Axionlike Particles from Hypernovae

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It was recently pointed out that very energetic subclasses of supernovae (SNe), like hypernovae and superluminous SNe, might host ultrastrong magnetic fields in their core. Such fields may catalyze the production of feebly interacting particles, changing the predicted emission rates. Here we consider the case of axionlike particles (ALPs) and show that the predicted large scale magnetic fields in the core contribute significantly to the ALP production, via a coherent conversion of thermal photons. Using recent state-of-the-art supernova (SN) simulations, including magnetohydrodynamics, we find that, if ALPs have masses $m_a \sim \mathcal{O}(10)$ MeV, their emissivity in such rare but exciting conditions via magnetic conversions would be over 2 orders of magnitude larger than previously estimated. Moreover, the radiative decay of these massive ALPs would lead to a peculiar delay in the arrival times of the daughter photons. Therefore, high-statistics gamma-ray satellites can potentially discover MeV ALPs in an unprobed region of the parameter space and shed light on the magnetohydrodynamical nature of the SN explosion.

Introduction.—Core collapse supernovae (SNe) are recognized as extremely efficient laboratories for light, feebly interacting particles, such as neutrinos, axionlike particles (ALPs), and dark photons [1,2]. Very often, supernova (SN) bounds significantly constrain the viable parameter space for these particles, providing guidance for their experimental searches. The enthusiasm for SNe as laboratories for astroparticle physics has grown in recent years, thanks to advancements in numerical simulations and the development of new neutrino detectors, which might provide important information about both the explosion mechanism of SNe and fundamental physics (see, e.g., Refs. [3–5] for recent reviews).

The current paradigm of the SN explosion mechanism is based on the neutrino-driven scenario, in which neutrino energy deposition revitalizes the stalled shock wave (see, e.g., Refs. [6,7]). However, there are some very energetic subclasses of supernovae, including hypernovae and superluminous SNe [8], which are highly unlikely to be explained by the conventional neutrino-driven explosion. Hypernovae and superluminous SNe are fascinating and rare events, with a typical Galactic rate $\sim 10^{-5}$ yr^{-1} [9], often advocated as an explanation of astrophysical conundrums such as *r*-process nucleosynthesis [10] or the source of Galactic positrons [11]. The most plausible scenario to account for these extreme events requires additional energy injection via the magnetohydrodynamically driven (MHD) explosions [12]. This situation has been explored in very recent dedicated numerical studies [13,14], finding that in this case the SN core might host ultrastrong magnetic fields ($B \gtrsim 10^{15}$ G [15]). The presence of a strong magnetic field is not the only possibility for such a subclass of SNe. For example, instead of considering a very energetic explosion, one may actually consider the possibility that the brightness of these events is due to the interaction of the ejecta with an extended and dense circumstellar envelope [16]. Nevertheless, the presence of strong magnetic field remains the most plausible scenario. Furthermore, such strong magnetic fields are also a plausible origin for more common types

of SNe, like type Ic-bl SNe [15], for which the analysis of the present Letter applies as well.

In light of these recent developments, it becomes essential to consider the question of whether the presence of potentially very large magnetic fields might influence significantly the emission of light, feebly interacting particles from SNe. It is known, for example, that if neutrinos possess a magnetic moment $\mu \gtrsim 10^{-13}\mu_B$, where μ_B is the Bohr magneton, strong magnetic fields may affect their flavor conversions in the SN core, leading to peculiar observational signatures (e.g., [17,18]).

However, the impact on other feebly interacting particles has never been studied before. Here we consider ALPs, which are among the most prominent and well-studied new physics candidates [5,19], and show that the presence of a magnetic field may have a significant impact on the emission of these particles if their mass falls in the $m_a \sim \mathcal{O}(10)$ MeV mass range. Our conclusions apply to a rare but spectacular class of events. As such, observations in the near future seem challenging and unlikely, but they would represent a completely new window on axion physics, SN explosion mechanisms, magnetic fields, early Universe cosmology, and much more.

Axionlike particles and supernovae.—ALPs are pseudo-scalar bosons a with a two-photon vertex, described by the Lagrangian [20]

$$\mathcal{L}_{a\gamma} = -\frac{1}{4} g_{a\gamma} F_{\mu\nu} \tilde{F}^{\mu\nu} a, \tag{1}$$

where $g_{a\gamma}$ is the effective ALP-photon vertex, $F_{\mu\nu}$ is the electromagnetic field, and $\tilde{F}^{\mu\nu}$ is its dual. These particles emerge in different extensions of the standard model (see Refs. [5,21] for recent reviews) and are the object of an intense experimental investigation [22,23]. Furthermore, astrophysical observations on different energy scales offer valuable opportunities to search for these particles [24–26].

