New constraint from supernova explosions on light particles beyond the Standard Model

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We propose a new constraint on light (sub-GeV) particles beyond the Standard Model that can be produced inside the protoneutron star core resulting from the core-collapse supernova explosion. It is derived by demanding that the energy carried by exotic particles being transferred to the progenitor stellar envelopes not exceed the explosion energy of $\lesssim 2 \times 10^{51}$ erg of observed supernovae. We show specifically that for the case of a dark photon which kinetically mixes with the SM photon and decays predominantly to an $e^{\pm}$ pair, a smaller mixing parameter of 1 order of magnitude below the well-established supernova cooling bound can be excluded. Furthermore, our bound fills the gap between the cooling bound and the region constrained by (non)observation of γ rays produced from supernovae for dark photons lighter than ~ 20 MeV. Our result also rules out the possibility of aiding successful supernova explosions by transferring energy from the supernova core to the shock with exotic particles.

I. INTRODUCTION

The Standard Model (SM) of particle physics has been the most successful theory that describes the fundamental properties and interactions between elementary particles. However, various hints from either the theoretical considerations or the cosmological and astrophysical observations point to the possibility that it is not a complete theory, and new particles beyond the SM (bSM) that only couple to the SM sector very weakly may exist.

Among the imperative searches and constraints of bSM particles, one important criterion comes from the observation of electron antineutrinos ($\bar{\nu}_e$) associated with the seminal core-collapse supernova (CCSN) event, SN1987A. The observed $\bar{\nu}_e$ burst duration of about 12 s, with individual energies up to 40 MeV, as well as the integrated total energy $\sim 5 \times 10^{52}$ erg [1–8], strongly supported the standard picture of neutrino cooling of the protoneutron star (PNS): The total gravitational binding energy, $E_G \sim 3 \times 10^{53}$ erg, released while forming a compact PNS with a mass $M_{\mathrm{PNS}} \sim 1.4 \, M_\odot$ and radius $R_{\mathrm{PNS}} \sim 10$ km is roughly equipartitioned by all six flavors of (anti)

neutrinos. Consequently, any bSM particles that can be produced inside the PNS and escape by taking away an energy comparable to E_G would have shortened the observed timescale of the $\bar{\nu}_e$ burst to be incompatible with the observation [9].

Constraints on various light bSM particles that may be produced in the hot and dense PNS core, based on the above argument, have been considered exhaustively in the literature, notably the axions [10–14], right-handed neutrinos [10,15,16], Majorons [17], Kaluza-Klein gravitons [18–20], Kaluza-Klein dilatons [18], unparticles [21,22], dark photons [23–26], dark matter [27–29], dilaton [30], saxion [31], Goldstone bosons [32,33], etc. Ideally, one should perform numerical simulations as in Refs. [19,34,35] to study the effects of a light bSM particle on the neutrino burst signal.

Other than affecting the PNS cooling, bSM particles produced inside the PNS may directly decay to photons, or indirectly produce the 511 keV lines via the pair annihilation by first decaying into $e^{\pm}$, outside the surface of the progenitor stars, $R_* \simeq 10^{14}$ cm. The (non)observation of γ rays associated with SN1987A, as well as the observed flux of 511 keV photons from the Milky Way, has been used to put constraints on bSM particles that couple electromagnetically to the SM sector [36–38]. Such derived bounds mostly complement those from the PNS cooling, because for bSM particles to decay outside R_*, the required coupling to the SM sector is usually not large enough to affect the PNS cooling.

In this paper, we propose a new constraint that bridges those from the PNS cooling and the γ-ray (non)observation. Our new constraint is based on a very basic fact: The known
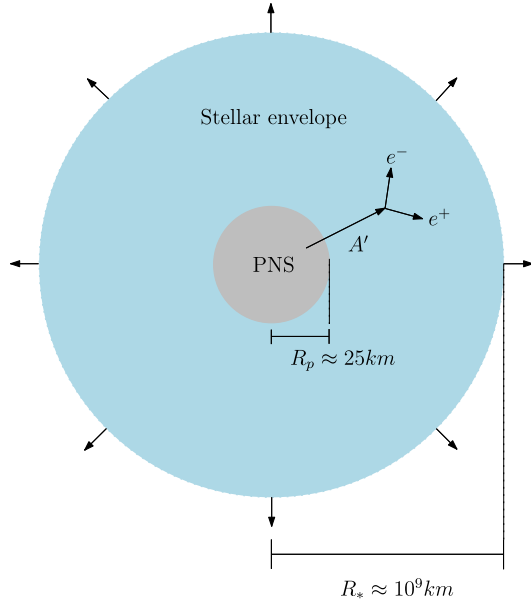
FIG. 1.   A schematic plot showing the energy deposition of the bSM particles produced from the PNS within a radius R_p into the stellar layers of the progenitor star with a radius R_*. Here we illustrate it with the example of a dark photon (A') decaying into an $e^\pm$ pair.

