Transient High-energy Gamma-Rays and Neutrinos from Nearby Type II Supernovae

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Abstract

A dense wind environment (or circumstellar medium) may be ubiquitous in regular Type II supernovae (SNe II) before explosion, the interaction of which with the SN ejecta could result in a wind breakout event. The shock generated by the interaction of the SN ejecta and the wind can accelerate protons and subsequently high-energy gamma-rays and neutrinos could arise from inelastic pp collisions. In this work, we present detailed calculations of gamma-ray and neutrino production for regular SNe II. The calculations are executed by applying time-dependent evolution of dynamic and proton distributions so that the emission can be shown at different times. Our results show, for the SN 2013fs-like wind environment, multi-GeV and a few hundred TeV gamma-rays are detectable with a time window of several days at ≲2–3 Mpc by Fermi/LAT and the Cerenkov Telescopes Array during the ejecta–wind interaction, respectively, and can be detected at a further distance if the wind environment is denser.

We find the contribution of wind breakouts of regular SNe II to diffusing neutrino flux is subdominant by assuming all SNe II are SN 2013fs-like, whereas for a denser wind environment the contribution could be conspicuous above 300 TeV.

Key words: acceleration of particles – neutrinos – shock waves – supernovae: general

1. Introduction

Type II supernovae (SNe II) originate from the explosion of hydrogen-rich supergiant stars. For Type IIn and superluminous SNe, massive stars may experience mass-loss episodes before they explode as SNe, forming a dense wind environment (or circumstellar medium) (Smith et al. 2007; Miller et al. 2009; Ofek et al. 2013a, 2013b; Margutti et al. 2014). Recently, thanks to rapid follow-up spectroscopy observations, SN 2013fs, a regular SN II, is suggested to have a dense wind environment that is produced by the progenitor prior to explosion at a high mass-loss rate of \( \sim 3 \times 10^{-3} \left( v_w/100 \text{ km s}^{-1} \right) \dot{M}_e \text{ yr}^{-1} \), where \( v_w \) is the assumed velocity of the wind (Yaron et al. 2017). Very recently, further early observations of dozens of rising optical light curves of SNe II candidates indicate that SN 2013fs is not a special case and circumstellar wind material should be ubiquitous for regular SNe II (Forster et al. 2018). The pre-explosion mass loss can have a rate of \( \dot{M} > 10^{-3} \left( v_w/100 \text{ km s}^{-1} \right) \dot{M}_e \text{ yr}^{-1} \) and last for years. So universally, for regular SNe II, a high-density circumstellar wind environment in the immediate vicinity of the progenitor can be caused by sustained mass loss of the progenitor before explosion.

After explosion, the interaction of the SN ejecta with the optically thick wind can result in a bright, long-lived wind breakout event, which may also cause a delay of the usual envelope breakout. The shock generated by the interaction of the SN ejecta and the wind can accelerate protons, and subsequent inelastic pp collisions between the accelerated protons and the shocked wind gas can give rise to signatures of neutrinos and gamma-rays (Katz et al. 2012; Murase et al. 2011; Zirakashvili & Ptuskin 2016; Petropoulou et al. 2017). The produced neutrinos and gamma-rays could be a crucial probe to determine the nature of these explosive phenomena, e.g., the properties of the progenitor and the acceleration of cosmic rays (Murase et al. 2014). In addition, the produced neutrinos could contribute to diffuse neutrino emission (Li 2018).

In this work, we present gamma-ray and neutrino emission during the ejecta–wind interaction by focusing on regular SNe II, which have a much higher event rate than Type IIn and superluminous SNe, and adopt typical parameter values based on SN 2013fs. The dynamics of the ejecta–wind interaction are calculated assuming a time-dependent (radius-dependent) evolution, and emission of gamma-rays and neutrinos is derived through detailed calculations. Modifications to the proton distribution due to cooling and injection at different radii are taken into account. Since a dense wind environment is common for regular SNe II, we derive the diffuse neutrino flux contributed by wind breakout by assuming all SNe II are SN 2013fs-like. This paper is organized as follows. In Section 2, we describe the dynamics of SN ejecta–wind interactions. We calculate the different timescales of protons in the shocked wind region in Section 3 and present the gamma-ray and neutrino emission in Section 4 as well as the contribution to diffuse neutrino emission. Discussions and conclusions are given in Section 5.