Core-collapse SNe are particularly powerful cosmic factories for ALPs [2,24]. Notably, the SN 1987A neutrino detection has been a milestone event also for axion physics [27–29]. The dominant production channel in SNe for light ALPs coupling exclusively to photons is the Primakoff emission [30], in which thermal photons are converted into ALPs in the electrostatic field of protons. Once produced, SN ALPs provide different signatures depending on their mass. Very light ALPs ($m_a \lesssim 10^{-10}$ eV) leave the star and convert into gamma rays in the magnetic field of the Milky Way [31,32]. Indeed, the lack of a gamma-ray signal in the Gamma-Ray Spectrometer on the Solar Maximum Mission in coincidence with the observation of the neutrinos emitted from SN 1987A provides a strong bound on the ALP coupling to photons [31,32]. The most recent analysis finds $g_{a\gamma} < 5.3 \times 10^{-12}$ GeV^{-1}, for $m_a < 4 \times 10^{-10}$ eV [33]. On the other hand, ALPs with $m_a \sim \mathcal{O}(0.1$–$100)$ MeV decay into photons, thus producing a

gamma-ray flux without the need for a galactic magnetic field. In this case, a similar analysis gives $g_{a\gamma} \lesssim 10^{-11}$ GeV^{-1} at $m_a \simeq 10$ MeV [34]. Intriguing opportunities to sharpen these bounds are offered by the detection of an ALP burst in future (extra)galactic SN explosions [35,36] or from the diffuse ALP flux from all past core-collapse SNe in the Universe [37]. If detected, the ALP flux would carry information on the deepest SN regions, complementary to the ones provided by neutrinos and gravitational waves [38,39].

Magnetic fields could play a significant role in the ALP production itself. ALPs may be produced by the scattering of photons on an external macroscopic magnetic field, rather than on the electrostatic field of charged particles as in Primakoff emission. Recent investigations have shown, for example, that the magnetic ALP production in the Sun may dominate over the Primakoff production, particularly in some energy ranges [40–42]. An equivalent analysis for the SNe is lacking because of the absence, until very recently, of numerical SNe simulations that also included a magnetic field. In this Letter, we fill this gap and investigate the role that ultrastrong magnetic fields in a SN core have in the emissivity of ALPs.

MHD supernova model.—For this purpose, we consider a fully relativistic 3D magnetorotational core-collapse SN simulation with multienergy neutrino transport [43]. The progenitor model is a 20 $M_{\odot}$ star with solar metallicity from Ref. [44], while the nuclear part is described by the equation of state SFHo [45]. Since the original progenitor model assumed neither rotation nor magnetic field during its evolution, a cylindrical rotation and a dipolarlike magnetic field was added in Ref. [43]. The resulting star has a central initial angular frequency and magnetic field strength of 1 rad s^{-1} and 10^{12} G, respectively. These lead to a dynamically relevant magnetic field of $B \gtrsim 10^{15}$ G inside the protoneutron star after bounce through the magnetic compression and winding effects. This model experiences the so-called magnetorotational explosion soon after bounce and the shock wave reaches ~ 4000 km from the center at the final simulation post-bounce time $t_{pb} = 370$ ms.

In Fig. 1, we show the angle-averaged radial profile of the magnetic field strength B (dashed red) at $t_{pb} = 370$ ms in the SN core ($r \lesssim 50$ km) for the SN model we are considering. In the very inner core $r \lesssim 5$ km the magnetic field assumes a constant value $B \simeq 4 \times 10^{15}$ G. It reaches its maximum $B_{max} \simeq 7 \times 10^{15}$ G at $r \simeq 10$ km, where the strong shear motion appears, while it decreases at larger radii, becoming smaller than 10^{15} G at $r \gtrsim 30$ km.