explosion energy of the CCSN of a progenitor star with $10~M_\odot \lesssim M_* \lesssim 20~M_\odot$ is $\simeq 1$ B, where B stands for bethe $\equiv 10^{51}$ ergs[39,40]. Most of this energy is carried by the kinetic energy of the expanding ejecta, with a mass of $\sim \mathcal{O}(10)~M_\odot$ and a velocity of $\sim 0.01c$ when we observed the emitted (quasi)thermal photons at $\geq \mathcal{O}(1)$ d after the core bounce [41]. In the absence of bSM physics, the prevalent theory is that the neutrinos emitted from the PNS within ~ 1 s after the core bounce can deposit a few percent of their energy to the stalled shockwave at $\sim \mathcal{O}(10^2)$ km to revive it [42]. The shock then wipes out the outer stellar envelopes at a speed of $\leq 0.1c$, giving rise to the observed explosion.

However, if bSM particles produced from the PNS can transfer the energy that they carry into the stellar envelopes or the shocked material before leaving the progenitor star, they would serve as a new energy source contributing to the total explosion energy (see Fig. 1 for a schematic plot). As a result, if this energy deposition mediated by bSM particles exceeds the observed explosion energy, after subtracting the gravitational binding of the stellar envelopes, such a bSM particle is then ruled out by CCSN observation.

Before working out a specific example, we first demonstrate analytically how these new bounds can improve the constraint derived from the PNS cooling. A well-known analytic criterion formulated by G. Raffelt of such states the following: For a novel cooling agent X that free-streams after production, its specific energy loss $\dot{\varepsilon}$ is bounded by [9]

$$\dot{\varepsilon}_X \lesssim \frac{L_\nu}{M_{\rm PNS}} \simeq 10^{19}~{\rm ergg}^{-1}~{\rm s}^{-1}, \qquad (1)$$

with $L_\nu \sim E_G/10 \simeq 3 \times 10^{52}~{\rm erg~s}^{-1}$ being the energy luminosity of all (anti)neutrinos and $\dot{\varepsilon}_X$ being evaluated at a typical core condition at ~ 1 s after the core bounce, with a temperature of $\simeq 30$ MeV and a density of $\simeq 3 \times 10^{14}~{\rm g~cm}^{-3}$.

The upper bound of the observed explosion energy of CCSNe associated with progenitor stars with zero-age main-sequence (ZAMS) masses between 10 and $20~M_\odot$ is mostly under $E_{\rm expl} = 2$ B (see, e.g., the compilations in Refs. [39,40,43]), while the typical binding energy of the stellar envelopes is $E_b \lesssim 1$ B (see later in this paper for details). Therefore, our proposed new constraint can be expressed by

$$K \cdot \dot{\varepsilon}_X \lesssim \frac{E_{\rm expl} + E_b}{\Delta t \cdot M_{\rm PNS}} \lesssim 10^{17}~{\rm ergg}^{-1}~{\rm s}^{-1}, \qquad (2)$$

where $\Delta t \simeq 10$ s, and $0 < K \leq 1$ denotes the efficiency of energy transfer into the region between a radius R_p, within which the particle X can be produced efficiently, and R_*. Comparing Eqs. (1) and (2), it is obvious that the new bound can exclude the bSM particle whose emissivity is ~ 2 orders of magnitude less than the one constrained by the PNS cooling, for cases where $K \sim 1$. For the rest of the paper, we consider a specific example of the dark photon that decays predominantly to an $e^\pm$ pair.