2. Dynamics

An SN explosion ejects the progenitor’s stellar envelope. Typically the ejecta has a bulk kinetic energy of \( E_k = 10^{51} \text{ erg} \) and a total mass of \( M_{ej} = 10 \dot{M}_e \text{ M}_\odot \), inducing a bulk velocity \( v_h = \sqrt{2E_k/M_{ej}} = 3.2 \times 10^8 \beta^{1/2} \Omega^{-1/2} \text{ cm s}^{-1} \). After the shock breakout from the stellar envelope, the energy of ejecta with velocity larger than \( v \) can be described by (Matzner & McKee 1999; Li 2018)

\[
E(\geq v) = E_k (v/v_h)^{-\chi}
\]

where \( \chi = 3 + 5/n \). For convective (radiative) envelopes, one has \( n = 3/2 \) (3) and \( \chi = 19/3 \) (14/3) (Matzner & McKee 1999). In this work, we adopt \( \chi = 6 \) for red supergiant stars.
stars, while we also use $\chi = 5$ for blue supergiant stars, which gives a negligible difference.

After the shock breakout from the stellar envelope, the interaction of the SN ejecta with the wind forms a forward shock with velocity $v_s$ propagating through the wind and a reverse shock crossing the SN ejecta. Since the gamma-ray and neutrino emission of the reverse shock are usually weaker than those of the forward shock (Murase et al. 2011), we neglect the contribution of the former. For a regular SN II 2013fs-like case, according to measurements, the wind profile is suggested to be $\rho(R) = AR^{-2}$ with $A = M/4\pi v_w = 1.5 \times 10^{15} \text{g cm}^{-3}$ for $M = 3 \times 10^{-3} M_\odot \text{yr}^{-1}$ and $v_w = 100 \text{km s}^{-1}$, and the wind is confined within, but could extend up to, $R_w \sim 10^{15} \text{cm}$ (Yaron et al. 2017). We assume the wind starts from the radius of the stellar envelope $r_w$ with a typical value of around hundreds of solar radii. At a radius $R \gtrsim R_w$ in the wind, the energy of shock swept-up wind material is $E_w = v_s^2 \int_0^R 4\pi v^2 dr \simeq 4\pi AR^2 v_s^2$, where $v_s$ is the shock velocity. The energy of the shocked wind is given by the SN ejecta with velocity $v > v_s$, so one obtains the dynamical evolution of the shock speed in the wind by making $E_s(v) = E(v)\big|_{v \approx v_s}$ (Li 2018),

$$v_s = \left(\frac{E_s(v_s)}{4\pi v_s^2}\right)^{1/2} R^{-1/8} = 6.9 \times 10^8 R_s^{-1/8} \dot{M}^{-1/8} \rho^{1/2} \rho^{-3/8} \text{cm s}^{-1}. \quad (2)$$

Note that this equation is for $v_s > v_w$ while if $v_s < v_w$, the dynamical evolution should be derived by making $E_s = E_0$. In the situation considered here, Equation (2) is always valid.

The shock precursor has a characteristic optical depth, $\tau_c = c/v_s$, estimated by equating the radiation diffusive velocity and the shock velocity. The shock breakout occurs when the optical depth of the wind material ahead of the shock is $\tau_w = \tau_c$, where at radius $R$ one has $\tau_w \simeq \rho/\rho_0 \sigma T \sim R$. As a result, the shock breakout radius can be written as $R_{br} = 2.2 \times 10^{13} \rho_0^{1/3} \dot{M}^{1/9} \rho^{-1/3} \text{cm}$. If this radius is smaller than the stellar envelope, i.e., $R_{br} \lesssim r_w$, the shock breakout will take place on the stellar surface. Hereafter, we adopt $r_w = R_{br}$ for simplicity.

