}(10)$ MeV. Due to the form of the in-medium photon polarization tensors, which we report in the Appendix A, Eqs. (A1) and (A2) that remain valid in the dense PNS environment as the hard-dense-loop approximation [97], we notice that resonances for the transverse mode occur when the DP mass is $\mathcal{O}(10)$ MeV. This is because the value of $\text{Re } \pi_{\text{T}}$ is on the order of the plasma frequency across the entire energy range. On the other hand, for longitudinal modes, $\text{Re } \pi_{\text{L}} \sim 3\omega_{\text{pl}}^2 m_{\gamma'}^2/\omega^2$, and a resonance energy $\omega \sim \omega_{\text{pl}}$ always exists for the low mass regime, $m_{\gamma'} < \omega_{\text{pl}}$.

The dominant production channel for DPs inside a PNS is bremsstrahlung in neutron-proton scatterings ($n + p \rightarrow n + p + \gamma'$), by which vector bosons can be radiated at the leading dipole order [98,99] when the c.m. is not consistent with the center of charge. We employ two valid approximations that make the computations simpler, the soft-radiation approximation [100,101] and the nondegenerate and nonrelativistic conditions of the nuclear medium (for details, see Refs. [98,102]). Under

these approximations, the scattering amplitude can be separated into the neutron-proton scattering part and the photon interaction with the dipole current. We exploit the neutron-proton scattering cross section derived from empirical data as shown in Ref. [98].

Furthermore, we take into account the multiple coherent scattering effects [103,104], which are significant in a high-density medium of a PNS. The absorption photon rate via the inverse neutron-proton bremsstrahlung is given by

$$\begin{aligned} \Gamma_i^{\text{ibr}} &= \frac{32\alpha n_n n_p}{3\pi \omega^3} \left(\frac{\pi T}{m_N} \right)^{3/2} \\ &\quad \times \langle \sigma_{np}^{(2)}(T) \rangle \left(\frac{\omega^2}{\omega^2 + \Gamma_{np}^2/4} \right) \times \left[\frac{m_{\gamma'}^2}{\omega^2} \right]_{i=\text{L}}, \quad (22) \end{aligned}$$

where $n_{n(p)}$ denotes the number density of neutrons (protons), and $\langle \sigma_{np}^{(2)}(T) \rangle$ is the thermally averaged neutron-proton scattering cross section taken from Ref. [98]. The last term appears only for the longitudinal component due to the current conservation. The factor of $\omega^2/(\omega^2 + \Gamma_{np}^2/4)$ with $\Gamma_{np} \approx n_n \langle \sigma_{np}^{(2)}(T) \rangle (3T/m_N)^{1/2}$ represents the Lorentzian shape ansatz accounting for the high-density behavior in neutron-proton bremsstrahlung scatterings [105].

We remark that negatively charged particles heavier than electrons, most notably muons [106] and charged pions [107], and, in nonminimal scenarios with additional (e.g., flavor-violating) couplings, hyperon-induced transition channels [108] may also contribute to DP production. Their thermal abundance in the PNS can reach the $\sim 1\%$ level relative to nucleons due to sizable chemical potentials comparable to their masses (~ 100 MeV). While their overall contribution to the total emissivity is subdominant, these species can populate the high-energy tail of the DP spectrum, potentially leaving observable signatures that depend sensitively on the spectral shape rather than the integrated luminosity. However, since the abundances of these particles carry significant model and equation-of-state (EOS) uncertainties, we do not include them in our numerical analysis; for a discussion of their possible impact, see Ref. [109], and Ref. [110] for the case of axions.

In the region far outside the PNS, DP production (and also absorption) is governed by the semi-Compton process, $e^- + \gamma \rightarrow e^- + \gamma'$, which is less sensitive to the medium density than bremsstrahlung scatterings. While semi-Compton scatterings in these outer layers contribute negligibly to the total DP luminosity, they can nonetheless have a nontrivial impact on supernova dynamics, particularly shock revival for a successful explosion [111]. In Appendix D, we present an approximate analytic derivation of the semi-Compton absorption rate, which is related to the production rate via detailed balance, under various relevant conditions.

The integral over the momentum range of the DP production rate in Eq. (21) can be divided into two regions,

around the resonance energy and outside the resonance energy, as follows:

$$\frac{d\mathcal{P}_{\gamma'}}{dV} = \frac{d\mathcal{P}_{\gamma'}}{dV}\bigg|_{\text{res}} + \frac{d\mathcal{P}_{\gamma'}}{dV}\bigg|_{\text{out}}. \quad (23)$$

Due to $\text{Im } \pi_i \ll \text{Re } \pi_i$ at the resonance energy, denoted by ω_{res} , we can approximate

$$\begin{aligned} & \frac{1}{(m_{\gamma'}^2 - \text{Re } \pi_i)^2 + |\text{Im } \pi_i|^2} \\ & \simeq \frac{1}{\left(\frac{\partial \text{Re } \pi_i}{\partial \omega}\right)^2 (\omega - \omega_{\text{res}})^2 + |\text{Im } \pi_i|^2} \\ & \simeq \delta(\omega - \omega_{\text{res}}) \times \pi \left(\left| \frac{\partial \text{Re } \pi_i}{\partial \omega} \right| |\text{Im } \pi_i| \right)^{-1}. \end{aligned} \quad (24)$$

Therefore, the production rate from the resonance reads

$$\frac{d\mathcal{P}_{\gamma'}}{dV}\bigg|_{\text{res}} \simeq \frac{\varepsilon^2 m_{\gamma'}^4}{2\pi} \int \frac{d\Omega_{\gamma'}}{4\pi} \sum_{i=\text{T,L}} g_i \left| \frac{\partial \text{Re } \pi_i}{\partial \omega} \right|^{-1} \frac{\omega_{\text{res}}^2 v_{\text{res}}}{e^{\omega_{\text{res}}/T} - 1}, \quad (25)$$

where $v_{\text{res}} = \sqrt{1 - m_{\gamma'}^2/\omega_{\text{res}}^2}$ and

$$\left| \frac{\partial \text{Re } \pi_i}{\partial \omega} \right| = \frac{m_{\gamma'}^2}{\omega_{\text{res}} v_{\text{res}}^2} \left(2 + \frac{m_{\gamma'}^2}{\omega_{\text{res}}^2} - \frac{3\omega_{\text{pl}}^2}{m_{\gamma'}^2} \left[\frac{m_{\gamma'}^2}{\omega_{\text{res}}^2} \right]_{i=\text{L}} \right). \quad (26)$$

The contribution from outside the resonance energy is explicitly calculated from Eq. (21) with the neutron-proton bremsstrahlung scattering rate given in Eq. (22).

A. Implications

Dark photons produced in CCSNe can affect several observable aspects of the explosion and its aftermath. Building on the DP production rate described above, we consider multiple supernova observables to derive constraints on the DP parameter space. For our fiducial analysis, we adopt the LS220-s18.88 supernova simulation profile, which represents the one-dimensional counterpart of a self-consistent three-dimensional explosion model with a PNS mass of $1.81M_{\odot}$ [112]. Figure 7 illustrates the time evolution of the temperature (upper panel) and the radial profiles of the temperature (red) and plasma frequency (blue) at 0.2 sec (dashed) and 1 sec (solid) postbounce (lower panel). The profiles at $t_{\text{pb}} \simeq 0.2$ s are particularly relevant for the DP case, since for some DP masses the cooling rate can be larger at earlier times due to the nature of resonant production (this is somewhat peculiar; another similar case is the Majoron model studied in Ref. [113]). We also compare our results with four additional models from the Garching supernova archive, thus bracketing the uncertainties associated with the SN modeling.

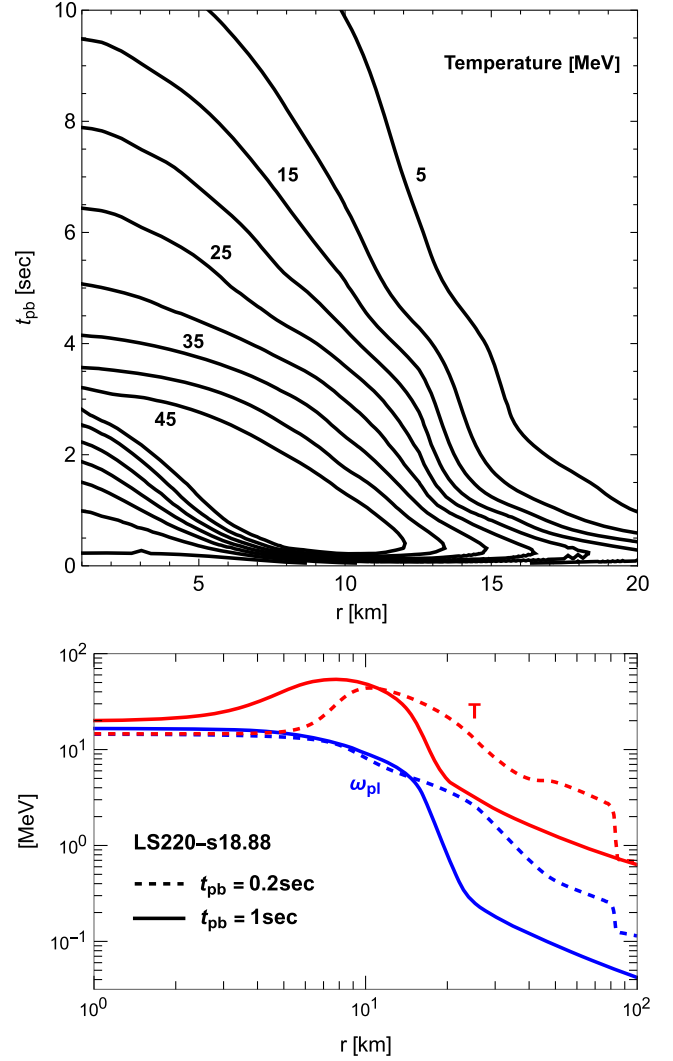


FIG. 7. Upper panel: time evolution of the temperature in the LS220-s18.88 supernova model; contour lines correspond to the indicated value in MeV. Lower panel: radial profiles of temperature (red) and plasma frequency (blue) at 0.2 sec (dashed) and 1 sec (solid) postbounce, respectively.

We also want to stress that no clear evidence of a compact remnant has yet been observed at the location of SN1987A, and in particular no pulsarlike signal. In principle, this could call into question a neutrino-driven explosion mechanism for SN1987A, although the remnant could simply be a nonpulsar neutron star, and the hot blob recently reported by the Atacama Large Millimeter Array radio telescope and the James Webb Space Telescope can be interpreted as a hint of such a source (see Sec. VIII. 6 in Ref. [114] for a more detailed discussion about this). In any case, among the arguments that we examine here, only the standard cooling argument and the SN failure argument rely on a specific, neutrino-driven explosion mechanism. By contrast, low-energy supernovae (LESNe), fireball formation, supernova γ -ray signals, and Galactic positron

injection are independent on the explosion mechanism and, taken together, supersede the standard cooling argument, making SNe a robust probe. Nevertheless, we recall that much above the free-streaming boundaries, self-consistent simulations would be desirable.

1. Cooling argument

The gravitational binding energy of the nascent PNS ($\approx 3 \times 10^{53}$ erg) is predominantly released in the form of neutrinos within the Standard Model. For the first few seconds after core bounce, neutrinos remain trapped in the dense stellar matter surrounding the PNS, forming a neutrino sphere. Energy diffusion through this region leads to a cooling phase lasting $\mathcal{O}(10)$ seconds. The observed duration of the neutrino burst from SN1987A, detected by Kamiokande-II, Irvine–Michigan–Brookhaven, and Baksan [115–117], agrees remarkably well with this prediction. To remain consistent with these observations, any exotic energy-loss mechanism must be subdominant to the standard neutrino cooling. This requirement is encapsulated in the so-called *Raffelt criterion* [110,118], which constrains the specific energy-loss rate of new particles, ensuring that the overall cooling timescale is not significantly shortened.

DPs produced in PNS can contribute to extra energy leakage if they escape the region cooled down by neutrino diffusion. In the course of propagation, DPs may either be absorbed into the nuclear medium through scatterings or decay into electron-positron pairs if the mass exceeds the two-electron mass threshold. This absorption process depends on the production location of a DP and its direction of travel. The optical depth for a DP produced at radius r with energy ω , traveling at an angle θ_γ relative to the radial direction, is given by

$$\tau_{\text{cool}} = \int d\hat{l} \hat{\Gamma}_\gamma \left(\sqrt{r^2 + l^2 + 2rl \cos \theta_\gamma} \right) / v', \quad (27)$$

where l represents the path length from the production point ($l \geq 0$), and $\hat{\Gamma}_\gamma$ is the DP absorptive rate. The effective cooling rate due to DPs is then expressed as

$$\mathcal{P}_\gamma|_{\text{cool}} = \int dV \frac{d\mathcal{P}_\gamma}{dV} e^{-\tau_{\text{cool}}}, \quad (28)$$

with the volume emissivity $d\mathcal{P}_\gamma/dV$ given by Eq. (21).