II. NEW CONSTRAINT ON DARK PHOTON

We consider the minimal extension of the SM with a $U(1)'$ dark sector. The dark photon (A') is the gauge boson of the broken $U(1)'$ symmetry which kinetically mixes with the hypercharge boson. When the dark photon mass is much smaller than the electroweak symmetry breaking scale, the mixing is effectively only with the photon (A). The effective Lagrangian for the photon–dark photon system is (see, e.g., Ref. [44] for the transformation from the dark photon gauge eigenstates to the mass eigenstates)

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{4} F'_{\mu\nu} F'^{\mu\nu} + \frac{1}{2} m^2_{A'} A'_\mu A'^\mu$$
$$- e \sum_f q_f (A_\mu + \epsilon A'_\mu) \bar{f} \gamma^\mu f. \qquad (3)$$

Here f is a SM fermion with electric charge q_f, and $m_{A'}$ and ϵ are the mass and the kinetic mixing parameter of the dark photon in the physical basis, respectively. Strategies for dark photon searches at colliders and fixed-target experiments, existing constraints on $(\epsilon, m_{A'})$, as well as anticipated sensitivities of planned experiments, can be found in the reports [45,46].

The in-medium physical eigenstates are quite distinct from those in vacuum due to the presence of the photon polarization tensor $\Pi = \Pi_R + i\Pi_I$ (see, e.g., Ref. [47]) in the inverse propagator matrix of the photon-dark photon system. As a consequence, in hot or dense stars, the collective effects of the stellar plasma can significantly change

the dark photon production rate [48,49]. References [25,26] found that plasma effects in the PNS qualitatively weaken the supernova cooling bound at dark photon masses below ~10 MeV.

In this work, we calculate the dark photon production rate, following closely Refs. [25,26]. For a dark photon weakly coupled to the thermal bath—i.e., when $\epsilon \ll 1$—one can invoke the in-medium effective kinetic mixing parameter

$$\epsilon_m^2 = \frac{\epsilon^2}{(1 - \Pi_R/m_{A'}^2)^2 + (\Pi_I/m_{A'}^2)^2} \quad (4)$$

for the transverse (T) and the longitudinal (L) polarizations separately. In CCSNe, the real part of the photon polarization tensor, $\Pi_{R|L,T}$, is dominantly generated by the electrons, which are relativistic and degenerate inside the neutrino sphere R_ν. The imaginary part $\Pi_{I|L,T}$ is determined mainly by the rates of the nuclear bremsstrahlung and the Compton scattering processes. Transversely and longitudinally polarized dark photons can thus be produced in the corresponding channels ($pn \to pnA'$, $pp \to ppA'$, and $\gamma e^- \to e^- A'$) through the effective in-medium mixing with the photon. Since for ϵ_m, the condition $\Pi_I \ll \Pi_R$ generally holds throughout the PNS environment, the resonant emission of longitudinal dark photons is open for $m_{A'} < \omega_p$, where ω_p is the photon plasma mass. Resonant emission of transverse dark photons, on the other hand, is only possible for $m_{A'}$ in a narrow range around ω_p.

Dark photons are reabsorbed in the supernovae mainly by the decay process $A' \to e^+ e^-$ when it is kinematically allowed. As pointed out in Ref. [26], in the PNS core region, dark photon decay is prevented due to the high electron chemical potential, unless $m_{A'}$ is larger than twice the effective electron mass in the plasma [50]. In this work, we are interested in the case in which the dark photon can escape the production region and decay freely in the stellar layers. The produced $e^\pm$ then quickly interact with the medium and lose their kinetic energy of $\sim$10–100 MeV to the surroundings in a length scale much shorter than R_* [51]. This effectively leads to an efficient transfer of the thermal energy from the PNS core region to the stellar envelope [$K \simeq 1$ in Eq. (2)].

For a given dark photon mass $m_{A'}$ and kinetic coupling ϵ, the total energy carried by the dark photons to a distance $R \geq R_p$ is calculated by [25]

$$L_{A'}(R, m_{A'}, \epsilon) = \sum_{L,T} \int_{r=0}^{R_p} \int_{\omega=m_{A'}}^{\infty} dr d\omega 4\pi r^2 e^{-\tau_{L,T}(r,\omega,R)}$$
$$\cdot \frac{\omega^3 v^3}{2\pi^2} e^{-\frac{\omega}{T(r)}} \epsilon_{m|L,T}^2(r, \omega)$$
$$\cdot [\Gamma_{iBr|L,T}(r, \omega) + \Gamma_{sC|L,T}(r, \omega)], \quad (5)$$