### 3. Particle Acceleration and Energy Loss

The SN shock may be radiation-mediated when the optical depth of Thomson scattering $\tau > \tau_c$, so that particle acceleration is prohibited (Katz et al. 2012; Murase et al. 2011). Particle acceleration can occur once the radiation starts to escape and the shock is expected to be collisionless, say, at $R > R_{\text{br}}$ (Waxman & Lobel 2001; Waxman & Katz 2016). The differential proton density accelerated and injected at radius $R (R > R_{\text{br}})$ is assumed to be a power law with a highest-energy exponential cutoff,

$$N_p^{\text{ inj}}(E_p, R) = N_0(R) E_p^{-\gamma} \exp(-E_p/E_{p, \text{max}}). \quad (3)$$

In order to find the highest proton cutoff energy $E_{p, \text{max}}$, one needs to evaluate the acceleration, dynamical, and cooling timescales of the protons. The magnetic field strength in the shocked wind can be estimated by $B = \sqrt{8\pi \epsilon_B \rho v_s^2} = 13.5 \epsilon_B R_{15}^{1/2} \dot{M}^{1/8} \rho^{3/8} \dot{M}^{1/2} \rho^{-3/8} \text{G}$, where $\epsilon_B$ is the equipartition parameter of the magnetic energy with a typical value $\epsilon_B = 0.01 c/\beta_s - 2$. The shock acceleration timescale is given by $t_{\text{acc}} = \kappa E_p / \beta_s^2 e B_0$, where $\beta_s = v_s/c$, and $\kappa$ indicates the uncertainty in the acceleration estimate. To determine the maximum proton energy, we adopt two values, i.e., $\kappa = 20/3$ and $\kappa = 1$ for Bohm diffusion and some theoretical prediction beyond the Bohm limit (e.g., Malkov & Diamond 2006). The dynamical timescale is $t_{\text{dyn}} \sim R/v_s$. The timescale of $pp$ cooling can be given by $t_{\text{pp}} = [0.5 \rho_{\text{pp}} R_{15} \text{cm}]^{-1}$, where $\rho_{\text{pp}}$ is the $pp$ collision cross section, $\rho_{\text{pp}} = \rho_{\text{sw}}/m_p^3$ is the number density of the shocked wind, and $\rho_{\text{sw}} = 4 \rho$. The timescale of proton synchrotron cooling in the magnetic field $B$ is $t_{\text{sync}} = 9(\gamma_p - 1) m_p^2 c^5 / 4 \epsilon_B B_{15}^2 \rho_{15}^{-2} \rho_{15}^{-3}$, where $\gamma_p = E_p / m_p c^2$ is the proton Lorentz factor and $\beta_p$ is the proton velocity in units of light speed.

Other proton cooling processes related to a low-energy radiation field, e.g., the photomeson ($\gamma p$) interaction and the Bethe–Heitler process (BH; $p + \gamma \rightarrow p + e + e^-$) are taken into account. For the regular SN II 2013fs-like case, the estimated bolometric luminosity based on multiband photometry is around a few $10^{52}$ erg s$^{-1}$ and the blackbody temperature is around a few $10^9 \text{K}$ (Yaron et al. 2017). In this work, a low-energy photon field is adopted as a blackbody distribution with temperature $kT = 2 R_{15}^{-1/2} \text{eV}$ so that the bolometric luminosity can be around the observational value. The photospheric radius can be evaluated from $\tau(\gamma_{0\text{B}}) = 1$, where $\tau(\gamma) = \int_{R_w}^{R} \rho \sigma T (\gamma - 1) dR$ is the optical depth of the material from $R$ to $R_w$, so one obtains $R_{\text{ph}} \approx 6 \times 10^{14} \text{cm} \sim R_w$, which indicates the assumption of a blackbody distribution of low-energy photons is approximately valid in this situation. All relevant timescales are plotted in Figures 1 and 2 for two representative radii, $R = 10^{14} \text{cm}$ and $R = 10^{15} \text{cm}$, respectively. $E_{p, \text{max}}$ can be obtained by letting $t_{\text{acc}} = \min(t_{\text{dyn}}, t_{\text{pp}}, t_{\text{sync}}, t_{\text{sync}})$. As we can see in two figures, $E_{p, \text{max}}$ is mainly determined by the $pp$ collision timescale and below $E_{p, \text{max}}$, the main energy-loss process is always $pp$ cooling for the adopted parameters. The photomeson and BH processes tend to be neglected as they basically operate at higher energies than $E_{p, \text{max}}$. Owing to $t_{\text{pp}} \propto \rho_{\text{sw}}^{-1} \propto R^2$ and $t_{\text{dyn}} \propto R^9/8$, at the smaller radius $pp$ cooling will be more dominant than dynamical evolution.