For DP masses below $2m_e$, the optical depth is primarily determined by scatterings near a PNS due to relatively higher density, allowing integration over arbitrarily large distances. On the other hand, for masses above $2m_e$, DPs can decay into electron-positron pairs outside the PNS, which leads to a significant contribution to the optical depth without suppression at large radii. In our analysis, we set an upper integration limit for radius $r_{\text{up}} = 20$ km (i.e., the corresponding upper limit of

$l = \sqrt{r_{\text{up}}^2 - r^2 \sin^2 \theta_\gamma} - r_{\text{up}} \cos \theta_\gamma$), which appropriately accounts for the geometric size of the PNS and the neutrino sphere.

The DP absorption rate is written as

$$\hat{\Gamma}_\gamma \simeq \frac{\varepsilon^2 m_\gamma^4}{(m_\gamma^2 - \text{Re } \pi_i)^2 + |\text{Im } \pi_i|^2} \times (1 - e^{-\omega/T}) [\Gamma_i^{\text{ibr}} + \Gamma_i^{\text{sc}} + \Gamma_i^{e\bar{e}}] \quad (29)$$

for each polarization mode $i = \text{T, L}$. Inside the PNS, the absorption of DPs is controlled by the inverse neutron-proton bremsstrahlung, the rate (Γ_i^{ibr}) of which is described in Eq. (22). On the other hand, the semi-Compton scattering process, $e^- \gamma' \rightarrow e^- \gamma$, dominates the DP absorption outside of the PNS due to its weaker power dependence of the number density (i.e., $\propto n_e$). For DP masses above $2m_e$, an additional absorptive channel opens up via decay into electron-positron pairs. Ignoring the positron abundance, the decay rate $\Gamma_i^{e\bar{e}}$ is given by

$$\Gamma_i^{e\bar{e}} \simeq \frac{1}{16\pi \sqrt{\omega^2 - m_\gamma^2}} \int_{x_-}^{x_+} dx \frac{|\bar{\mathcal{M}}_i|^2}{e^{\frac{\mu_e - \omega x}{T}} + 1}, \quad (30)$$

where μ_e is the chemical potential of electrons, and $x_\pm = (1 \pm \sqrt{(1 - 4m_e^2/m_\gamma^2)(1 - m_\gamma^2/\omega^2)})/2$ defines the integration range. The squared matrix elements read

$$|\bar{\mathcal{M}}_i|^2 = 8\pi\alpha \times \begin{cases} (m_\gamma^2 + 2m_e^2) - \xi(x) & [i = \text{T}] \\ 2\xi(x) & [i = \text{L}], \end{cases} \quad (31)$$

with $\xi(x) = (2x(1-x)\omega^2 - m_\gamma^2/2)m_\gamma^2/(\omega^2 - m_\gamma^2)$; for the detailed derivation, see Ref. [102].

Figure 8 shows the DP luminosity as a function of the kinetic mixing parameter ε for several representative masses. In the low-coupling regime, where DPs freely escape from the PNS, the luminosity scales as ε^2 . At larger couplings, the typical optical depth in Eq. (27) exceeds unity, leading to a suppression of the effective DP luminosity due to increased trapping and reabsorption. For lighter masses, a peak appears at intermediate couplings, where resonantly produced DPs are also resonantly reabsorbed near their production region, as shown by the red and blue lines in Fig. 8 for $m_\gamma = 1$ MeV and 2 MeV, respectively. This implies that at higher couplings, the DP luminosity becomes dominated by nonresonant production processes. In this context, heavier DPs (such as the case with $m_\gamma = 30$ MeV shown by the green line) exhibit only a single peak due to the absence of resonance effects. For DP masses below the $2m_e$ threshold, where decay into electron-positron pairs is kinematically forbidden, absorption is dominated by scattering processes; above this threshold,

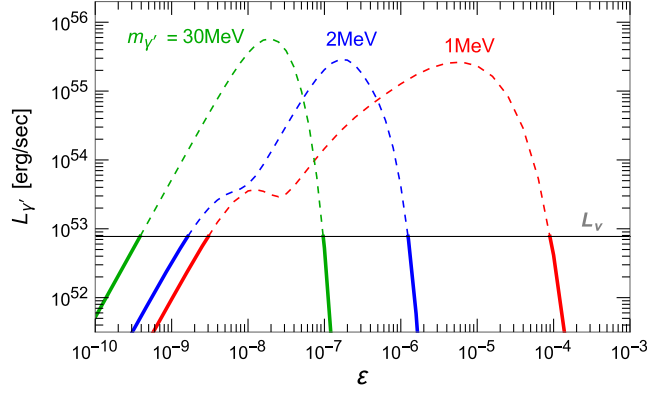


FIG. 8. Dark photon luminosity as a function of the kinetic mixing parameter ϵ for masses of 1 MeV (red line), 2 MeV (blue line), and 30 MeV (green line), evaluated at 1 sec postbounce. The gray line indicates the neutrino luminosity.

the additional decay channel into electron-positron pairs significantly enhances the optical depth.

Consequently, we derive cooling constraints from the SN1987A neutrino observations by comparing the effective DP cooling rate at 1 sec postbounce with the corresponding neutrino luminosity (see, e.g., the gray line in Fig. 8). The red region in Fig. 9 illustrates the parameter space where the effective DP emission rate exceeds the diffusive neutrino luminosity, based on the LS220-s18.88 simulation profile. We additionally show cooling constraints obtained from other supernova models of the Garching group [119,120] with different progenitor masses and/or nuclear equations of state: LS220-s20.0 (blue dashed), SFHo-s18.8 (green dashed), and SFHo-s18.6 (cyan dashed). The LS220 EOS employs a Lattimer-Swesty compressible liquid-drop model with an incompressibility of $K = 220$ MeV, while

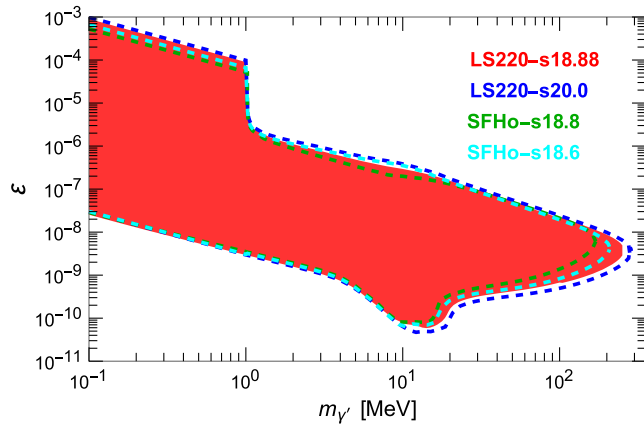


FIG. 9. Constraints from the SN1987A cooling argument for different SN models: LS220-s18.88 (red-shaded region), LS220-s20.0 (blue dashed line), SFHo-s18.8 (green dashed line), SFHo-s18.6 (cyan dashed line).

the SFHo EOS [121] is based on relativistic mean-field theory and yields a more compact PNS with slightly higher central densities and temperatures. The comparison across these simulations, covering progenitor masses between 18.6 and $20.0 M_{\odot}$, allows us to gauge the sensitivity of our results to uncertainties in the progenitor structure and the nuclear microphysics. Overall, the resulting constraints are largely consistent across most of the parameter space, with some deviations only in the trapping regime for $m_{\gamma} \lesssim 1$ MeV—where the result is sensitive to the target density for scattering—and in the heavy-mass region, where the constraint depends more strongly on the maximum temperature of the profile.

We also compare our updated results with those previously reported in the literature [102]. Figure 10 shows this comparison, where the black solid line represents the result from Ref. [102] obtained using their fiducial supernova profile. As pointed out by the authors of Ref. [110], their evaluation of the optical depth assumes a purely radial propagation [i.e., $\cos \theta_{\gamma} = 1$ in Eq. (27)] with an additional geometric correction. This prescription tends to overestimate the optical depth, leading to a stronger constraint in the trapping regime. Nevertheless, their bound for $m_{\gamma} \lesssim 1$ MeV is less stringent than ours because their fiducial profile contains a higher electron density outside the PNS, which enhances the optical depth through the semi-Compton scattering contribution. For comparison, we confirm that adopting a purely radial optical depth treatment (blue dashed line) indeed reproduces this overestimation.

Finally, we note that the Stefan-Boltzmann (SB) approach provides a good approximation to the cooling bound in the trapping regime for low-mass DPs, where m_{γ}

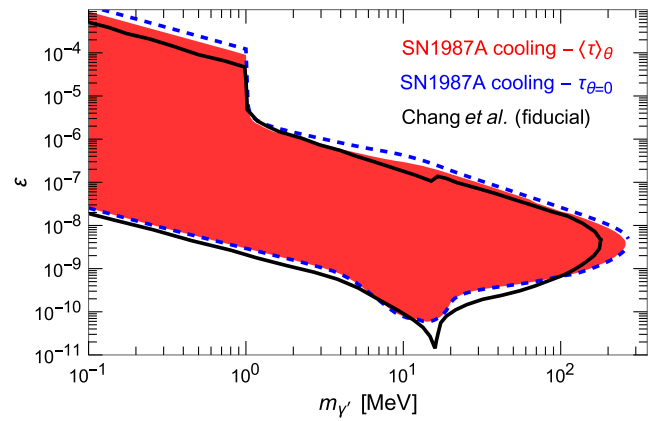


FIG. 10. Our SN1987A cooling constraints (red-shaded region) using the Garching LS220-s18.88 model, compared with previous results from Ref. [102]. We also show results (blue dashed line) obtained by evaluating the optical depth assuming purely radial propagation. As pointed out by Ref. [110], such an incorrect procedure leads to an overestimation of the constraints.

is smaller than the typical temperature at the PNS boundary ($T \sim 10$ MeV). The differential SB luminosity [110] reads

$$\frac{d\mathcal{P}_{\gamma'}^{\text{SB}}}{d\omega} = \sum_{i=\text{T,L}} 4\pi R_\omega^2 \frac{g_i}{8\pi^2} \frac{\omega^3}{e^{\omega/T(R_\omega)} - 1}, \quad (32)$$

where R_ω denotes the radius at which the radial optical depth becomes $2/3$ (i.e., $[\tau_{\text{cool}}]_{\cos\theta_{\gamma'}=1} = 2/3$). Since the DP absorption rates from bremsstrahlung and semi-Compton processes depend on the DP energy ω , each energy corresponds to a different diffusion radius R_ω . In particular, the absorption rate decreases with increasing ω , implying that the dominant emitted DP spectrum in the trapping regime shifts toward higher energies as ϵ increases [102]. For very heavy DPs, however, the SB approximation ceases to be valid. In this regime, the dominant production region lies deep inside the PNS, near the maximum temperature, where the transport of energy even in the trapping limit remains essentially ballistic rather than diffusive.

2. Low-energy supernova explosion

Even if DPs successfully escape the PNS, they can still interact with materials in the outer layer of the progenitor. Such interactions may deposit energy into the mantle, subsequently increasing the total kinetic energy of SN explosive ejecta [122–124].

The explosion energy of a CCSN is inferred from light-curve properties and spectroscopic analysis. While the typical range of the SN explosion energy is around 1 bethe (1 bethe = 10^{51} erg), some SN cases imply an order of magnitude lower explosion energy, on the order of 0.1 bethe [125,126]. We exploit observations of low-energy supernova explosions to constrain DPs, following Ref. [123], which pioneered this method using the axion-photon coupling as an illustrative example.

DPs deposit energy in the mantle through scatterings with materials or via their decay into visible particles. Scattering processes, which are proportional to the number density of charged particles, become significantly less efficient outside the PNS, where the density drastically decreases. Therefore, the absorption of DPs via scatterings would be inefficient in transferring their energy to the mantle. On the other hand, the decay of DPs is almost independent of the thermal properties of the outer stellar medium (e.g., electron degeneracy).

For DP masses below $2m_e$, decays occur only via the radiative three-photon channel, which is extremely suppressed and leads to lifetimes too long to contribute to SN explosion energy. On the other hand, for $m_{\gamma'} > 2m_e$, their decay into electron-positron pairs is kinematically allowed, and the lifetime can be short enough to deposit their energy into the mantle. Hence, low-energy supernova explosion observations can give a relevant constraint for $m_{\gamma'} > 2m_e$.

We follow the procedure in Ref. [127] to derive the constraint from low-energy SN explosion observations. Based on the energy-deposition framework developed in Refs. [122,123], Ref. [127] provides an improved treatment that properly accounts for the effective energy transfer from feebly interacting particles produced inside the PNS. In particular, it offers a consistent description of the strong-coupling regime, where DPs are efficiently absorbed or decay within the dense PNS medium, leading to diffusive energy transport. The energy deposited in the mantle by DP decays can be expressed as [127]

$$\mathcal{E}_{\text{mantle}} = \int dt \left[\mathcal{P}_{\text{PNS}}^{(1)} - \mathcal{P}_{\text{PNS}}^{(2)} - \mathcal{P}_{\text{prog}}^{(1)} - \mathcal{P}_{\text{prog}}^{(2)} \right]. \quad (33)$$

Here, $\mathcal{P}_{\text{PNS}}^{(1)} = \mathcal{P}_{\gamma'}|_{\text{cool}}$ represents the total DP luminosity generated inside the PNS that is available for energy deposition, while $\mathcal{P}_{\text{PNS}}^{(2)}$ correspond to the portion escaping the progenitor. The terms $\mathcal{P}_{\text{prog}}^{(1)}$ and $\mathcal{P}_{\text{prog}}^{(2)}$ denote the DP production in the mantle, with decays occurring inside the PNS and outside the progenitor, respectively.