assuming that the nucleons and electrons are in local thermal equilibrium at temperature $T(r)$. Under this condition, the total production and the total absorption rate of SM photons of energy ω are related by $\Gamma_{\text{prod}} = e^{-\omega/T(r)} \Gamma_{\text{abs}}$, where Γ_{abs} is determined by Γ_{iBr} and Γ_{sC}, the inverse bremsstrahlung and the semi-Compton process rates, respectively. For Γ_{iBr}, we adopt the soft-radiation approximation and neglect many-body effects in the nuclear medium, as Ref. [24]. Therefore, dark photons are created through in-medium kinetic mixing with the SM photons at the rate $\Gamma'_{\text{prod}|L,T} = e^{-\omega/T(r)} \epsilon_{m|L,T}^2 \Gamma_{\text{abs}|L,T}^{(\text{in eq.})}$. The photon velocity in medium is $v = \sqrt{1 - m_{A'}^2/\omega^2}$. The term $e^{-\tau(r,\omega,R)}$ takes into account dark photon attenuation between radius r and R. We calculate the optical depth for a dark photon produced at radius r with energy ω, which travels radially outward to R by

$$\tau_{\text{radial out}}(r, \omega, R) = \left[\int_r^{R_p} d\tilde{r} \Gamma'_{\text{abs}|L,T} + (R - R_p) \cdot \Gamma'_{e^+ e^-} \right]$$
$$\qquad (6)$$

and include a correction factor to relate $\tau(r)$ to $\tau_{\text{radial out}}(r)$ as suggested by Ref. [25]. The dark photon absorption rate Γ'_{abs} receives contributions from the inverse bremsstrahlung processes, semi-Compton scattering, and decay to $e^\pm$ pairs. We have checked that outside R_p, Pauli blocking can be ignored, and one can use the decay rate in vacuum for $\Gamma'_{e^+ e^-}$.

The supernova cooling bound is determined by $L_{A'}(R_p) \leq L_\nu$ [cf. Eq. (1)] in the dark photon $(m_{A'}, \epsilon)$ parameter space. Our new bound, Eq. (2), is by requiring that the energy deposited by the decay of A' between R_p and R_* be smaller than the sum of the observed SN explosion energy and the total gravitational binding energy between these two radii:

$$E_d(R_p) \equiv [L_{A'}(R_p) - L_{A'}(R_*)] \cdot \Delta t \leq E_{\text{expl}} + \Delta E_g(R_p).$$
$$\qquad (7)$$

Here $\Delta E_g(R) \equiv E_g(R_*) - E_g(R)$, with

$$E_g(R) \equiv \int_0^R dr \frac{G\rho(r) M_{\text{enc}}(r)}{r} 4\pi r^2, \quad (8)$$

the gravitational binding energy inside radius R, where $M_{\text{enc}}(r)$ is the total mass enclosed in the region inside r. We fix the emission duration $\Delta t = 10$ s, which is the typical timescale of the PNS cooling.[1] Note that in Eq. (7), we have

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[1]Since the quantity $E_{\text{expl}} + \Delta E_g(R_p)$ on the rhs of Eq. (7) is only $\sim 10^{51}$ erg (see below), much smaller than the total binding energy of the PNS, $E_G \sim 3 \times 10^{53}$ erg, for dark photons that just carry and deposit an energy $E_d(R_p) \gtrsim E_{\text{expl}} + \Delta E_g(R_p)$, they would only alter the cooling behavior of the PNS core by ~1%. Therefore, the PNS cooling timescale of ~10 s should not be affected, and our derived bound based on $\Delta t = 10$ s is robust.
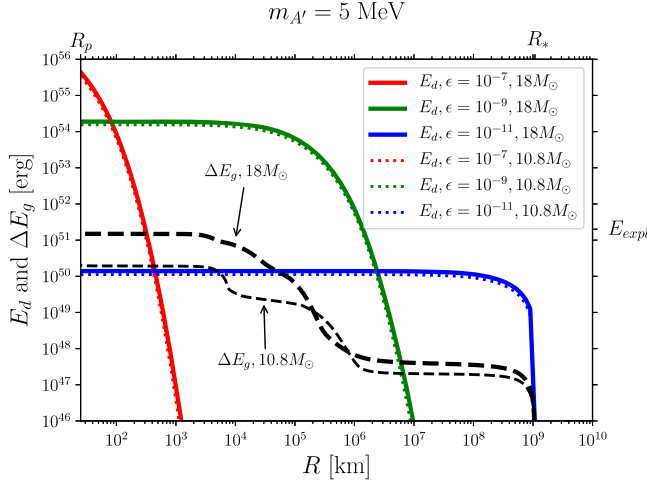
FIG. 2. Energy deposition $E_d(R)$ by dark photons to stellar envelopes outside radius R, for various dark photon parameters and for supernovae with progenitor masses of 18 $M_\odot$ (thick solid curves) and 10.8 $M_\odot$ (thin dotted curves). Also shown are the corresponding gravitational binding energy $\Delta E_g(R)$ outside R in both cases (thick and thin dashed curves).