We also show the angle-averaged radial profiles for the temperature T (blue), plasma frequency ω_{pl} (orange), and screening scale κ_s (black) in the inner SN core. The plasma frequency, defined as $\omega_{pl} \simeq 16.3$ MeV$(Y_e \rho_{14})^{1/3}$ [46], where $\rho_{14} = \rho/10^{14}$ g cm^{-3} and Y_e is the electron
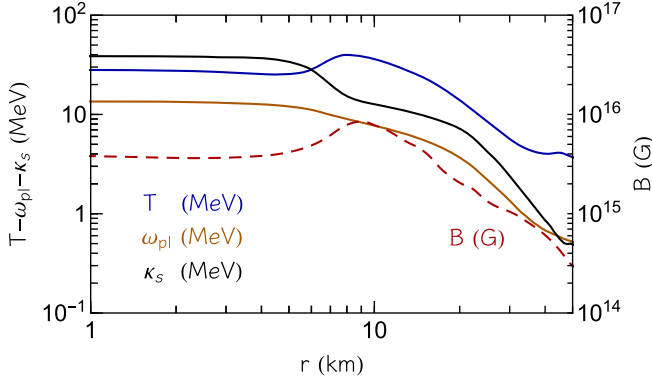
FIG. 1.   Angle-averaged radial profile for the temperature T (blue), plasma frequency ω_{pl} (orange), screening scale κ_s (black), and magnetic field strength B (dashed red) at $t_{pb} = 370$ ms.

fraction, plays the role of an "effective photon mass." The screening scale κ_s accounts for the correlation effects of the charged particles in the stellar medium. Since protons are only partially degenerate, a good estimate of the screening scale is given by $\kappa_s^2 = 4\pi\alpha n/T$, with n being the proton number density (see discussion in Ref. [33] for the effects of partial degeneracy).

ALP Primakoff emission.—The main ALP production channel in a SN core for ALPs with masses $m_a \lesssim 80$ MeV is the Primakoff process. Since electrons are highly degenerate in the SN core, their phase space is Pauli blocked and their contribution to the ALP production is negligible. Therefore, the most substantial contribution to the Primakoff process comes from free protons, which are only partially degenerate in the SN core. With these assumptions, the Primakoff production rate for massive ALPs is [29,47]

$$\Gamma_P = 2g_{a\gamma}^2 \frac{T\kappa_s^2}{32\pi} \frac{p}{E} \left\{ \frac{[(k+p)^2 + \kappa_s^2][(k-p)^2 + \kappa_s^2]}{4kp\kappa_s^2} \right.$$
$$\times \ln\left[\frac{(k+p)^2 + \kappa_s^2}{(k-p)^2 + \kappa_s^2}\right] - \frac{(k^2 - p^2)^2}{4kp\kappa_s^2}\ln\left[\frac{(k+p)^2}{(k-p)^2}\right] - 1 \left.\right\},$$

(2)

where the factor 2 accounts for the photon polarization states, E and p are, respectively, the energy and momentum of the ALP, and k is the momentum of the photon. In the conditions of interest for us, we can neglect the proton recoil energy and take the photon energy equal to the ALP energy, $\omega = E$.

ALP emission via magnetic conversions.—The rate for photon-ALP conversion in a magnetized plasma was recently calculated in Refs. [40–42]. ALPs can be resonantly produced in an external B field when the dispersion relations of longitudinal (L) or transverse (T) photon degrees of freedom match the dispersion relation of the ALP. The resonant emission rate is [40,42,48]

$$\Gamma_B = m_a^2 B_\parallel^2 g_{a\gamma}^2 \frac{\pi}{2\omega^2} Z_L \delta[\omega - \omega_0^L(k)]$$
$$+ \omega^2 B_\perp^2 g_{a\gamma}^2 \frac{\pi}{2\omega^2} Z_T \delta[\omega - \omega_0^T(k)],$$

(3)

where the factor $Z_{L,T}$ can be interpreted as renormalizing the coupling to the axion, $B_\parallel$ and $B_\perp$ are the B field components parallel and orthogonal to the photon momentum, and $\omega_0^{T,L}(k)$ is the dispersion relation for the considered electromagnetic mode [2] (see the Appendix for further details). Any excitation of the electromagnetic field in a medium, both longitudinal and transverse, is called "plasmon."

A word of caution is in order. The presence of very strong magnetic fields, $B \simeq 10^{16}$ G, could potentially affect the electron wave functions and, moreover, generate QED nonlinear effects. However, SNe feature a strongly degenerate and ultrarelativistic plasma, so the Fermi momentum is much larger than the synchrotron frequency $\omega_B = eB/p_F$. In such environments, the gas behaves quasiclassically, as the magnetic field does not quantize the electron wave functions [49–51]. This hierarchy also prevents from the production of electron-positron pairs. Moreover, the electron mass is modified to $m_e^2 \simeq (e^2/2\pi^2)p_F^2 \gg (m_e^{vac})^2$ [52,53], which implies that the critical magnetic field $B_c = m_e^2/e$ is much larger than in vacuum. Thus, we can neglect also birefringence effects [54,55], which typically inhibit the conversion in intense magnetic fields in vacuum [56]. Therefore, we identify a previously overlooked condition for efficient ALP-photon conversion in exceptionally strong magnetic fields that does not rely on the ALP being nonrelativistic [57].