In the weak-coupling regime, where $\mathcal{P}_{\text{PNS}}^{(i)} \gg \mathcal{P}_{\text{prog}}^{(i)}$ and the mean free path of DPs exceeds the PNS radius (so that energy transport is ballistic), the deposited energy in the mantle reads

$$\mathcal{E}_{\text{mantle}} = \int dt \int dV \frac{d\mathcal{P}_{\gamma'}}{dV} (e^{-\tau_{\text{cool}}} - e^{-\tau_{R_*}}). \quad (34)$$

The DP volume emissivity ($d\mathcal{P}_{\gamma'}/dV$), the optical depth at the boundary of the PNS ($= \tau_{\text{cool}}$) are described above, and τ_{R_*} indicates the optical depth for the progenitor surface, which is dominated by the decay probability due to the large progenitor radius. The volume integration range for DP production is limited to the PNS region ($R_{\text{PNS}} = 20$ km). For $m_{\gamma'} > 2m_e$, we find

$$\tau_{R_*} \simeq \Gamma_{e\bar{e}} \frac{R_*}{v_{\gamma'}}, \quad (35)$$

with the vacuum decay rate of the electron-positron channel (neglecting the medium effect in the mantle) as follows:

$$\Gamma_{e\bar{e}} = \epsilon^2 \frac{\alpha m_{\gamma'}^2}{3 \omega} \left(1 + \frac{2m_e^2}{m_{\gamma'}^2} \right) \sqrt{1 - 4 \frac{m_e^2}{m_{\gamma'}^2}}. \quad (36)$$

In the small- ϵ regime of our interest, DPs freely escape the PNS (i.e., $\tau_{\text{cool}} \ll 1$) and decay predominantly within the progenitor (i.e., $\tau_{R_*} \gg 1$). Equation (34) then simplifies to

$$\mathcal{E}_{\text{mantle}} \simeq \int dt \int dV \frac{d\mathcal{P}_{\gamma'}}{dV}, \quad (37)$$

corresponding to the total energy emitted in DPs from the PNS. This allows a simple estimate of the bound by comparing the DP luminosity with the neutrino luminosity at 1 sec postbounce (Sec. IV A 1).

At large couplings, where the DP absorption or decay length becomes shorter than the PNS radius, the ballistic description in Eq. (34) may no longer apply, but the energy transport from the PNS would become diffusive [127] (i.e., $\mathcal{P}_{\text{PNS}}^{(1)} \sim \mathcal{P}_{\text{prog}}^{(1)} \gg \mathcal{P}_{\text{PNS}}^{(2)}, \mathcal{P}_{\text{prog}}^{(2)}$). Such a diffusive heating relies on the DP thermal conductivity, which is set by the mean free path $\lambda_{\gamma'} \simeq v'/\Gamma_{e\bar{e}} \propto \varepsilon^{-2}$. The resulting diffusive energy deposition rate can be approximated analytically (see Ref. [127] for detailed derivation), as follows:

$$\frac{d\mathcal{E}_{\text{mantle}}}{dt} = \frac{2R_{\text{PNS}}^2}{3\pi T^2} \left(-\frac{\partial T}{\partial r} \right) \int dE_{\gamma'} \frac{\lambda_{\gamma'} E_{\gamma'}^3 p_{\gamma'} e^{E_{\gamma'}/T}}{(e^{E_{\gamma'}/T} - 1)^2}, \quad (38)$$

evaluated at the PNS surface, $r = R_{\text{PNS}}$.

However, this approximation breaks down for DP masses above 50 MeV, since it assumes diffusive DP emission at the PNS radius where $T < 10$ MeV. For heavier DPs, their emission are dominated deeper inside the PNS, where the temperature is higher, and thus even in the strong-coupling regime the transport remains effectively ballistic. In this case, Eq. (37) remains valid.

To evaluate $\mathcal{E}_{\text{mantle}}$, we use the SN simulation profile for the LS220-s18.88 model, incorporating all postbounce time snapshots, and adopt a progenitor radius of $R_* \simeq 10^{14}$ cm. The cyan region in Fig. 11 (and the corresponding “LESNe” region in Fig. 1) is excluded by the condition of $\mathcal{E}_{\text{mantle}} < 0.1$ bethe. As expected, the low-energy SN explosion provides a more stringent constraint on DPs than the cooling argument.

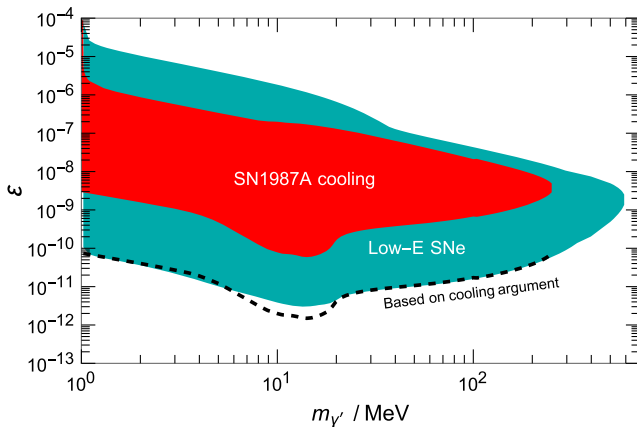


FIG. 11. Cyan-shaded region shows the constraint from low-energy supernova explosions, derived from energy deposition in the mantle due to dark photon decays. The black dashed line indicates the result obtained using the approach based on the cooling argument (red), with an effective emission duration of 2 sec.

The black curve in Fig. 11 indicates the result from the simplified approach described in Eq. (37) with the effective emission duration of 2 sec. This approximate bound agrees well with the full numerical result, except near $m_{\gamma'} \sim 10$ MeV, where transverse resonance production becomes significant. In the numerical calculation, this resonance is smeared out by accounting for the time-evolving profiles.

3. Absence of prompt SN1987A γ -ray signals

In standard core-collapse SN scenarios, any energetic photons produced within the PNS are promptly trapped and thermalized due to the extremely high density of the stellar medium. Consequently, no prompt γ -ray signal is expected to accompany the core collapse. However, new particles such as DPs, which interact only feebly with SM particles, can escape the progenitor. If they subsequently decay into photons, these events may yield detectable γ -ray signals, potentially observable with minimal delay relative to neutrino emission.

At the time when the diffusive neutrino burst from SN1987A was observed, the Gamma-Ray Spectrometer (GRS) aboard the Solar Maximum Mission satellite [128] was operational. It monitored the γ -ray flux in three energy bands: 4.1–6.4 MeV, 10–25 MeV, and 25–100 MeV [129]. The observation period extended for 223.2 sec postbounce before the satellite switched to calibration mode. No excess was detected in any energy range, leading to upper bounds on the γ -ray fluence, as follows:

$$\mathcal{F}_{\gamma} < 6.11 \text{ cm}^{-2} [4.1 - 6.4 \text{ MeV}], \quad (39)$$

$$\mathcal{F}_{\gamma} < 1.48 \text{ cm}^{-2} [10 - 25 \text{ MeV}], \quad (40)$$

$$\mathcal{F}_{\gamma} < 1.84 \text{ cm}^{-2} [25 - 100 \text{ MeV}]. \quad (41)$$

Since the region where DP-induced photons can arrive within the observation window is much closer to the progenitor than the SN1987A distance ($d_* = 51.4$, kpc), the expected γ -ray fluence from DP decays can be approximated as

$$\mathcal{F}_{\gamma} \simeq \int dr \int d \cos \theta_{\gamma} \int dE_{\gamma} \int d\omega \frac{1}{4\pi d_*^2} \frac{d\mathcal{N}_{\gamma'}}{d\omega} e^{-\Gamma_{\gamma'} r/v'} \times \frac{1}{v'} \frac{\partial^2 \Gamma_{\text{rad}}}{dE_{\gamma} d \cos \theta_{\gamma}} \Theta(t_{\text{delay}} < 223.2 \text{ sec}), \quad (42)$$

where $d\mathcal{N}_{\gamma'}/d\omega$ is the spectrum of the total number of DPs emitted from the SN, integrated over the entire PNS and all postbounce snapshots. Dividing by $4\pi d_*^2$ gives the DP fluence at Earth in the absence of decay. While $\Gamma_{\gamma'}$ is the total decay rate of DPs with the energy ω , Γ_{rad} denotes the partial decay rate into channels with a final-state photon. The variables E_{γ} and θ_{γ} are defined as the photon energy

and angle relative to the DP boost in its rest frame, respectively. The observed photon energy at Earth is then given by $E_\gamma \gamma' (1 + v' \cos \theta_\gamma)$ with $v' = \sqrt{1 - m_\gamma'^2/\omega^2}$ and $\gamma' = \omega/m_\gamma'$.

The nonvanishing DP mass and a nonstraight line of sight of radially emitted DPs (corresponding relative angle to the straight line is correlated with θ_γ) impose a delay in the arrival of photon signals from the DP decay, which reads

$$t_{\text{delay}} \simeq \frac{r}{v' \gamma'^2 (1 + v_\gamma' \cos \theta_\gamma)}. \quad (43)$$

The observational constraint, $t_{\text{delay}} < 223.2$ sec, in the last term of Eq. (42) leads to the integration limits. Furthermore, for the signals to be detectable, DPs must decay outside the photosphere, i.e., $r > R_*$.

We thus derive constraints on DPs from the nonobservation of *prompt* SN1987A γ -ray emission. For masses above the dielectron threshold, $m_\gamma' > 2m_e$, the decay $\gamma' \rightarrow e^+e^-$ is kinematically allowed, yielding a short DP lifetime with a total DP decay rate given by Eq. (36). The dominant radiative mode is final-state radiation (FSR), $\gamma' \rightarrow e^-e^+\gamma$, for which we adopt the differential rate from Appendix B of Ref. [130]. The resulting constraint on the DP parameter space, derived from the absence of an SN1987A γ -ray excess, is shown in orange in Fig. 1 and labeled as “SN1987A γ .”

Apart from SN1987A, one may also consider the γ -ray signals from Type Ic SNe, which have much more compact progenitors and therefore allow one to access different regions of parameter space by stacking many such events detected by Fermi-Large Area Telescope [131] (stacking multiple events is necessary, because if a clear timing of the onset of the explosion is not available, a single Type Ic SN cannot individually provide constraints). Furthermore, it would be interesting to study the impact of DPs on early time observations of SN explosions, such as those of SN2023ixf [132]. We leave these studies for the future.

4. Galactic positron injection

Injection of a significant number of positrons into the Milky Way can substantially enhance the diffuse γ -ray emission from the Galactic Center region. As these positrons propagate through the interstellar medium, they lose energy via Coulomb interactions with ambient electrons, emit bremsstrahlung radiation, and eventually annihilate. The annihilation of nonrelativistic positrons, predominantly through the formation of positronium (a bound state of an electron and a positron), gives rise to the observed 511 keV γ -ray line flux [133]. Precise measurements of this 511 keV line by the SPI spectrometer onboard the INTEGRAL satellite [134–137] have placed stringent limits on the Galactic nonrelativistic positron injection rate, constraining it to be no larger than $\mathcal{O}(10^{43})e^+s^{-1}$.

The average number of positrons injected into the Galaxy per CCSN, which eventually thermalize and contribute to the 511 keV emission through the decay of SN-produced dark photons outside the progenitor, can be expressed as

$$\mathcal{N}_{e^+}^{\text{SN}} = \sum_{i=\text{I,II}} r_i \int dt d\omega \frac{d\mathcal{N}_{\gamma'}}{dt d\omega} e^{-\tau_{R_*}} \int dE_{e^+} \frac{dN_{e^+}}{dE_{e^+}} P_{E_{e^+} \rightarrow m_e},$$

where $P_{E_{e^+} \rightarrow m_e}$ denotes the survival probability for a positron with initial energy E_{e^+} to become nonrelativistic through interactions with the ambient Galactic medium [138]. Here, dN_{e^+}/dE_{e^+} represents the positron spectrum from the decay of a boosted dark photon with energy ω ; the general transformation between the boosted and rest-frame spectra is presented in Appendix E.

There are two main CCSN subclasses that contribute to this rate: Type II and Type Ib/c. They are distinguished by the presence (Type II) or absence (Type Ib/c) of hydrogen features in their spectra, which reflects the difference in progenitor envelope size and composition. Accordingly, the relative weight r_i in the sum in Eq. (44) accounts for the occurrence rate of each SN type in the Galaxy.

Following Ref. [139], we adopt the upper limit on the number of positrons injected per CCSN,

$$\mathcal{N}_{e^+}^{\text{SN}} < 1.4 \times 10^{52}, \quad (44)$$

assuming the CCSN rate in the Milky Way Galaxy of 2.3 per century [140]. For analysis, we take the progenitor radii of $R_* = 5 \times 10^{12}$ cm for Type Ib/c and 10^{14} cm for Type II, with the relative fraction $r_{\text{I}} = 33\%$ and $r_{\text{II}} = 67\%$ [140], respectively. Since the core collapse and the subsequent PNS formation are essentially identical for both classes, we assume an equivalent DP production mechanism. By integrating over all the time snapshots, the yellow-shaded region in Fig. 12 (and the corresponding “511 keV line” region in Fig. 1) is excluded by the condition in Eq. (44).