The time-dependent (or radius-dependent) energy injection rate of the shocked wind can be given by $L_{\text{sw}} = 4\pi R^2 u_{\text{sw}} v_s$, where $u_{\text{sw}} = \dot{M} / \rho_{\text{sw}} v_s^2$ is the energy density of the swept-up wind by the shock. The accelerated protons typically carry a fraction $\xi = 0.1 \xi_c$ of the shock energy, i.e., $L_p = \xi L_{\text{sw}}$, so one has $L_p = 4 \pi \xi A v_s^3$. As a result, the energy density of the accelerated protons can be described by $u_{p, \text{max}}^{\text{ inj}}(R) = L_p / 4 \pi A v_s = \xi v_s^2 \rho(R)$. The normalization factor $N_0(R)$ of the distribution of the injected protons can be derived using

$$u_{p, \text{max}}^{\text{ inj}}(R) = \int E_p N_p^{\text{ inj}}(E_p, R) dE_p. \quad (4)$$

### 4. Gamma-Ray and Neutrino Production

In our numerical calculations, the distribution of secondaries produced by $pp$ collisions is obtained by following the semi-analytical method provided by Kelner et al. (2006). A detailed treatment of secondaries from $pp$ collisions can be found in the Appendix. We denote the emissivity of gamma-rays or neutrinos as $N_{\gamma}(E_\gamma) = \tilde{\mathcal{F}}(N_{p}(E_p), n_{\text{pp}})$ by invoking an operator $\tilde{\mathcal{F}}$, where $\tilde{\mathcal{F}}, i_\gamma = \gamma$ or $\nu$. As suggested by Liu et al. (2018a), if we

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8 We neglect the contribution of possible He abundance.
consider a group of protons with a distribution of \( N_p^{\text{in}}(E_p, r) \) injected at a radius \( r \), when they propagate to a radius \( R \) the differential number density changes to

\[
N_p(E_p, r; R) = N_p^{\text{in}}(E_p, r) \exp[-(1 - 2^{1-\xi})\tau_{\text{pp}}(E_p, r, R) - (s - 1)\tau_{\text{ad}}(r, R)].
\]

Here, the main energy-loss processes as shown in Figure 1 and Figure 2, i.e., \( pp \) collisions and the adiabatic cooling, are taken into account during proton propagation. \( \tau_{\text{pp}}(E_p, r, R) = \sigma_{\text{pp}}(E_p)c \int_0^{(r/R)} n_{\text{se}}(\tilde{r})\, d\tilde{t} \) indicates the optical depth of \( pp \) collisions of protons injected at \( r \) propagating to \( R \). \( \tau_{\text{ad}}(r, R) = \int_0^R \frac{v_{\text{i}}(\tilde{r})\, d\tilde{r}}{\tilde{r}} = \int_0^R \frac{d\tilde{r}}{\tilde{r}} = \ln(R/r) \) is related to the adiabatic cooling of protons moving from \( r \) to \( R \). The differential luminosity of secondaries from \( pp \) collisions for the shock front at \( R \) can be derived by integrating over all radii \( (r_* < r < R) \), i.e.,

\[
L_i(E_i, R) = E_i^2 \int_{r_*}^R \mathcal{F}_i[N_p(E_p, r; R), n_{\text{se}}(R)]\, 4\pi r^2\, dr.
\]

High-energy gamma-rays produced by \( pp \) interactions will be attenuated by the low-energy photon field through \( \gamma + \gamma \rightarrow e^+e^- \) and absorbed by low-energy protons through the BH process in the emission region (Murase et al. 2011). Gamma-rays escaping from the emission region should be multiplied by a factor of \( [1 - \exp(-\tau_{\gamma\gamma} - \tau_{\text{BH}})]/\tau_{\gamma\gamma} + \tau_{\text{BH}} \), where \( \tau_{\gamma\gamma}(E_\gamma) \simeq R \int \sigma_{\gamma\gamma}(E_\gamma, \epsilon)N_e(\epsilon)\, d\epsilon \) and \( \tau_{\text{BH}} \simeq R\sigma_{\text{BH}}n_{\text{se}} \).