FIG. 3. Shaded region: Excluded parameter space of dark photon derived using the observed SN explosion energy for progenitor masses of 18 and 10.8 $M_\odot$. The black dashed curve shows the bound determined by the PNS cooling argument for 18 $M_\odot$. Also shown is the excluded region inferred from the (non)observation of γ rays (dotted green curve), taken from Ref. [38].

neglected the kinetic energy of the shocked material, as well as that of the stellar envelope, which contribute at most $\sim 10\%$ of $E_{\rm expl}$.

The dark photon deposited energy $E_d(R_p)$ and the gravitational binding energy of the stellar envelope $\Delta E_g(R_p)$ depend on the structure of the PNS and the mass of the stellar progenitor. We examine two cases using the radial profile of the mass density, temperature, electron fraction, and electron chemical potential obtained by SN simulations of progenitor stars with 10.8 $M_\odot$ and 18 $M_\odot$ masses [52], chosen at $t = 1$ s after the core bounce. As those SN simulations do not contain the structure of the outermost hydrogen layer of the progenitor star, we extend the profile to R_* using the pre-SN structure provided by Ref. [53]. For both cases, we have used the same $R_p = 25$ km (slightly larger than R_ν) so as to encompass all the dark photon resonant production sites.

Figure 2 shows the comparison of $E_d(R)$ calculated with $m_{A'} = 5$ MeV and a few selected $\epsilon = 10^{-7}$, 10^{-9}, and 10^{-11}, to $\Delta E_g(R)$ for both progenitor masses. Different progenitor masses only lead to distinct $\Delta E_g(R)$ for $R > R_p$, but not $E_d(R)$, because the PNS structure is almost independent of the progenitor mass. For a given $m_{A'}$, dark photons with larger (smaller) ϵ carry more (less) energy away from the PNS and decay to $e^\pm$ at smaller (larger) radii above R_p. For $\epsilon = 10^{-7}$ and 10^{-9}, the energy deposition by the dark photon decay far exceeds the gravitational binding energy of the envelope by several orders of magnitude and can therefore be ruled out by our criterion. With $\epsilon = 10^{-11}$, dark photons only carry $\sim 10^{50}$ erg of energy away from the PNS and therefore cannot be ruled out by our constraint.
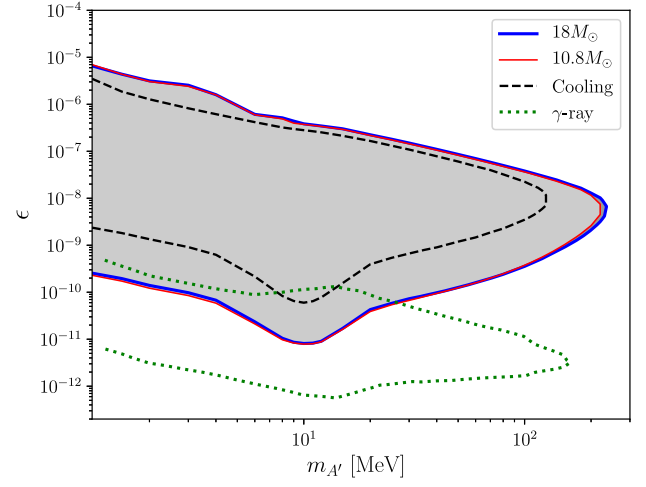
In Fig. 3, we show the contour plot for regions excluded by our new constraint, and that excluded by the PNS cooling, computed as aforementioned. In addition, we show the excluded region by the γ-ray (non)observation from Ref. [38]. The regions excluded by the observed SN explosion energy are nearly identical for both the 10.8 $M_\odot$ and 18 $M_\odot$ progenitors, because $E_d(R_p)$ are almost the same and $\Delta E_g(R_p) \ll E_d(R_p)$ for most of the excluded region other than those very close to the boundary (see Fig. 2). Their shapes closely follow and enclose that from the PNS cooling constraint. Their upper boundaries denote the ϵ value for which dark photons of mass $m_{A'}$ are produced copiously but also reabsorbed strongly inside radius R_p. Besides, as expected by our analytic estimate, this new consideration extends the excluded region to a lower ϵ by roughly 1 order of magnitude (which corresponds to a factor of ~ 100 in terms of dark photon emissivity) for a given $m_{A'}$. Note that as it largely overlaps with the γ-ray bound in the small-small-ϵ regime, they form together a robust bound covering nearly 6 orders of magnitude for $m_{A'} \lesssim 20$ MeV.