Analogous to the low-energy supernova explosion constraint discussed in Sec. IVA 2, one can estimate the Galactic positron injection constraint, based on a simple rescaling of the SN cooling limits. In the regime where DPs freely escape the PNS, the constraint can be rescaled as

$$\mathcal{N}_{e^+}^{\text{SN}} \sim \sum_{i=\text{I,II}} r_i t_{\text{SN}} \left(\frac{\varepsilon}{\varepsilon_{\text{cool}}} \right)^2 \frac{L_\nu}{\langle \omega \rangle} e^{-\tau_{R_*}(\langle \omega \rangle)}, \quad (45)$$

where $\varepsilon_{\text{cool}}$ denotes the kinetic mixing corresponding to the SN1987A cooling constraint (the boundary of the red-shaded region in Fig. 12), and L_ν is the total neutrino luminosity. We take an effective duration of $t_{\text{SN}} = 2$ sec and assume a monochromatic DP spectrum with the average energy $\langle \omega \rangle$; one can evaluate this latter

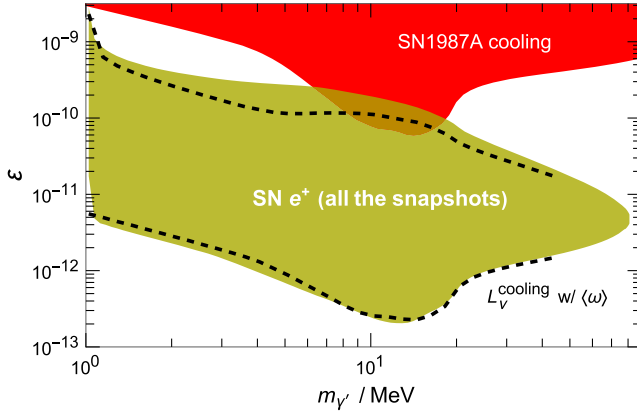


FIG. 12. Yellow-shaded region is excluded by Galactic positron injection leading to the observed 511 keV line. The black dashed lines represent the approximate estimate obtained by rescaling the SN1987A cooling constraint (red-shaded region).

numerically, but a decent approximation (especially for DP masses above ~ 20 MeV) is $\langle \omega \rangle \sim m_{\gamma'} + T$. The black dashed lines in Fig. 12 show this approximation, which agrees well with the full numerical result except around $m_{\gamma'} \sim 10$ MeV, where the broad transverse-resonance spectrum becomes relevant. For masses below 3 MeV, the production is dominated by longitudinal resonance, which imposes a sharply peaked DP spectrum centered near $\langle \omega \rangle$.

While here we have considered only the annihilation of nonrelativistic positrons, predominantly through the formation of positronium, one can also consider relativistic positrons losing energy by scattering on free and bound electrons in the interstellar medium, and then leading to a continuum emission above 511 keV [141]. This type of radiation leads to constraints (using Imaging Compton Telescope and Energetic Gamma Ray Experiment Telescope data) which are slightly better or slightly worse than the ones considered here, depending on the treatment of the background (see Fig. 9 of Ref. [142]).

5. Fireball

The fireball model was originally proposed to explain the characteristics of γ -ray burst (GRB), intense flashes of energetic photons peaking in the sub-MeV range [143]. In this model, a compact source drives a relativistic outflow that carries energy either through the kinetic energy of ultrarelativistic particles or through electromagnetic Poynting flux. The initially optically thick plasma surrounding the source is referred to as a *fireball*. As the outflow expands and becomes optically thin, its internal energy is released as γ -rays. (See Ref. [143] and references therein for a comprehensive review.)

Novel particles produced in CCSNe and decaying electromagnetically—such as the DPs considered here—can act as an “engine” for forming a fireball around the progenitor star [144]; see Refs. [145,146] for an analogous

study in the axionlike particle scenario. Following the approach of Ref. [145], we assess whether DPs could trigger fireball formation.

When DPs decay into electron–positron pairs, $\gamma' \rightarrow e^+e^-$, they inject a dense population of charged particles just outside the stellar surface. The characteristic radius at which most decays occur is $R_{\text{FB}} \simeq \frac{v_{\gamma'}}{\Gamma_{e^+e^-}}$, where $\Gamma_{e^+e^-}$ is the DP decay rate to an e^+e^- pair, and the spread in the decay radii defines the shell thickness Δ , which typically satisfies $\Delta \ll R_{\text{FB}}$. The number density of electrons (and positrons) then reads

$$n_{e^\pm} \approx \frac{\mathcal{N}_{\gamma'} e^{-R_*/R_{\text{FB}}}}{4\pi R_{\text{FB}}^2 \Delta}, \quad (46)$$

where $\mathcal{N}_{\gamma'}$ denotes the total number of DPs emitted from the PNS.

A fireball can form only if the electron–positron plasma is sufficiently dense to maintain thermal equilibrium and chemical equilibrium. The first requirement is that the optical depth to pair annihilation, $e^-e^+ \rightarrow \gamma\gamma$, be large:

$$n_{e^\pm} \sigma_{e^+e^- \rightarrow \gamma\gamma} \Delta \gg 1, \quad (47)$$

where $\sigma_{e^+e^- \rightarrow \gamma\gamma}$ is the annihilation cross section. In the comoving frame of the expanding plasma, the c.m. energy of the $e^\pm$ pairs is set by the DP mass, yielding the approximation $\sigma_{e^+e^- \rightarrow \gamma\gamma} \simeq 4\pi\alpha^2/m_{\gamma'}^2$ [147]. A second condition ensures that number-changing processes such as bremsstrahlung, $e^-e^+ \rightarrow e^-e^+\gamma$, are efficient, as follows:

$$2n_{e^\pm} \sigma_{ee \rightarrow ee\gamma} \Delta \gg 1, \quad (48)$$

with $\sigma_{ee \rightarrow ee\gamma} = 8\alpha^2[\log(m_{\gamma'}/3m_e) + \gamma_E + 5/4]/m_e^2$ [147], assuming an initial plasma temperature of order $m_{\gamma'}$. Substituting Eq. (46) into Eqs. (47) and (48) shows that both conditions are independent of Δ . The black-shaded region labeled “Fireball [Pioneer Venus Orbiter (PVO)]” in Figs. 1 and 13 indicates the portion of the DP parameter space in which either of the above conditions is satisfied and a fireball can form around the SN progenitor.

Fireball formation has important implications for interpreting several observational constraints on DPs [130]. In deriving bounds from the SN1987A γ -ray search by the GRS, we assumed that photons from final-state radiation escape freely and retain their original spectrum. If a fireball forms, however, this radiation is trapped and thermalized, and is instead released as a GRB-like signal in the sub-MeV range—outside the GRS sensitivity window. Similarly, fireball formation affects the Galactic positron-injection constraints from SN DPs: if bremsstrahlung efficiently increases the particle number [Eq. (48)], the plasma cools while maintaining energy conservation, and most positrons annihilate once the fireball becomes optically thin at $T \ll \text{MeV}$ [130,143,145,146]. As a result, the SN1987A

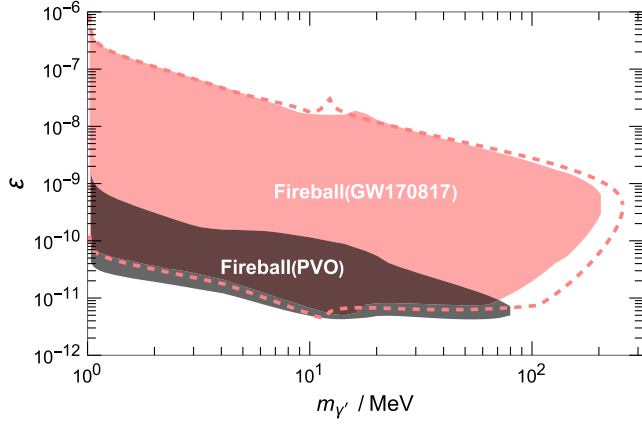


FIG. 13. Fireball constraints on the dark photon scenario from both the *Pioneer Venus Orbiter* at the time of SN1987A (dark gray-shaded region) and x-ray observations of GW170817 by Calorimetric Electron Telescope/CALET Gamma-ray Burst Monitor (CALET/CGBM), Konus-Wind, and Insight-Hard X-ray Modulation Telescope/High Energy X-ray telescope (HXMT/HE) (pink-shaded region for the LS220-s18.88 SN profile and pink dashed curve for the one-zone model).

γ -ray and Galactic positron bounds do not apply within the black-shaded “fireball” region in Figs. 1 and 13.

Beyond evading other constraints, the thermalized emission from a fireball can itself provide a probe of new particles. At the time of SN1987A, the *Pioneer Venus Orbiter* (PVO) was sensitive to photons in the 0.2–2 MeV range—consistent with the expected thermalized fireball spectrum. The nonobservation of a γ -ray excess by PVO constrains the total fireball energy to $E_{\text{FB}}^{\text{SN1987A}} \lesssim 3 \times 10^{43}$ erg [145]. In the dark gray region of Figs. 1 and 13, DPs would deposit more energy than this limit, leading to an exclusion.

A similar mechanism can occur in binary neutron-star mergers if a long-lived remnant forms before collapsing into a black hole. The event GW170817 likely produced a hypermassive neutron star surviving for ~ 1 sec [148,149]. To estimate DP emission in this context, we consider two representative thermal profiles: (i) the LS220-s18.88 SN profile at $t_{\text{pb}} = 8$ s, and (ii) a simplified one-zone model with $T = 18$ MeV, $R = 16$ km, and duration $\delta t = 1$ sec following Ref. [146]. For both, we evaluate the thermalization criteria in Eqs. (47) and (48), adopting Eq. (46) for $n_{e^\pm}$ and requiring decays occur beyond $R_{\text{min}} = 1000$ km. X-ray observations of GW170817 by CALET/CGBM, Konus-Wind, and Insight-HXMT/HE place an upper limit of $E_{\text{FB}}^{\text{GW170817}} \lesssim 3 \times 10^{46}$ erg on any GRB-like fireball emission. The corresponding exclusion based on the LS220-s18.88 profile is shown as the pink-shaded region in Figs. 1 and 13, while the pink dashed curve denotes the result for the one-zone model.

6. Diffuse γ -ray background

CCSNe have occurred isotropically throughout cosmic history, and the emission of DPs from each event

cumulatively contributes to a cosmological DP population. The contribution to the comoving DP density per redshift interval is given by

$$\frac{dn_{\gamma'}}{dz} = \frac{dn_{\text{cc}}}{dz} \int d\omega \frac{dN_{\gamma'}}{d\omega}, \quad (49)$$

where dn_{cc}/dz represents the CCSNe event rate per comoving volume. Here, we ignore any attenuation from the intergalactic medium, since the optical depth at these energies and low redshift is negligible. The subsequent radiative decay of DPs tones up the diffuse cosmic γ -ray background, which should not exceed the measured value [92]. In the following, we consider only DPs with masses below $2m_e$ and their direct radiative decay into photons. For larger masses, the DPs can decay into electron-positron pairs, with associated final state radiation. The associated bounds in this mass region, however, are already excluded by other probes and we will not include them.

For the masses of interest here (smaller than $2m_e$) DPs remain relativistic at the time of decay, and the resulting diffuse photon flux can be written as

$$\begin{aligned} \frac{d\Phi_{\gamma}}{dE_{\gamma}} &\simeq \frac{c}{4\pi} \int_0^\infty dz \frac{dn_{\text{cc}}}{dz} \int_{E_{\gamma}(1+z)}^\infty d\omega \frac{dN_{\gamma'}}{d\omega} \left(\frac{1+z}{\omega} \right) \\ &\times \left(1 - \exp \left[- \int_0^z dz' \Gamma_{\gamma'}(z') \left| \frac{dt}{dz'} \right| \right] \right) \\ &\times \left(\frac{\Gamma_{\text{rad}}}{\Gamma_{\gamma'}} \right) \frac{dN_{\gamma}}{dx} \Big|_{x=E_{\gamma}(1+z)/\omega}, \end{aligned} \quad (50)$$

where $\Gamma_{\gamma'}(z')$ is the total DP decay rate evaluated at the redshifted energy $\omega' = \omega(1+z')/(1+z)$ and Γ_{rad} is the partial decay rate for the radiative channel. The Hubble time is expressed as $|dt/dz| = 1/(1+z)H$ with the Hubble rate H , as follows:

$$H(z) = H_0 \sqrt{\Omega_{\text{m}}(1+z)^3 + \Omega_{\Lambda}}, \quad (51)$$

where H_0 is the present Hubble constant, and Ω_{m} and Ω_{Λ} are the matter and dark energy densities, respectively. We neglect photon absorption during propagation, as the CCSN rate dn_{cc}/dz is concentrated at relatively low redshifts ($z \lesssim 3$). The last term dN_{γ}/dx represents the photon energy spectrum per radiative decay of a boosted DP with the initially produced energy ω in terms of the covariant ratio of $x = E_{\gamma}(1+z)/\omega \in [0, 1]$. The explicit derivation of dN_{γ}/dx from the given energy spectrum in the rest frame of the DP, denoted by $dN_{\gamma}/d\hat{x}$ with $\hat{x} = E/(m_{\gamma}/2)$ for the photon energy E in the rest frame of the DP, is provided in Appendix E. In the highly boosted limit (i.e., $\omega \gg m_{\gamma}$), the boosted spectrum simplifies to

$$\frac{dN_{\gamma}}{dx} \simeq \int_x^1 \frac{d\hat{x}}{\hat{x}} \frac{dN_{\gamma}}{d\hat{x}}. \quad (52)$$

For masses below MeV, DPs are dominantly produced by resonant conversion of longitudinal plasmons at the energy around the plasma frequency, approximately 20 MeV, which results in DPs being relativistic during decay. Therefore, the description above is relevant with $d\mathcal{N}_{\gamma'}/d\omega \approx \mathcal{N}_{\gamma'}\delta(\omega - \omega_R)$.