Two optical depths are calculated numerically in this work and, for simplicity, the cross section of the BH process, \( \sigma_{\text{BH}} \), is adopted approximately as a fixed value of 10 mb. The low-energy photon field, \( N_e \), is assumed to be a blackbody distribution as adopted above. Very high-energy photons will be attenuated due to the cosmic microwave background and extragalactic background light (EBL) by a factor \( e^{-\tau_{\text{CMB}} - \tau_{\text{EBL}}} \).

The model of EBL is based on Finke et al. (2010). Note that we neglect the contribution of secondary electrons produced by \( pp \) collisions even though the highest-energy electrons may radiate \( \sim \text{GeV} \) photons by synchrotron radiation in the adopted magnetic field. This is because the secondary electrons that can contribute \( \sim \text{GeV} \) photons are produced by protons with energies around the cutoff energy \( E_{p,\text{max}} \), where the proton luminosity is already significantly smaller than that of relatively low energy protons for an index \( s = 2 \) or softer due to an exponential cutoff. Another reason is that the emissivity of gamma-rays is about twice that of electrons produced during the \( pp \) interaction, so the electron synchrotron at the GeV band is subdominant.

The production of gamma-rays and neutrinos is presented in Figure 3. As we can see, gamma-ray emission above \( \sim 10\text{ GeV} \) will be suppressed significantly by the low-energy blackbody photon field, so a different setup of the low-energy photon field can give rise to a different gamma-ray flux. In this work, the low-energy photon field for a regular SN II is based on observations of SN 2013fs. Due to the absorption of the low-energy photon field in the shocked wind, the spectrum shows significant suppression in the energy range \( \sim 10\text{ GeV} - 100\text{ TeV} \), while the influence of the absorption of BH process is very weak, which can just be seen (the difference between the red solid line and the red dotted line below 10 GeV in Figure 3) at the early stage when the density of low-energy protons is high. In Figure 3, the sensitivities of Fermi/LAT and the CTA (Cerenkov Telescopes Array) are shown to compare with the gamma-ray emission. At \( R = 10^{14} \text{ cm} \), the duration of emission is \( t_d \approx R/v_\gamma \sim 10^{5.5} \text{ s} \), while at \( R = 10^{15} \text{ cm} \) one has \( t_d \sim 10^{6.5} \text{ s} \). For typical parameter values, i.e., \( \epsilon_B = 0.01, \xi = 0.1, \alpha_\ell = \delta = \kappa = 1 \), at 10 Mpc, high-energy gamma-rays are difficult to observe using current and next-generation telescopes for an SN 2013fs-like case. However, at a distance \( \lesssim 2 - 3\text{ Mpc} \), GeV gamma-rays can be detected by Fermi/LAT and gamma-rays around a few hundred TeV can be detected by the CTA. Note that through early-time spectrum modeling...
(e.g., Yaron et al. 2017) or early-time lightcurve modeling (e.g., Förster et al. 2018), one can only obtain the density, profile, and extended radius of the wind, while the mass-loss rate $M$ and the mass-loss duration $t_w$ before the SN explosion are estimated by assuming a wind velocity $v_w$ (Morozova et al. 2017). Yaron et al. (2017) achieved $M = 3 \times 10^3 M_\odot$ yr$^{-1}$ by assuming $v_w = 100$ km s$^{-1}$, while in Förster et al. (2018), $M$ has a comparable value but $v_w$ is much smaller (the terminal wind velocity is assumed to be $10$ km s$^{-1}$), indicating a much larger wind density (i.e., larger $\alpha$). To determine the gamma-ray radiation, we also tried a larger $\alpha$. For a denser wind environment (e.g., $\alpha = 3$ shown by the orange thin solid line in Figure 3), the gamma-ray flux is significantly enhanced and may be still detectable for a farther source.