Comparison of ALP fluxes.—The ALP emissivity for a given process is found by integrating the production rate over the photon thermal spectrum

$$Q_a = \int \frac{d^3\mathbf{k}}{(2\pi)^3} \frac{\Gamma\omega}{e^{\omega/T} - 1}$$

(4)

and the ALP luminosity by integrating the previous expression over the SN model,

$$L_a = \int d\Omega \int Q_a(r,\Omega)r^2 dr = \int d\omega\omega \frac{dN_a}{d\omega dt},$$

(5)

where $dN_a/d\omega dt$ is the differential ALP production rate. The latter is shown in Fig. 2 for three different values of the ALP mass, $m_a = 2$, 5, and 13 MeV. The plasmon-ALP conversion clearly dominates over Primakoff, especially for energies $\omega \lesssim 200$ MeV, if the ALP mass is in the range $4 \lesssim m_a \lesssim 14$ MeV. For $m_a \lesssim 4$ MeV, the resonance conversion takes place in the external regions of the star, where the magnetic field is too weak to lead to a sizable production. On the other hand, if the mass is too large, $m_a \gg 14$ MeV, the conversion happens in the very inner
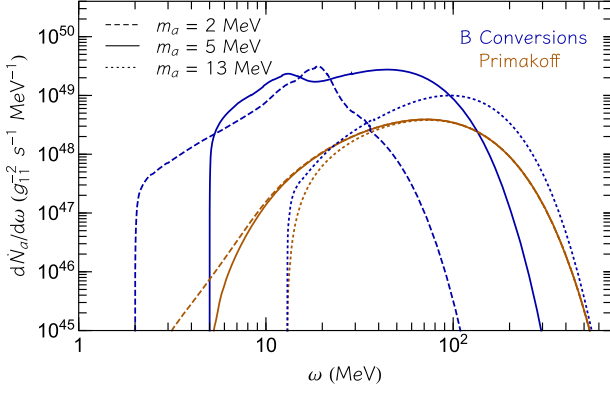
FIG. 2. Differential ALP production rate for the Primakoff (orange) and the B-conversion (blue) processes, for $m_a = 2, 5,$ 13 MeV (dashed, solid, and dotted lines, respectively) and $g_{a\gamma} = 10^{-11}$ GeV^{-1}, at $t_{pb} = 370$ ms. In the case of $m_a = 2,$ 5 MeV, at $\omega \gtrsim 10$ MeV, we notice a small peak due to the presence—in a very narrow energy range—of longitudinal modes conversion.

core ($r \lesssim 5$ km), where the plasmon population is suppressed. Notice that the contribution of the L mode is relevant only in a very narrow energy range, at low energy. At large energies, which—as we shall see below—are the most interesting for a possible detection, the transverse are the only relevant modes. Integrating the differential flux gives the total luminosity. We find that the B-conversion luminosity is larger than the Primakoff luminosity, $L_{a,B} > L_{a,P}$, for $4 \lesssim m_a \lesssim 14$ MeV, with a peak luminosity $L_{a,B}^{\max} \simeq 10^2 L_{a,P} \simeq 10^{49} (g_{a\gamma}/10^{-11}\,\text{GeV}^{-1})^2\,\text{ergs}^{-1}$ for $m_a \simeq 10$ MeV.

After being produced in the SN core, ALPs with $\mathcal{O}(1)$ MeV mass can decay into two photons with energy $E_\gamma = \omega/2$, producing a large gamma-ray flux that can be observed with existent [58] and planned [59–61] gamma-ray detectors. This flux will feature a characteristic time distribution, as the distance traveled by the ALP before decaying plus the distance traveled by daughter photons is larger than the distance of the SN from Earth.