In the highly boosted limit, the spectrum in the lab frame becomes

$$\frac{dN_{\gamma}}{dx} = \frac{5239 - 34300x^3 + 46575x^4 - 17514x^5}{1020}. \quad (53)$$

The factor of 3 accounts for the total number of photons per the decay. Substituting into the general expression yields the diffuse photon flux from SN-originated DPs, as follows:

$$\begin{aligned} \frac{d\Phi_{\gamma}}{dE_{\gamma}} \approx & \frac{c}{4\pi} \int_0^{\infty} dz \frac{dn_{cc}}{dz} \mathcal{N}_{\gamma'} \Theta\left(\frac{\omega_R}{1+z} - E_{\gamma}\right) \left(\frac{1+z}{\omega_R}\right) \\ & \times \left(1 - \exp\left[-\int_0^z dz' \frac{\hat{\Gamma}_{3\gamma}}{\omega'/m_{\gamma'}} \left|\frac{dt}{dz'}\right|\right]\right) \frac{dN_{\gamma}}{d\hat{x}} \Big|_{\hat{x}=E_{\gamma}(1+z)/\omega_R}, \end{aligned} \quad (54)$$

where ω_R denotes the approximately monochromatic DP energy from resonant production.

The blue-shaded region in Fig. 14 (and the corresponding “Diffuse γ (SN)” region in Fig. 1) illustrate the constraint derived from the diffuse γ -ray background. The cosmic core-collapse rate is related to the comoving star formation rate $\dot{\rho}_*(z)$, through $dn_{cc}/dz = \dot{\rho}_*(z)k_{cc}$ with the coefficient k_{cc} for core-collapse ratio per mass. For analysis, we adopt the star formation rate from Ref. [150] and take $k_{cc} = (143M_{\odot})^{-1}$ [151], which yields the most

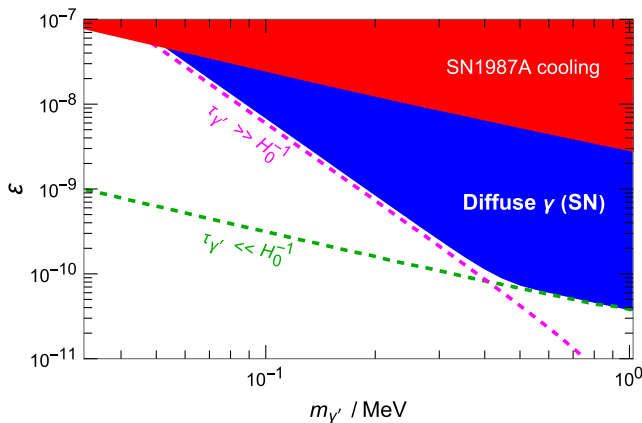


FIG. 14. Blue-shaded region is excluded by the extragalactic γ -ray background, contributed by radiative decay of dark photons from cosmic core-collapse supernovae. The magenta dashed and green dashed lines represent approximate analytical limits corresponding to lifetimes much longer and much shorter than the age of the Universe, respectively.

conservative (i.e., lowest) estimate of the cosmic core-collapse rate among those discussed in Ref. [110].

The expression in Eq. (54) can be further simplified in terms of the DP lifetime, following the approach for axions in Ref. [110]. For short-lived DPs ($\Gamma_{\gamma'} \gg H_0$), most decay shortly after production, and the exponential suppression can be neglected. Since the cosmic SN rate peaks near $z \approx 1$, we approximate

$$\begin{aligned} \left[\frac{d\Phi_{\gamma}}{dE_{\gamma}}\right]_{\Gamma_{\gamma'} \gg H_0} \approx & \frac{c}{4\pi} \left[\frac{dn_{cc}}{dz}\right]_{z=1} \mathcal{N}_{\gamma'} \Theta\left(\frac{\omega_R}{2} - E_{\gamma}\right) \\ & \times \left(\frac{2}{\omega_R}\right) \frac{dN_{\gamma}}{d\hat{x}} \Big|_{\hat{x}=2E_{\gamma}/\omega_R}. \end{aligned} \quad (55)$$

Considering the approximately measured diffuse photon flux in the γ -ray spectrum range as $E_{\gamma}^2(d\Phi_{\gamma}/dE_{\gamma}) \simeq 2 \times 10^{-3} \text{ MeV cm}^{-2} \text{ s}^{-1} \text{ ster}^{-1}$, the dashed green line in Fig. 14 shows the resulting limit in the short-lifetime regime, consistent with the full numerical calculation for $m_{\gamma'} \gtrsim 0.5 \text{ MeV}$, scaling as $\epsilon \propto m_{\gamma'}^{-1}$.

In the long-lived regime ($\Gamma_{\gamma'} \ll H_0$), most DPs survive until late times, and the photon flux is instead approximated as

$$\begin{aligned} \left[\frac{d\Phi_{\gamma}}{dE_{\gamma}}\right]_{\Gamma_{\gamma'} \ll H_0} \approx & \frac{c}{4\pi} \int_0^{\infty} dz \frac{dn_{cc}}{dz} \mathcal{N}_{\gamma'} \Theta\left(\frac{\omega_R}{1+z} - E_{\gamma}\right) \\ & \times \left(\frac{1+z}{\omega_R}\right) \int_0^z dz' \frac{\hat{\Gamma}_{3\gamma}}{\omega'/m_{\gamma'}} \left|\frac{dt}{dz'}\right| \\ & \times \frac{dN_{\gamma}}{d\hat{x}} \Big|_{\hat{x}=E_{\gamma}(1+z)/\omega_R}. \end{aligned} \quad (56)$$

The corresponding estimate is shown by the dashed magenta curve in Fig. 14, which exhibits a scaling $\epsilon \propto m_{\gamma'}^{-3}$.

In principle, for $m_{\gamma'} > 2m_e$ one could also consider constraints from photons produced via FSR in the decay $\gamma' \rightarrow e^-e^+\gamma$, similarly to the SN1987A γ -ray analyses. In the diffuse case, however, the corresponding limits are already ruled out by either SN1987A observations or by the 511 keV line constraints, so we do not show them.

7. Can dark photons prevent SN explosions?

Traditional SN1987A DP bounds come from cooling arguments: too much energy loss from the PNS via DP emission would shorten the observed neutrino burst. Moreover, for most cases, the emission of new particles is computed around ~ 1 sec after bounce and it comes from the central core, where the PNS is hot and dense enough, and the typical emission of particles such as the QCD axion is at its maximum. However, DP energy loss is rather peculiar due to the DP resonant production; thus cooling can be largest in the gain layer below the stalled shock wave

and counteract the usual shock rejuvenation by neutrino energy deposition and thus prevent the explosion [111]. In fact, in the Bethe-Wilson framework [93] of CCSN, a *gain region* exists between the gain radius and the stalled shock front, $r_{\text{gain}} < r < r_{\text{shock}}$, where neutrino absorption dominates over emission and deposits energy that helps re-energize the shock. DPs, in contrast, provide an additional energy-loss channel that drains energy from the same region, potentially suppressing shock revival if the cooling is sufficiently strong. In this regime, arguments based on subsequent PNS cooling become irrelevant: the mere requirement that the explosion occurs already leads to more restrictive and robust limits, independent of the SN1987A neutrino signal.

To investigate the impact of this cooling channel, one can compute the differential DP luminosity ($d\mathcal{P}_{\gamma'}/dr$) as a function of radius and compared it with neutrino heating and cooling at various times before shock revival [152].

The net neutrino heating rate per unit volume in the gain region can be approximated as [152]

$$Q_{\nu}^{+} \approx \frac{160 \text{ MeV}}{\text{s}} \frac{\rho}{m_u} \frac{L_{\nu_e,52}}{r_7^2 \langle \mu_{\nu} \rangle} \left(\frac{T_{\nu_e}}{4 \text{ MeV}} \right)^2, \quad (57)$$

where ρ is the matter density, m_u the atomic mass unit, $L_{\nu_e,52}$ the ν_e luminosity in units of $10^{52} \text{ erg s}^{-1}$, r_7 the radius in 10^7 cm , and $\langle \mu_{\nu} \rangle$ the neutrino flux factor approaching unity in the free-streaming limit. This estimate assumes equal luminosities and spectral shapes for ν_e and $\bar{\nu}_e$ ($L_{\nu_e} \simeq L_{\bar{\nu}_e}$), Fermi-Dirac spectra without degeneracy ($\langle \epsilon_{\nu_e}^2 \rangle \simeq 21 T_{\nu_e}^2$), and typical composition $Y_n + 2Y_p \simeq 1$. The temperature T_{ν_e} is defined at the neutrinosphere, the energy-sphere—where neutrinos decouple energetically from the background.

Along the same lines, one can derive the corresponding neutrino cooling rate per unit volume [152],

$$Q_{\nu}^{-} \approx \frac{145 \text{ MeV}}{\text{s}} \frac{\rho}{m_u} \left(\frac{T}{2 \text{ MeV}} \right)^6, \quad (58)$$

where T is the local matter temperature and ρ the density in the gain region. This expression follows from

$$Q_{\nu}^{-} = \frac{\sigma_0 T^6 (3\alpha^2 + 1)}{8\pi^2 m_e^2 m_u} [Y_p \mathcal{F}_5(\eta_e) + Y_n \mathcal{F}_5(-\eta_e)], \quad (59)$$

with $\mathcal{F}_5$ the relativistic Fermi integral; see Eq. (32) in Ref. [152]. In the second line of Eq. (58), we assumed $Y_n + Y_p \simeq 1$ and low electron degeneracy, appropriate for shock-heated layers rich in $e^{\pm}$ pairs. Unlike the heating rate, which scales as $T_{\nu_e}^2$ with the neutrino temperature at the neutrinosphere, the cooling rate depends on the sixth power of the *local* temperature. The gain radius, r_{gain} , is defined by the balance condition $Q_{\nu}^{+} = Q_{\nu}^{-}$, yielding [152]

$$r_{\text{gain},7} \left(\frac{T_g}{2 \text{ MeV}} \right)^3 \simeq 1.05 \sqrt{\frac{L_{\nu_e,52}}{\langle \mu_{\nu} \rangle_g}} \left(\frac{T_{\nu_e}}{4 \text{ MeV}} \right), \quad (60)$$

where $T_g \equiv T(r_{\text{gain}})$ is the local temperature at the gain radius.