The diffuse neutrino intensity from all SNe II wind breakouts in the universe can be given by integrating the contributions of individual wind breakout events at different cosmological epochs,

$$d\phi = c \int R(z) \frac{dN}{dE'}(1 + z) \frac{dt}{dz} dz,$$

(7)

where $dN/dE' = \int_{v_w}^{\infty} L_\nu(E', t)/E'^2 dt$ with $E' = E(1 + z)$.

$dN/dE'$, the latter term dedicated to expressing the total neutrino production for an individual wind breakout event; $L_\nu(E', t)$ can be found from Equation (6). Also, $dz/dt = H_0(1 + z)\left(\Omega_M(1 + z)^3 + \Omega_\Lambda\right)^{-1/2}$ and we adopt $\Omega_M = 0.27$, $\Omega_\Lambda = 0.73$ and $H_0 = 70$ km s$^{-1}$ Mpc$^{-1}$ in our calculations. $R(z) = R(0)S(z)$ is the SNe II event rate at redshift $z$, where $R(0)$ is the local event rate and $S(z)$ is the redshift evolution of an event rate that is assumed to follow the star formation rate (Yüksel et al. 2008). The volumetric rates of nearby core-collapse SNe are measured to be $0.7 \times 10^{-4}$ Mpc$^{-3}$ yr$^{-1}$ (Li et al. 2011), most of which are SNe II. Our result is presented in Figure 4. By assuming that all SNe events are SN 2013fs-like, the diffuse neutrino flux from wind breakouts of SNe II is subdominant in the diffuse neutrino flux detected by IceCube, with a contribution of around a few percent. However, if the wind environment is denser, e.g., $\alpha = 3$, and the maximum energy of accelerated protons is optimistic, i.e., $\kappa = 1$, the contribution of the wind breakouts of SNe II to the diffuse neutrino flux could be conspicuous above 300 TeV. The diffuse neutrino flux obtained in the numerical calculations is consistent with the analytical estimation of Li (2018). The detailed contribution to the diffuse neutrino flux is also down to the spectral index of protons, and a softer proton distribution will slightly reduce this contribution.

5. Discussion and Conclusion

5.1. High-energy Gamma-Rays and Neutrinos

In this work, we have studied the gamma-ray and neutrino emission during the interaction of SN ejecta with dense wind which may be expelled almost simultaneously with optical/infrared light. For an SN 2013fs-like wind, the ratio of shock velocity to bulk velocity is $v_\phi/v_b \approx 2H_0^{-1/2},(z^{-1/2}, \Omega^{-1/2})$, so one obtains the fraction of shock energy in the bulk ejecta energy,

$$\eta = E(>v_\phi) = (v_\phi/v_b)^{-6} \approx 9 \times 10^{-3}R_4^{3/4}(z^{-3/4})^{1/4}h^{-3/4}. $$

(8)