Following Refs. [34,62], we estimate the time delay of an ALP-originated photon arriving at Earth, with respect to a massless particle coming directly from the SN, as

$$t = \frac{L_1}{\beta} + L_2 - d_{SN}, \qquad (6)$$

where β is the ALP velocity, $L_1 = 64\omega\pi/(g_{a\gamma}^2 m_a^4) \times \sqrt{1 - (m_a/\omega)^2}$ is the ALP decay length, and $L_2 = -L_1 \cos\alpha + \sqrt{d_{SN}^2 - L_1^2 \sin^2\alpha}$, with α as the angle between the photon and the ALP momentum, is the distance covered by the daughter photon before reaching Earth. In the relativistic limit, $\sin\alpha \simeq (\omega/m_a)^{-1}$ [34]. Finally, d_{SN} is the Earth-SN distance, which we fix to $d_{SN} = 10$ kpc and, in all cases of interest for us, is much larger than L_1.

The flux of daughter photons can be observed over the interval of several days with gamma-ray instruments such as the *Fermi* Large Area Telescope (LAT) [58]. In general, the number of events per unit time from the daughter photons expected in a gamma-ray detector with a delay time τ can be estimated as

$$\dot{N}_\gamma(\tau) = \frac{2}{4\pi d_{SN}^2} \left(\frac{dt(\omega_\tau^*)}{d\omega}\right)^{-1} \frac{dN_a(\omega_\tau^*)}{d\omega} A(\omega_\tau^*/2)$$
$$\times e^{-R^*/L_1}(1 - e^{-d_{SN}/L_1}), \qquad (7)$$

where $dN_a/d\omega$ is the SN ALP production rate, $dt/d\omega$ is obtained from Eq. (6), $R^* = 10^{14}$ cm [63] is the effective SN radius, and the factor e^{-R^*/L_1} ensures that only ALPs decaying outside the SN are counted. Analogously, the factor $(1 - e^{-d_{SN}/L_1})$ selects ALPs decaying before reaching Earth, and $A(\omega)$ is the detector effective area. Notice that Eq. (7) should be evaluated at $\omega = \omega_\tau^*$, where the ALP energy ω_τ^* is obtained by solving Eq. (6) with $t = \tau$.

Here, to give a realistic assessment of the potential to detect photons from ALP decay, we consider the specific case of *Fermi* LAT with Pass 8 event selection [64,65]. In Fig. 3, we show the number of expected events per unit time $\dot{N}_\gamma$—with 2σ Poissonian fluctuation contours—as a function of the time delay τ. The blue band refers to photons from the decay of Primakoff ALPs, and the orange band refers to photons from ALPs produced in the SN magnetic field. In both cases, we are fixing the ALP mass to $m_a = 5$ MeV and the ALP-photon coupling to $g_{a\gamma} = 10^{-11}$ GeV^{-1}.

Though the two signals have a qualitatively similar time evolution, the difference between them is manifest. Since ALPs produced through magnetic conversion are considerably less energetic than Primakoff ALPs, their daughter photons arrive with a more substantial delay. Thus, the photon flux in the first couple of days, being dominated by Primakoff ALPs, has a very weak dependence on the magnetic field, and would allow one to fix the ALP coupling to photons. The photon flux of the following days, after the Primakoff contribution decline, could then provide information about the intensity of the magnetic field. Therefore, reconstructing the time evolution of the signal allows us to gain confidence about the presence of a strong magnetic field in the SN and to measure its intensity.

As a final remark, we underline that the parameter region we are probing can also be constrained from cosmological observations based on big-bang nucleosynthesis. The cosmological argument, however, can be (at least partially) circumvented in low-reheating cosmological scenarios [66]. Therefore, a SN ALP detection in the mass range probed in this Letter would have important implications also for cosmology, pointing toward a nonstandard cosmological scenario. This shows a nice synergy between astrophysical and cosmological arguments to probe ALPs.
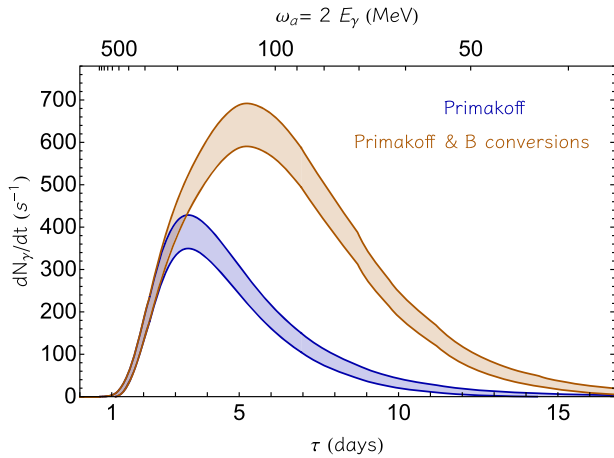
FIG. 3. Number of events per unit time with 2σ Poissonian fluctuation contours expected to be observed by *Fermi* LAT, as a function of the time delay, obtained considering only Primakoff (blue), and both Primakoff and B conversions (orange), for $m_a = 5$ MeV and $g_{a\gamma} = 10^{-11}$ GeV^{-1}. On the upper x axis, the ALP energy corresponding to each delay time τ is shown.