Figure 15 illustrates the time evolution of the plasma frequency at the gain and the shock radii (upper panel) and the corresponding radial profiles of differential luminosities at 0.2 sec postbounce for the supernova explosion model LS220-s18.88; the DP parameters are fixed at $m_{\gamma'} = 0.3 \text{ MeV}$ and $\varepsilon \sim 10^{-8}$ corresponding to the supernova cooling limit. We compute the DP emissivity and compare it with the neutrino cooling and heating rates. A resonance occurs within the gain region at $r \sim 80 \text{ km}$, where DP cooling largely exceeds the effective neutrino heating. In such a case, a successful supernova explosion within the Bethe-Wilson framework can be prevented,

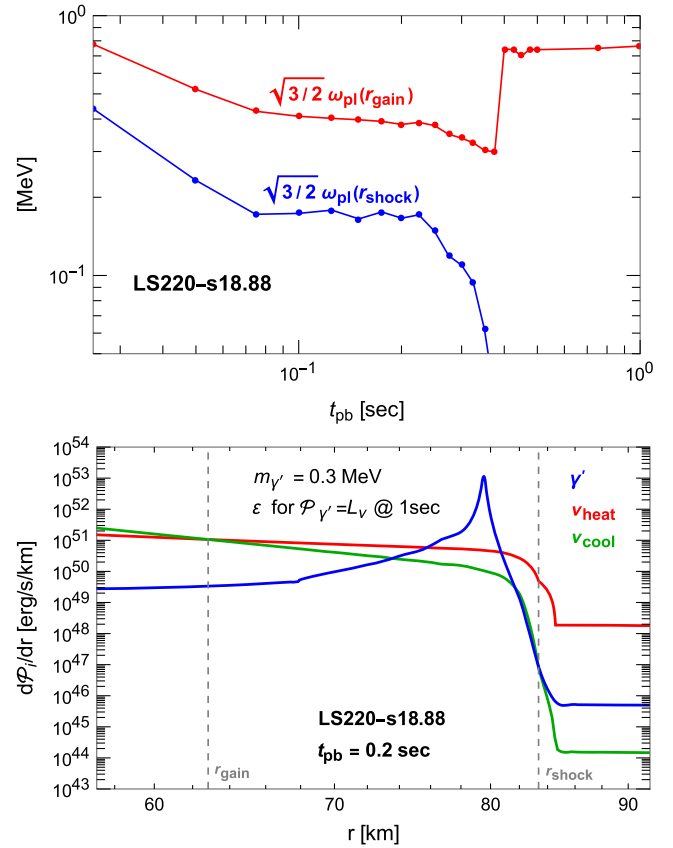


FIG. 15. Upper panel: time evolution of the plasma frequency at the gain radius (red line) and at the shock radius (blue line), multiplied by a factor of $\sqrt{3/2}$. Lower panel: differential luminosity profiles as a function of radius at 0.2 sec postbounce for dark photons (blue line), neutrino heating (red line), and neutrino cooling (green line). For the dark photon emission, we adopt $m_{\gamma'} = 0.3 \text{ MeV}$ and the corresponding kinetic mixing at the supernova cooling limit. These plots were reproduced with permission from Ref. [111].

as also confirmed for some parameters by the 2D simulations in Ref. [153]. The resulting excluded parameter space for $0.1 \text{ MeV} \lesssim m_{\gamma'} \lesssim 0.5 \text{ MeV}$ is shown in Fig. 1 (“SN Failure”), although most of it is already ruled out by cosmological probes. More interesting regions of the parameter space could be probed studying the impact on the explosion for larger DP masses, $m_{\gamma'} \gtrsim 0.5 \text{ MeV}$, for which resonant production of T modes occurs in deeper regions below r_{gain} , where it could alter the diffusive neutrino flux and affect the explosion mechanism. Such a case, however, deserves self-consistent numerical simulations, which we plan to perform in the future. Further details can be found in Ref. [111].

V. CONSTRAINTS FROM ACCELERATORS

Although this is not the focus of our work, we find it useful to include for comparison also bounds from collider probes, given that DPs can affect both precision observables and direct searches at accelerators.

One of the earliest probes arises from the anomalous magnetic moments of the electron and muon [15], since a kinetically mixed dark photon induces loop corrections proportional to $a\epsilon^2/2\pi$. For $m_{\gamma'} > 2m_e$, dark photons decay to charged fermions and depending on parameters, the decay can be prompt or displaced, requiring complementary search strategies. *Electron fixed-target and beam-dump experiments* are sensitive to dark photons in the MeV–GeV range [7,25,38,154,155]. Fixed-target setups search for narrow resonances in the invariant-mass distribution of e^+e^- pairs produced in $eZ \rightarrow eZA'$ reactions, while thick-target beam-dump experiments (e.g., E137, E141, E774 [156,157]) probe smaller mixings by exploiting long-lived dark photons decaying downstream of the dump. *Proton-beam dumps* also produce dark photons via meson decays such as $\pi^0, \eta \rightarrow \gamma A'$ [13,14], providing sensitivity to similar parameter regions. *Electron-positron colliders* (e.g., BABAR, Belle, Belle-II, KLOE) search for $e^+e^- \rightarrow \gamma A'$ followed by $A' \rightarrow f^+f^-$, using bump-hunt techniques in the invariant-mass spectrum [30]. The production cross section scales as $\sigma_{\gamma A'} \propto \epsilon^2/E_{\text{beam}}^2$, making low-energy colliders particularly competitive for sub-GeV dark photon masses. At higher energies, *proton-proton colliders* such as the LHC can probe heavier dark photons via Drell-Yan production, prompt low-mass dimuon resonances [158,159] or (more model dependent) exotic Higgs decays [29,160], while LHCb extends coverage to smaller masses through displaced-vertex searches in meson decays [161,162]. *Rare meson decays* such as $\pi^0, \eta, K, \phi \rightarrow \gamma A'$ [8] provide additional, though currently subleading, constraints. Finally, dark sectors may contain a dark Higgs boson (h') enabling associated production ($e^+e^- \rightarrow A'h'$), leading to distinctive multilepton or displaced signatures studied at BABAR [30], Belle [163], KLOE-2 [164], and Belle-II [165].

In Fig. 1, the gray-shaded region shows the compilation of constraints obtained using the DarkCast tool [35–37] and Ref. [38], as well as forecasts for Search for Hidden Particles facility at CERN [38,166]. For the masses and couplings of interest for this work, the relevant bounds come from beam-dump experiments.

VI. DISCUSSIONS AND CONCLUSION

In this work, we have re-examined the *irreducible* constraints on DPs. On the cosmological side, we fixed the reheating temperature of the Universe to $T_{\text{RH}} = 6 \text{ MeV}$, which corresponds to the lowest value still compatible with BBN according to the most recent evaluation of light-element abundances and N_{eff} [52]. Within this cosmological setup, we consistently derived all related observables, including diffuse and Galactic γ -ray emission, CMB anisotropies, variations of N_{eff} , and the impact on primordial element abundances. After addressing these cosmological implications, we turned to the production of DPs in CCSN environments and derived the associated constraints, such as those from neutrino cooling, prompt and diffuse γ -ray emission, fireball formation, Galactic positron injection, low-energy supernova energetics, and the possible stalling of the SN shock wave. We also surveyed current collider limits, which provide a complementary probe to the cosmological and astrophysical bounds, particularly in the regime of large DP masses and sizable kinetic mixing. Taken together, these results delineate the most robust and model-independent exclusion region for dark photons over many orders of magnitude in mass, highlighting the complementarity between laboratory, astrophysical, and cosmological probes.

Finally, we recall that our constraints apply to the minimal DP model, and that any coupling to a dark sector [167] may alter some of the bounds discussed. For example, if the DP can decay invisibly into dark sector particles, all limits from fireball formation, γ -rays, 511 keV line, LESNe must be rescaled. Similarly, cosmological bounds assume that production is driven by freeze-in via kinetic mixing with the Standard Model; however, the relic abundance can be significantly modified in the presence of a richer dark sector. We leave the study of these different scenarios for future work.

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Views and opinions expressed are however those of the authors only and do not necessarily reflect those of the European Union. The European Union cannot be held responsible for them.

DATA AVAILABILITY

The data that support the findings of this article are openly available [51,62,67,72–74,77,90–92,112,120].

APPENDIX A: PROPERTIES OF THE SM PLASMA

In this appendix, we summarize well-known results about the SM plasma. We focus on the scenario in which a maximal temperature is as low as $\mathcal{O}(1)$ MeV, as required to ensure the consistency of the standard cosmology; in this paper, we take 6 MeV. Thus, the light degrees in the SM thermal bath contain photons, electrons, and neutrinos. Although there exists a relatively small baryon asymmetry in the form of nucleons, its impact on our discussion is negligible due to its massiveness. Accordingly, the real parts of $\pi_{T,L}$ are mainly contributed by the coherent scatterings with surrounding electrons/positrons and read [168,169]

$$\text{Re } \pi_T = \omega_{\text{pl}}^2 \left[1 + \frac{1}{2} G(v_*^2 k^2 / \omega^2) \right], \quad (\text{A1})$$

$$\text{Re } \pi_L = \omega_{\text{pl}}^2 \frac{m_\gamma^2}{\omega^2} \frac{1 - G(v_*^2 k^2 / \omega^2)}{1 - v_*^2 k^2 / \omega^2}, \quad (\text{A2})$$

where ω and k denote the energy and momentum of the external DP state, respectively. The plasma frequency ω_{pl} is given by

$$\omega_{\text{pl}}^2 = \sum_{i=e^\pm} \frac{4\alpha}{\pi} \int_0^\infty dp f_i \frac{p^2}{E} \left(1 - \frac{1}{3} v^2 \right), \quad (\text{A3})$$

where f_i denotes the phase space distribution of i with their energy E and momentum p , $v = p/E$, and v_* is the typical electron velocity given by

$$v_*^2 = \frac{\int_0^\infty dp f_p p (5v^3/3 - v^5)}{\int_0^\infty dp f_p p (v - v^3/3)}. \quad (\text{A4})$$

The function G is defined as

$$G(x) = \frac{3}{x} \left[1 - \frac{2x}{3} - \frac{1-x}{2\sqrt{x}} \log \frac{1+\sqrt{x}}{1-\sqrt{x}} \right]. \quad (\text{A5})$$

APPENDIX B: THERMAL PRODUCTION IN THE EARLY UNIVERSE

In this appendix, we provide details about resonant and nonresonant production of DPs.

1. Resonant production

We start providing the analytic derivation of the thermal production of DPs at the resonance point. The following derivation is valid for narrow resonance width cases, which are typically satisfied.

a. Transverse modes

For the transverse polarization, $\text{Re } \pi_T$ is of the order of ω_{pl}^2 across the entire momentum range. Since the typical energy range of produced DPs is an order of temperature, which is larger than ω_{pl} , we can employ the relativistic approximation for the energy-momentum of DPs (i.e., $\omega \simeq k$) leading to

$$\text{Re } \pi_T \simeq \omega_{\text{pl}}^2 \left[1 + \frac{1}{2} G(v_*^2) \right]. \quad (\text{B1})$$

Although each momentum of transverse DPs may experience resonance at different epochs in principle, the time difference within the dominant energy range of temperature is shorter than the Hubble scale; taking momentum-independent $\text{Re } \pi_T$ in Eq. (B1) is a valid simplification. Therefore, the resonance of the transverse polarization occurs approximately at the specific temperature T_{res} . Plugging Eq. (5) with $\text{Re } \pi_T$ in Eq. (B1) into the Boltzmann equation in Eq. (7), we can derive

$$\begin{aligned} [Y_{\gamma'}]_{T,L}^{\text{res}} &\simeq \frac{\varepsilon^2 m_{\gamma'}^4}{\pi^2} \int dT \frac{1}{sTH} \left(1 - \frac{1}{3} \frac{d \log g_{*s}}{d \log T^{-1}} \right) \\ &\times \frac{|\text{Im } \pi_T|}{(\partial \text{Re } \pi_T / \partial T)^2 (T - T_{\text{res}})^2 + |\text{Im } \pi_T|^2} \\ &\times \int d\omega \frac{\sqrt{\omega^2 - m_{\gamma'}^2}}{e^{\omega/T} - 1} \\ &\simeq \varepsilon^2 m_{\gamma'}^2 \frac{\pi^3}{12\zeta(3)} \left[\frac{n_\gamma}{s} \frac{(\partial \text{Re } \pi_T / \partial T)^{-1}}{T/m_{\gamma'}^2} \right. \\ &\times \left. \frac{1}{TH} \left(1 - \frac{1}{3} \frac{d \log g_{*s}}{d \log T^{-1}} \right) \right]_{T=T_{\text{res}}}. \end{aligned} \quad (\text{B2})$$

Here, we approximate $((\partial \text{Re } \pi_T / \partial T)^2 (T - T_{\text{res}})^2 + |\text{Im } \pi_T|^2)^{-1} \simeq \delta(T - T_{\text{res}}) \times \pi (\partial \text{Re } \pi_T / \partial T)^{-1} / |\text{Im } \pi_T|$, which is valid within the narrow resonance width condition $\text{Im } \pi_T \ll \text{Re } \pi_T$. Moreover, we neglect $m_{\gamma'}$ for integral over ω due to $\omega_{\text{pl}} \ll T$, enabling it to be expressed by the photon equilibrium number density n_γ .

b. Longitudinal modes

The resonance feature for the longitudinal polarization differs from the case of the transverse polarization. According to Eq. (A2), the real part of the longitudinal polarization can be approximated as

$$\text{Re } \pi_L \simeq m_{\gamma'}^2 \frac{\tilde{\omega}^2}{\omega^2}, \quad (\text{B3})$$

where $\tilde{\omega}^2 = \mathcal{O}(\omega_{\text{pl}}^2)$.² This indicates that the resonance for the longitudinal polarization, $\text{Re } \pi_L = m_{\gamma'}^2$, occurs at the specific energy $\omega = \tilde{\omega}$ at any temperature, as long as $m_{\gamma'} > \omega_{\text{pl}}$. Following the transverse case above, the contribution to the DP abundance from the resonance of the longitudinal polarization reads

$$\begin{aligned} [Y_{\gamma'}]_{\text{L}}^{\text{res}} &\simeq \frac{\epsilon^2 m_{\gamma'}^2}{2\pi^2} \int dT \frac{1}{sTH} \left(1 - \frac{1}{3} \frac{d \log g_{*s}}{d \log T^{-1}} \right) \\ &\times \int d\omega \frac{\sqrt{\omega^2 - m_{\gamma'}^2}}{e^{\omega/T} - 1} \frac{|\text{Im } \pi_L / m_{\gamma'}^2|}{(2/\tilde{\omega})^2 (\omega - \tilde{\omega})^2 + |\text{Im } \pi_L / m_{\gamma'}^2|^2} \\ &\simeq \frac{\epsilon^2 m_{\gamma'}^2}{2\pi} \int dT \frac{1}{sTH} \left(1 - \frac{1}{3} \frac{d \log g_{*s}}{d \log T^{-1}} \right) \frac{\tilde{\omega} \sqrt{\tilde{\omega}^2 - m_{\gamma'}^2}}{2(e^{\tilde{\omega}/T} - 1)}, \end{aligned} \quad (\text{B4})$$

where we approximate $|\text{Im } \pi_L / m_{\gamma'}^2| / ((2/\tilde{\omega})^2 (\omega - \tilde{\omega})^2 + |\text{Im } \pi_L / m_{\gamma'}^2|^2) \simeq \delta(\omega - \tilde{\omega}) \times \pi (\tilde{\omega}/2)$.