As a result, the wind breakouts of SN II shocks can convert a fraction $\xi \approx 9 \times 10^{-4} R_4^{3/4}(z^{-3/4})^{1/4}h^{-3/4}$ of the bulk energy into accelerated protons. These undergo significant cooling by $pp$ interactions and transfer almost their total energy to secondaries. For gamma-rays, under the typical parameters of the SN 2013fs-like case, $\sim$GeV and a few $\times$100 TeV gamma-rays can be detected at $\approx 2–3$ Mpc by Fermi/LAT and
the CTA during the ejecta–wind interaction, respectively. For the SN II wind breakout as the point source of neutrinos, at 10 Mpc, the flux is \( \sim 3 \times 10^{-10} \text{ GeV cm}^{-2} \text{ s}^{-1} \), which could be within the sensitivity level of the future IceCube Gen2 (Aartsen et al. 2017b); at closer distances, or for events within the Galaxy, the neutrinos can be detected by current IceCube (Murase 2018). Furthermore, the efficiency of the \( pp \) interaction is proportional to the number density of the wind, i.e., \( \propto \mathcal{A} \), so the fluxes of secondaries are proportional to \( \mathcal{A}^{7/4} \). Consequently, if an SN has a denser wind environment (\( \mathcal{A} > 1 \)), it can be still detectable even if farther away. By assuming all SNe II are SN 2013fs-like, we have determined the per-flavor diffuse neutrino flux from SN II wind breakouts as \( \sim 5 \times 10^{-10} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \), which is comparable with the estimated diffuse neutrino flux from SNe IIn even though the event rates of regular SNe II are much larger than those of SNe IIn (Petropoulou et al. 2017). One possible reason is that here we consider an SN ejecta with an intermediate velocity distribution, so the fraction of total ejecta energy converting to shock is quite small. However, for a denser wind, e.g., \( \mathcal{A} = 3 \), the diffuse neutrino flux from wind breakouts could reach a comparable level with the observed IceCube diffuse neutrinos above 300 TeV. Moreover, under the assumption that the low-energy photon field in the optical/ infrared energy band has a luminosity of a few \( \times 10^{42} \text{ erg s}^{-1} \), emitted gamma-rays with energies from tens of GeV to tens of TeV are significantly absorbed in the emission region. So in this case, the diffuse gamma-ray emission accompanying the diffuse neutrino emission can be estimated as \( \sim 1 \times 10^{-9} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \) without considering the cascade in intergalactic space. Such a diffuse gamma-ray flux is typically lower than that of the diffuse isotropic gamma-ray background (Ackermann et al. 2015).

For high-energy gamma-rays, \( \text{Fermi}/\text{LAT} \) and the CTA are able to detect signatures of the wind breakouts of SNe II at 2–3 Mpc for a time window of several days. This size is comparable with that of a local galaxy cluster. The expected SN II event rate in a local galaxy cluster is \( \sim a \) few over ten years (Mannucci et al. 2008). The search for accompanying gamma-rays for past nearby SNe II located in the field of view of \( \text{Fermi}/\text{LAT} \) could be a test of wind breakout, and a follow-up observation by \( \text{Fermi}/\text{LAT} \) and the CTA in the future will be valuable.

### 5.2. Lower-energy Radiation

In addition to high-energy gamma-rays, we give a brief discussion of radiation at other wavelengths. For \( \tau_e \lesssim \tau_{\nu} \equiv c/\nu_i \), the shock is expected to be collisionless, and its energy is \( \eta E_k \sim 10^{46} \text{ erg} \), only \( \xi \approx 0.1 \xi_{-1} \) of which is assumed to be converted to relativistic particles. Thus, most of the energy of the shock is thermal energy. The temperature of thermal protons in the immediate downstream of shock can be estimated as \( kT_p = 3m_p\nu_i^2/16 \approx 93 \text{ keV} \) for typical values of relevant parameters in this work. The electron temperature should not be larger than the equipartition temperature (\( \approx 47 \text{ keV} \)) but it is still uncertain due to the unknown efficiency by which protons transfer energy to electrons in collisionless shocks. However, since collisionless shock heating is typically faster than Coulomb collisional processes (Katz et al. 2012), a lower limit for the electron temperature can be obtained by assuming the shock is collisional (in other words, there is no collisionless heating). In the absence of collisionless shock heating, the electron temperature is obtained from the balance between Coulomb heating and cooling processes. If the fastest cooling process is inverse Compton scattering off the radiation field, one has \( kT_e \sim 40U^{-2/3}n_{e,10}^{2/3}(kT_p/100 \text{ keV})^{2/3} \text{ keV} \) (Waxman & Loeb 2001; Katz et al. 2012; Murase et al. 2011), where \( U = L/(4\pi R_i^2 v_d) \approx 10^{3} L_{43} R_i^{-2} v_d^{-1} \text{ erg cm}^{-3} \) is the energy density of
the low-energy radiation field and $v_d = c/\tau_w$ is the diffusion velocity of light. Consequently, X-rays can naturally be expected for electrons with energies of tens of keV via inverse Compton or thermal bremsstrahlung (Chevalier & Irwin 2012; Pan et al. 2013). Since the electron energy is from Coulomb heating of protons, the radiation efficiency of shocked material can be obtained by comparing the proton cooling timescale $\tau_p \sim 2 \times 10^4 U_{\gamma,10}^{-3/5} n_{sw,10}^{-8/5} (kT_p/100 \text{ keV})^{-3/5} \text{ s}$ with the dynamical timescale $R/v_h \sim 10^6 \text{ s}$ (Katz et al. 2012). As a result, the cooling of the shocked material is efficient and contributes thermal X-rays, whose luminosity is about $(1 - \xi)/\xi$ of that of non-thermal gamma-rays if we neglect their external absorption, i.e., $\sim 10^{43} \text{ erg s}^{-1}$.