Conclusions.—Very energetic supernovae, such as hypernovae and superluminous SNe, host intense magnetic fields in their cores. We have shown that these fields would trigger a resonant ALP production via conversion of photons. Furthermore, we have demonstrated that if ALPs have a mass in the range $4 \lesssim m_a \lesssim 14$ MeV, the B conversion would overcome the Primakoff production by over an order of magnitude, becoming the dominant ALP production mechanism in SNe with $B \sim \mathcal{O}(10^{15})$ G.

These massive ALPs, with $m_a \sim \mathcal{O}(10)$ MeV and $g_{a\gamma} \sim 10^{-11}$ GeV^{-1}, have a decay length much shorter than the SN distance and so would produce an observable photon flux on Earth. Depending on the axion coupling, the flux might be observable by gamma-ray instruments such as *Fermi* LAT for well over a week. Furthermore, the temporal distribution of the flux is peculiar of the production mechanisms, providing a transparent way to recognize if ALPs are produced through conversions in a magnetic field or through the Primakoff process. Indeed, standard signals associated with violent events such as hypernovae and superluminous SNe are expected to peak at larger energies, above the GeV [67] and have a short duration. For this reason, hypernovae and superluminous SNe are often advocated as an explanation for gamma-ray bursts, which are extremely powerful signals that last from seconds to minutes [68]. Finally, since the magnetic production mechanism is very sensitive to the ALP mass, an accurate analysis of the time delay would allow us to infer m_a.

The discovery of a gamma-ray flux from a Galactic SN with the temporal properties described in this Letter would have enormous consequences both for ALP and SN physics. On the one hand, it would point out the existence of ALPs and allow us to pin down its mass; on the other, it

would be the smoking gun for the presence of ultrastrong magnetic fields in the SN core. Furthermore, the signal from axion decay will also be spread in angles away from the line of sight [34]; this may be particularly relevant in the future for closer SNe and gamma-ray instruments with better angular resolution.

We stress once more that our study mainly applies to a rare class of events. As such, observations in the near future seem unlikely but would offer an exciting possibility to learn about axions, as well as high-energy astrophysics and early Universe cosmology. This confirms once more SNe as multimessenger laboratories for particle physics and shows how much could be learned about SNe from the detection of ALP-induced processes.

Appendix: Plasmon dispersions and renormalization factors.—In this Appendix, we recollect the dispersion relations for the electromagnetic excitations in a QED plasma [2]. For an isotropic medium, in the Lorenz gauge, these dispersion relations for the longitudinal and transverse modes read

$$\sqrt{\omega^2 - v_*^2 k^2} = \omega_0^L(k) \equiv \sqrt{\omega_p^2[1 - G(v_*^2 k^2/\omega^2)]}, \quad \text{(A1a)}$$

$$\sqrt{\omega^2 - k^2} = \omega_0^T(k) \equiv \sqrt{\omega_p^2[1 + 1/2G(v_*^2 k^2/\omega^2)]}, \quad \text{(A1b)}$$

where G is an auxiliary function, defined as $G(x) = (3/x)\{1 - (2x/3) - [(1-x)/2\sqrt{x}]\ln[(1+\sqrt{x})/(1-\sqrt{x})]\}$, and v_* is the typical velocity for the electron and positron in

the plasma [2]. For the SN profile we considered, $v_* \gtrsim 0.99$ throughout the entire star.

The renormalization factors are also needed to evaluate Eq. (3),

$$Z_L = \frac{\omega^2}{\omega^2 - k^2} \frac{2(\omega^2 - v_*^2 k^2)}{3\omega_p^2 - (\omega^2 - v_*^2 k^2)}, \qquad \text{(A2a)}$$

$$Z_T = \frac{2(\omega^2 - v_*^2 k^2)}{\omega^2[3\omega_p^2 - 2(\omega^2 - k^2)] + (\omega^2 + k^2)(\omega^2 - v_*^2 k^2)}. \qquad \text{(A2b)}$$

One finds that Z_T is always very close to unity, while Z_L can strongly deviate from 1.

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