Figure 16 illustrates $[Y_{\gamma'}]_{\text{T,L}}^{\text{res}}$ for the transverse (solid) and longitudinal (dashed) polarizations.

2. Nonresonant production

We provide the DP production rate for the following processes: $e^\mp + \gamma \rightarrow e^\mp + \gamma'$ (Compton-like scatterings), $e^- + e^+ \rightarrow \gamma + \gamma'$ (annihilation), and $e^- + e^+ \rightarrow \gamma'$ (coalescence). The following expressions and derivations can also be straightforwardly generalized to other charged particles as scattering targets.

In the following derivations, we neglect plasma effects, as they are irrelevant under the conditions where these processes become important. Although this assumption is

²When electrons/positrons are relativistic (i.e., $T \gg m_e$), $\tilde{\omega}^2$ is an order of ω_{pl}^2 increased by a factor $\log[\omega^2/m_{\gamma'}^2]$, which is never a large number for the cases of interest.

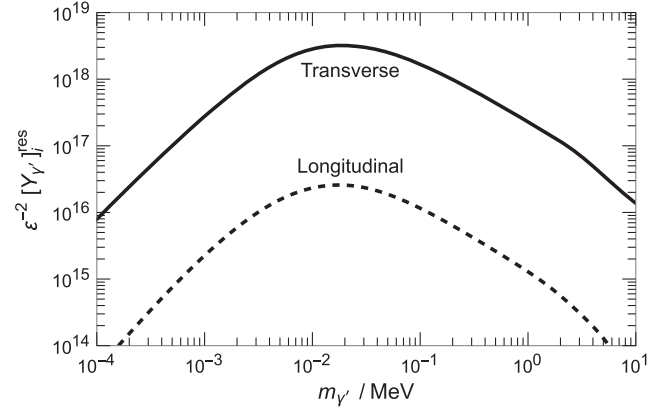


FIG. 16. Freeze-in dark photon abundance partially contributed from the resonance, $[Y_{\gamma'}]_i^{\text{res}}$, for i polarization.

strictly valid only for temperatures satisfying $\omega_{\text{pl}} < m_{\gamma'}$, the IR-dominated nature of DP production renders this approximation reliable for our purposes (see the discussion in the main text). In this limit, plasma effects can be ignored, and the polarization sum simplifies to

$$\sum_{\text{T,L}} \epsilon_{\gamma'}^\mu \epsilon_{\gamma'}^{*\nu} \rightarrow -g^{\mu\nu} + \frac{k^\mu k^\nu}{m_{\gamma'}^2}, \quad (\text{B5})$$

where the contribution of the last term vanishes due to electromagnetic current conservation.

For a generic $2 \rightarrow 2$ process, $1 + 2 \rightarrow 3 + \gamma'$, the DP production rate is given by

$$\begin{aligned} \gamma_{\gamma'} &= \int \prod_{i=1,2,3} \frac{g_i d^3 \vec{p}_i}{(2\pi)^3 2E_i} \frac{g_{\gamma'} d^3 \vec{k}}{(2\pi)^3 2\omega} f_1 f_2 (1 \pm f_3) \\ &\times |\mathcal{M}|^2 (2\pi)^4 \delta^4(p_1 + p_2 - p_3 - p_{\gamma'}), \end{aligned} \quad (\text{B6})$$

where $p_i = (E_i, \vec{p}_i)$, f_i is the distribution function and g_i denotes the internal degrees of freedom of species i . Here, $p_{\gamma'} = (\omega, \vec{k})$ is the four-momentum of the DP, and $\mathcal{M}$ is the corresponding matrix element. The $\pm$ sign accounts for Bose enhancement or Pauli blocking in the final state, though these effects are negligible in the analysis below.

Since plasma corrections are neglected, the phase-space integration can be evaluated in the c.m. frame following the standard method in Ref. [50]. Assuming the Boltzmann statistics, $f_i = e^{-E_i/T}$, the DP production rate can be written by

$$\gamma_{\gamma'} = \frac{g_1 g_2}{32\pi^4} T \int_{s_{\text{min}}}^{\infty} ds \frac{\lambda(s, m_1, m_2)}{\sqrt{s}} K_1\left(\frac{\sqrt{s}}{T}\right) \sigma_{12 \rightarrow 3\gamma'}(s) \quad (\text{B7})$$

where $\sigma_{12 \rightarrow 3\gamma'}(s)$ is the corresponding cross section averaged over the angular dependence in terms of the squared c.m. energy s with the integral range of

$s_{\min} = \text{Max}[(m_1 + m_2)^2, (m_3 + m_{\gamma'})^2]$, and the Källén function λ is defined as

$$\lambda(x, y, z) = (x - (y + z)^2)(x - (y - z)^2). \quad (\text{B8})$$

The corresponding differential cross section reads

$$\frac{d\sigma_{12 \rightarrow 3\gamma'}}{d\Omega} = \frac{1}{64g_i g_j \pi^2 s} \sqrt{\frac{\lambda(s, m_3, m_{\gamma'})}{\lambda(s, m_1, m_2)}} |\mathcal{M}|^2. \quad (\text{B9})$$

For the Compton-like scattering process $e^\mp + \gamma \rightarrow e^\mp + \gamma'$, the cross section is given by

$$\begin{aligned} \sigma_{\text{sc}} = & \frac{\epsilon^2 e^4}{16\pi s} \frac{\sqrt{\lambda(s, m_e, m_{\gamma'})}}{s(s - m_e^2)^3} \left(s^3 + (15m_e^2 + 7m_{\gamma'}^2)s^2 \right. \\ & + m_e^2(2m_{\gamma'}^2 - m_e^2)s + m_e^6 - m_e^4 m_{\gamma'}^2 \\ & - (s^2 - (6m_e^2 + 2m_{\gamma'}^2)s - 3m_e^4 + 2m_e^2 m_{\gamma'}^2 + 2m_{\gamma'}^4) \\ & \left. \times \frac{4s^2}{\sqrt{\lambda(s, m_e, m_{\gamma'})}} \text{Coth}^{-1} \left[\frac{m_{\gamma'}^2 - m_e^2 - s}{\sqrt{\lambda(s, m_e, m_{\gamma'})}} \right] \right). \quad (\text{B10}) \end{aligned}$$

The cross section for the electron-positron annihilation process reads

$$\begin{aligned} \sigma_{\text{ann}} = & \frac{\epsilon^2 e^4}{4\pi s^{3/2}} \frac{1}{\sqrt{s - 4m_e^2}(s - m_{\gamma'}^2)} \left(-s^2 - 4m_e^2 s - m_{\gamma'}^4 \right. \\ & + \frac{2\sqrt{s}(s^2 + 4m_e^2 s - 8m_e^4 - 4m_e^2 m_{\gamma'}^2 + m_{\gamma'}^4)}{\sqrt{s - 4m_e^2}} \\ & \left. \times \text{Tanh}^{-1} \left[\sqrt{\frac{s - 4m_e^2}{s}} \right] \right). \quad (\text{B11}) \end{aligned}$$

Note that the factor $(s - m_{\gamma'}^2)^{-1}$ of the denominator in the production rate may cause IR divergence if $m_{\gamma'} > 2m_e$. Following the prescription in Ref. [45], this can be regulated by replacing

$$(s - m_{\gamma'}^2)^{-1} \rightarrow (s - m_{\gamma'}^2 + 2m_{\gamma'} m_\gamma(T))^{-1}, \quad (\text{B12})$$

where $m_\gamma(T)$ is the thermal photon mass from electron-photon plasma and is approximately given by $m_\gamma^2 \approx \omega_{\text{pl}}^2$.

When kinematically allowed, inverse-decay (i.e., coalescence) processes, $e^- + e^+ \rightarrow \gamma'$, can also contribute significantly to DP production. Using detailed balance and following Ref. [50], the corresponding reaction density can be expressed in terms of the *equilibrium* DP phase space and its vacuum decay width as

$$\gamma_{\gamma'}(T) = \frac{g_{\gamma'} m_{\gamma'}^2 T}{2\pi^2} K_1\left(\frac{m_{\gamma'}}{T}\right) \Gamma_{e^+ e^-}, \quad (\text{B13})$$

where $g_{\gamma'} = 3$ for a massive vector, K_1 is the modified Bessel function, and the decay width at the rest frame reads

$$\Gamma_{e^+ e^-} = \frac{\alpha \epsilon^2}{3} m_{\gamma'} \left(1 + \frac{2m_e^2}{m_{\gamma'}^2} \right) \sqrt{1 - \frac{4m_e^2}{m_{\gamma'}^2}}. \quad (\text{B14})$$

Here, we provide a useful analytical result for the DP abundance from electron-positron coalescence; for details, see Ref. [170]. Performing the Boltzmann equation in Eq. (7) with defining $x \equiv m_{\gamma'}/T$, we obtain

$$\begin{aligned} Y_{\gamma'} & \simeq \int_{T_0}^{T_{\text{RH}}} \frac{\gamma_{\gamma'}}{sHT} dT \\ & = \frac{45}{(1.660)4\pi^4} \frac{g_{\gamma'} M_{\text{Pl}} \Gamma_{e^+ e^-}}{g_{*s} g_*^{1/2} m_{\gamma'}^2} \int_{x_{\text{RH}}}^{\infty} x^3 K_1(x) dx, \quad (\text{B15}) \end{aligned}$$

with $x_{\text{RH}} \equiv m_{\gamma'}/T_{\text{RH}}$. Since contributions at $T \ll m_{\gamma'}$ are exponentially suppressed, the lower integration limit T_0 is irrelevant. It is convenient to factor out the full-range integral $\int_0^\infty x^3 K_1(x) dx = 3\pi/2$ and define

$$S(x) \equiv \frac{2}{3\pi} \int_x^\infty u^3 K_1(u) du, \quad S(0) = 1. \quad (\text{B16})$$

Then, the yield takes the compact form

$$Y_{\gamma'} = \frac{135}{1.6608\pi^3} \frac{g_{\gamma'} M_{\text{Pl}} \Gamma_{e^+ e^-}}{g_{*s} g_*^{1/2} m_{\gamma'}^2} S\left(\frac{m_{\gamma'}}{T_{\text{RH}}}\right). \quad (\text{B17})$$

For $x \gg 1$, where $K_1(x) \simeq \sqrt{\pi/(2x)} e^{-x}$, one finds

$$S(x) \simeq \frac{2}{3\sqrt{2\pi}} x^{5/2} e^{-x} (x \gg 1), \quad (\text{B18})$$

which makes the expected Boltzmann suppression explicit when $m_{\gamma'} \gg T_{\text{RH}}$. A smooth analytic fit valid for all x that reproduces $S(0) = 1$ and Eq. (B18) is

$$S(x) \simeq [0.266 + 0.734e^{-x}]e^{-x}(1 + x^2)^{5/4}. \quad (\text{B19})$$

APPENDIX C: HIGHER INITIAL TEMPERATURE LIMIT

We extend the analysis presented in the main text, where we considered the minimal cosmological initial temperature of $T_i = 6$ MeV, to explore higher-temperature scenarios as a more general case. As an illustrative example, we take $T_i = 100$ MeV and examine how the resulting cosmological bounds on DPs are modified.

Since the DP production rate governed by renormalizable interactions is IR-dominated, the contribution of nonresonant scattering processes (see Appendix B 2) to

the thermal freeze-in abundance is largely insensitive to T_i when the DP mass is below this temperature. Consequently, as shown in Fig. 3, for $m_{\gamma'} \lesssim 0.74$ MeV corresponding to $\sqrt{3/2}\omega_{\text{pl}}(6 \text{ MeV})$, the thermal DP abundance is essentially identical for both $T_i = 6$ MeV (black solid line) and $T_i = 100$ MeV (gray dashed line). At higher masses, however, the abundance increases due to additional production channels. Resonant conversion from photons occurs up to $\sqrt{3/2}\omega_{\text{pl}}(10 \text{ MeV}) \simeq 12$ MeV; while this mechanism dominates for subthreshold masses as $m_{\gamma'} < 2m_e$, its contribution becomes subdominant (about 10% of the total) for heavier DPs, where kinematically allowed coalescence processes begin to dominate. At even higher masses, additional coalescence channels involving both leptons and hadrons become kinematically accessible. Their contribution to the total DP production rate can be expressed using Eq. (B13), with the relevant decay widths for each channel. For leptonic coalescence processes such as $\mu^- \mu^+ \rightarrow \gamma'$, the decay width follows Eq. (B14) with the appropriate lepton mass. For hadronic processes like $\pi^- \pi^+ \rightarrow \gamma'$, the decay width reads [76]

$$\Gamma_{\pi\pi} = \frac{1}{12} \alpha_{\text{eff}}^{\pi\pi} \epsilon^2 m_{\gamma'} \left(1 - 4 \frac{m_\pi^2}{m_{\gamma'}^2} \right)^{3/2}, \quad (\text{C1})$$

where $\alpha_{\text{eff}}^{\pi\pi}$ denotes the effective scalar QED coupling [76,171]. We adopt the empirical fitting functions for $\alpha_{\text{eff}}^{\pi\pi}$ in the cases of the charged pions [171] and the charged kaons [172], based on measurements of the $e^- e^+ \rightarrow \pi^- \pi^+$ and $e^- e^+ \rightarrow K^- K^+$ from the BABAR Collaboration.