III. DISCUSSIONS

We have shown that the observed explosion energy of the CCSNe can be used to derive important constraints on light bSM particles that may be copiously produced from the PNS core. For dark photons that kinetically mix with SM photons, we show that our new bound excludes a larger parameter space than that derived using the observed neutrino burst from SN1987a. Moreover, it overlaps with the region recently obtained using the (non)observation

of γ rays produced by supernovae. Therefore, all three constraints together exclude a large range of parameter space that is not accessible by current terrestrial experiments or by cosmological observation. Although we have only considered the explicit example of dark photons, constraints on other bSM particles such that may effectively transfer energy from the PNS to the stellar layers—e.g., the sub-GeV axionlike particles [37] and MeV sterile neutrinos [54]—can be similarly derived.

Besides, in new physics models, it is sometimes speculated that the light bSM particles escaping the PNS may deposit energy into the gain region behind the stalled supernova shock to revive it and facilitate supernova explosions, in case neutrino heating is not effective. Such a scenario would typically require that light bSM particles provide an additional heating rate of $\sim$ a few times 10^{51} erg s^{-1}, similar to that from neutrino heating (see, e.g., Fig. 1 in Ref. [55]). As this rate is about 1 order of magnitude smaller than the luminosity L_ν used to derive the SN cooling bound, the corresponding parameter space cannot be excluded by the cooling bound. However, our new constraint dictates that the average luminosity of any bSM particles emitted from PNS cannot exceed $\sim$1% of L_ν, if they can deposit energy above PNS [see Eqs. (1) and (2)]. It thus rules out the possibility of a light bSM particle reviving the SN shock, because otherwise the continuous energy injection to stellar layers during the PNS cooling will lead to explosions that are too energetic.

Several uncertainties may affect the exact excluded region derived with the simple argument presented in this work for dark photons. For example, improved description of the dark photon emission from the nuclear bremsstrahlung beyond the soft radiation approximation adopted here, incorporating the somewhat uncertain condition of the PNS core temperature, density, and composition (see, e.g., Ref. [56]), as well as the time-dependence of the PNS

structure and the stellar envelope profile, may introduce some minor corrections. Nevertheless, we would like to emphasize that these uncertainties would affect all the derived bounds, including that from the PNS cooling and the γ-ray (non)observation. Therefore, the main message of this paper remains solid despite these uncertainties: the observed explosion energy of core-collapse supernovae places improved constraint on bSM particles that are able to transfer energy efficiently from the PNS core to the stellar mantle.

On the other hand, a detailed SN light-curve modeling taking into account ejecta driven by bSM particles can potentially provide even better constraints. For instance, even if the bSM particles only unbind the outermost part of the stellar envelope with an energy smaller than E_{expl} (see, e.g., the case with $\epsilon = 10^{-11}$ in Fig. 2), the standard neutrino-driven mechanism can still work to eject the entire inner layers. Depending on their relative velocity, those two ejecta may collide at times of a few days after the core collapse and lead to very luminous events not compatible with observations. The hydrogen layer of the stellar envelope may also be driven off by the energy deposition from bSM particles with a speed much larger than typical SN ejecta velocity, resulting in an electromagnetic precursor prior to the main supernova peak lights, or reducing the line feature of hydrogen. All these aspects require more dedicated work beyond the scope of this paper and deserve further exploration.

ACKNOWLEDGMENTS

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