Building upon the derivation of the thermal DP abundance, we perform the same cosmological analyses as outlined in Sec. III. The results for $T_i = 100$ MeV are illustrated in Fig. 17. For $m_{\gamma'} < 2m_e$, the constraints from overproduction (gray-shaded region) and Galactic/extragalactic photon observations (red- and green-shaded regions, respectively) remain nearly identical to those in the $T_i = 6$ MeV case, except for an enhancement in the range $m_{\gamma'} \gtrsim 0.74$ MeV due to additional resonant contributions to the DP abundance. For intermediate masses, $2m_e \lesssim m_{\gamma'} \lesssim 100$ MeV, the bounds from ΔN_{eff} , CMB energy injection, and BBN (mainly due to deuterium dissociation) are also comparable to those obtained for $T_i = 6$ MeV.

At higher masses, $m_{\gamma'} \gtrsim 100$ MeV, the CMB constraint extends up to $m_{\gamma'} \sim 300$ MeV, while new bounds emerge from hadronic effects during BBN once $m_{\gamma'}$ exceeds the pion threshold, $2m_\pi$. These additional BBN constraints originate from the injection of mesons (e.g., $\pi^\pm$) produced in DP decays. Such mesons can alter the neutron-proton interconversion rates via reactions such as $p + \pi^- \leftrightarrow n + \pi^0$ and $n + \pi^+ \leftrightarrow p + \pi^0$, thus modifying the neutron freeze-out fraction prior to nucleosynthesis [173]. Since charged pions have short lifetimes ($\tau_\pi \sim 10^{-8}$ s) and are rapidly

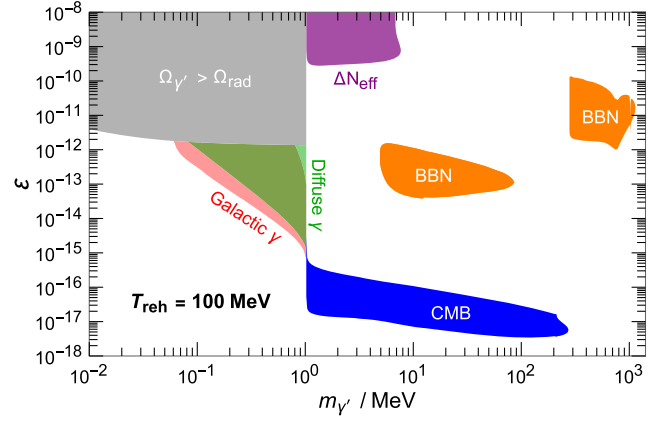


FIG. 17. Constraints on the dark photon model with $T_{\text{reh}} = 100$ MeV. The blue-shaded region is excluded by CMB. The orange-shaded regions correspond to BBN. The purple-shaded region shows ΔN_{eff} constraints. The gray-shaded region is excluded by the overproduction of dark photons. The red- and green-shaded region is ruled out by Galactic/extragalactic x-ray flux.

thermalized in the primordial plasma, their number density can be approximated by a quasistatic equilibrium form of $n_{\pi^\pm} \approx \tau_\pi \Gamma_{\gamma'} n_{\gamma'}$. The resulting modification of the neutron-proton ratio can be evaluated by incorporating this meson population into the Boltzmann equations governing weak interactions; for details, see Refs. [173,174]. This alters the freeze-out neutron fraction and consequently the predicted ^4He abundance, leading to the BBN constraint shown as the orange-shaded region in the high-mass regime of Fig. 17.

In Fig. 18, we compare our updated BBN constraint (orange-shaded regions) with the earlier result from Ref. [76] (brown dashed lines). For lower masses, where

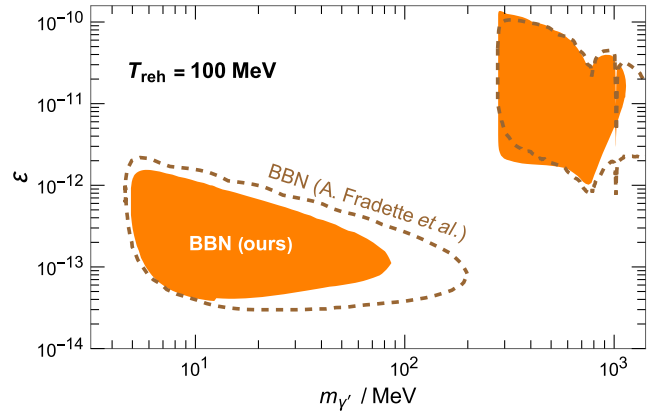


FIG. 18. BBN constraints for an initial temperature of $T_{\text{reh}} = 100$ MeV (orange-shaded regions) compared with the previous BBN result from Ref. [76] (brown dashed lines). We have verified that, when adopting the same values for the measured D/H ratio and nuclear reaction rates as in Ref. [76], our results below the GeV scale agree very well, except for a 10% discrepancy of unknown origin around $m_{\gamma'} \sim 100$ MeV.

reduction of the relic deuterium abundance is induced by the radiative DP decay, our bound is slightly weaker due to the inclusion of sizable theoretical uncertainties in the nuclear reaction rates affecting deuterium synthesis. At higher masses, where hadronic energy injection occurs, our result is in good agreement with Ref. [76] up to 0.8 GeV, above which the thermal DP abundance is exponentially suppressed. We have verified numerically that, when we adopt the same values for the measured D/H ratio and nuclear reaction rates as in Ref. [76], our results below the GeV scale agree, except for a 10% discrepancy of unknown origin for $m_{\gamma'} \sim 100$ MeV.

APPENDIX D: COMPTON-LIKE SCATTERING RATE IN THE SUPERNOVA ENVIRONMENT

We compute the Compton-like scattering rate for DPs, $e^- + \gamma' \rightarrow e^- + \gamma$, which contributes to both DP absorption and, by detailed balance, their inverse production process, under conditions relevant to core-collapse supernovae. As discussed in Sec. II, DP interactions follow the same structure as electromagnetic interactions, up to an overall factor of the kinetic mixing parameter ϵ . The detailed balance relation connects the absorption and emission rates [see Eq. (4)], allowing the Compton rate for photons to be reinterpreted. Hence, we derive the Compton scattering rate for photons involving the dispersion of DPs, and one can easily apply these results to evaluate the DP interaction rate in CCSNe. In what follows, we focus on the transverse polarization modes; the longitudinal mode contribution can be obtained by multiplying the transverse result by a factor of $m_{\gamma'}^2/\omega^2$, which follows from current conservation.

The general expression for the absorption via Compton-like scatterings reads

$$\begin{aligned} \Gamma_{\text{sc}}(\omega) = & \frac{2}{2\omega} \int \frac{d^3 \vec{p}_i}{(2\pi)^3 2E_i} f_i \int \frac{d^3 \vec{p}_f}{(2\pi)^3 2E_f} (1 - f_f) \\ & \times \int \frac{d^3 \vec{p}_\gamma}{(2\pi)^3 2E_\gamma} |\mathcal{M}|^2 (2\pi)^4 \delta^4(k + p_i - p_f - p_\gamma), \end{aligned} \quad (\text{D1})$$

where $p_{i(f)} = (E_{i(f)}, \vec{p}_{i(f)})$ denotes the four-momentum of the incident (scattered) electron, k is the four-momentum of the incoming photon with the DP dispersion, and $p_\gamma = (E_\gamma, \vec{p}_\gamma)$ is that of the outgoing photon state. For the final photon state, we neglect the statistical factor and the plasma correction (i.e., $E_\gamma \simeq |\vec{p}_\gamma|$). In the relativistic regime, the leading-order squared matrix element, averaged over initial spin degrees of freedom, reads

$$|\mathcal{M}|^2 = 32\pi^2 \alpha^2 \left(\frac{p_i \cdot k}{p_i \cdot p_\gamma} \right). \quad (\text{D2})$$

Although a full analytic expression valid across all regimes is cumbersome, useful approximate results can be obtained for conditions characteristic of CCSNe. Below we discuss two representative limits relevant for DP absorption and production.

DPs produced in the PNS may be reabsorbed in the outer stellar layers through Compton-like scattering, $\gamma' + e^- \rightarrow \gamma + e^-$, since the bremsstrahlung process, being proportional to the square of the target density, becomes inefficient in the dilute outer regions. In this case, the relevant DP energies are comparable to or higher than the PNS temperature of 30–50 MeV. At the outer PNS boundary ($r \gtrsim 20$ km), the temperature decreases significantly, while the matter remains degenerate ($\mu_e \gg T$). For the parameter space of interest, we typically have $\omega \gtrsim \mu_e$ in this case of absorption outside the PNS.

In this degenerate limit, the Fermi-Dirac distribution may be approximated by a Heaviside step function, implying that electrons are filled up to μ_e , and scattering is only allowed for final electrons with $E_f > \mu_e$. Performing the phase-space integration under these conditions yields

$$\Gamma_{\text{sc}} \approx \frac{\alpha^2 \mu^2}{2\pi \omega} \log \left(\frac{\mu \omega}{m_e^2} \right). \quad (\text{D3})$$

This expression provides a reliable estimate of the Compton-like absorption rate for DPs escaping the PNS core and is the dominant absorption mechanism outside the neutrino sphere.

DP production in the outer supernova layers can also impact the explosion dynamics, potentially preventing successful shock revival (see Sec. IV A 7). Analogous to the absorption discussed above, the DP production outside the PNS is governed by Compton-like scatterings. In these regions, such as the gain region, the electron density is much lower, and the temperature exceeds the chemical potential ($T \gg \mu_e$), corresponding to the classical limit. Pauli blocking can then be neglected, and the rate can be expressed as

$$\Gamma_{\gamma'} = 2 \int \frac{d^3 p_i}{(2\pi)^3} f_i \sigma_{\text{sc}} v_{\text{mol}}, \quad (\text{D4})$$

where $v_{\text{mol}} = (s - m_e^2)/2\omega E_i$ is the Møller velocity, and the cross section takes the well-known form

$$\sigma_{\text{sc}} = \frac{2\pi \alpha^2}{s} \log \left(\frac{s}{m_e^2} \right). \quad (\text{D5})$$

Assuming the Maxwell-Boltzmann for electrons as a good approximation in this case results in

$$\Gamma_{\text{sc}} \approx \frac{\alpha^2 T^2}{\pi \omega} \log \left(\frac{\omega T}{m_e^2} \right). \quad (\text{D6})$$

Note that Eq. (D6), derived under the assumptions $T \gg \mu_e$ and $\omega \sim T$, remains approximately valid even in the degenerate regime ($\mu_e \gg T$) as long as $T \sim \omega$, up to a factor of about 2. This is because only electrons near the Fermi surface, within a narrow energy width $\sim T$, participate effectively in the scattering. The dominant contribution arises from backward scatterings in the c.m. frame, where the outgoing electron acquires a higher energy, making Pauli suppression less. In this regime, however, the Compton-like process typically becomes subdominant compared to bremsstrahlung, owing to the higher target density dependence of the latter. Finally, we stress that for all the results of this paper, we have used the complete, fully numerical production and absorption rates for DPs. However, the analytical formulas reported here agree extremely well with the numerics across all the parameter space of interest.

APPENDIX E: SPECTRUM OF THE DECAY OF DARK PHOTONS IN THE BOOSTED FRAME

We provide the energy spectrum of daughter particles from the decay of boosted dark photons, in terms of the spectrum in the rest frame.

Let us consider the dark photon decay channel into degenerate particles; for instance, $\gamma' \rightarrow 3\gamma$ or $\gamma' \rightarrow e^-e^+$. Denoting the decay rate at the rest frame, $\Gamma_{\gamma'}$, the spectrum of a daughter particle i in the boosted dark photon frame with the energy ω is given by

$$\left. \frac{dN_i}{dE} \right|_{E=\omega} = \frac{N_i}{2\sqrt{\gamma'^2 - 1}} \int_{E_0^{\min}}^{E_0^{\max}} dE_0 \frac{1}{\sqrt{E_0^2 - m^2}} \times \frac{1}{\Gamma_{\gamma'}} \frac{d\Gamma_{\gamma'}}{dE_0} \Theta\left(\frac{m_{\gamma'}}{2} - E_0\right), \quad (\text{E1})$$

where $\gamma' = \omega/m_{\gamma'}$, m is the mass of daughter particles, and N_i indicates the number of degree of i . The integration range of E_0 reads

$$E_0^{\max, \min} = E\gamma' \left(1 \pm \sqrt{1 - \gamma'^{-2}} \sqrt{1 - m^2/E^2} \right). \quad (\text{E2})$$

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