Relativistic electrons including secondary electrons (from $pp$ collisions) and primary electrons (co-accelerated with protons by the shock) can contribute to non-thermal X-ray, radio, and MeV gamma-ray emission, and the electromagnetic cascade initiated by the absorbed high-energy gamma-rays in the emission region can also contribute (Murase et al. 2018). Accurate calculations of these are challenging since inelastic Compton scattering by thermal electrons, along with other complexities proposed in Waxman & Katz (2016), plays a crucial role in determining the distributions of electrons and photons, but they are not generally in thermal equilibrium. However, in the X-ray energy band, non-thermal contributions tend to be subdominant because the radiation of thermal electrons is efficient and their energy is typically larger than that of relativistic electrons, most of which is still thermal energy, as we mentioned above. Radio emission could arise from secondary and primary electrons, but it may be suppressed by free–free absorption, synchrotron self-absorption, and the Razin–Tsytovich process, and modified by Comptonization of thermal electrons (Murase et al. 2014). Soft X-rays are expected to be up-Comptonized and gamma-rays to be degraded by thermal electrons to some extent, depending on the opacity of Compton scattering. The typical photon energy may be comparable to that of thermal electrons, i.e., a few tens of keV. Soft X-ray and radio emission has been reported in some SNe II (e.g., Pooley et al. 2002; Chevalier et al. 2006, and references therein), but their luminosities are usually weak, ranging from $10^{37}$ to almost $10^{42} \text{ erg s}^{-1}$ (Dwarkadas 2014), implying suppression due to Comptonization. Although we mainly focus on high-energy gamma-ray emission, and detailed discussions of X-ray and radio emission are beyond the scope of this work, in the future, in addition to gamma-rays, observational constraints on X-ray and radio emission could help to check the ejecta–wind interaction model for regular SNe II and to include the wind environment. In particular, X-ray missions such as the HXMT (Xie et al. 2015) and the Einstein probe (Yuan et al. 2015) should significantly improve the prospects for the detection of accompanying X-ray emission which, in addition to high-energy radiation, will help us to understand the progenitor nature of SNe II.

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Appendix

Secondaries Produced by $pp$ Collisions

Basically, we follow the semi-analytical method provided by Kelner et al. (2006) (see also Kafexhiu et al. 2014; Liu et al. 2018b). The differential production in unit energy and unit time is given by

$$\mathcal{F}_i(E_i) = c n_{sw} \int_{E_i}^{\infty} \sigma_{pp}(E_p) N_p(E_p) F_i \left( \frac{E_i}{E_p}, E_p \right) \frac{dE_p}{E_p},$$

where $i$ can be $\gamma$ or $\nu$, and the cross section $\sigma_{pp}(E_p) = 34.3 + 1.88 L + 0.25 L^2 \text{ mb}$ with $L = \ln(E_p/1 \text{ TeV})$. $F_i$ is the
spectrum of secondary $\gamma$ or $\nu$ in one collision, which can be found in Equations (58), (62), and (66) of Kelner et al. (2006). The above analytical presentation works for $E_p > 100$ GeV, while for $E_p < 100$ GeV the spectra of secondaries can be continued to low energies using the $\delta-$ functional approximation for the energy of produced pions (Aharonian & Atoyan 2000):

$$F(E) = 2cn_{\text{cr}} \frac{n}{K_\pi} \int_{E_{\text{min}}}^{\infty} \sigma_{pp} \left( m_p + \frac{E_\pi}{K_\pi} \right) \sigma_{\nu} \left( m_\nu + \frac{E_\nu}{K_\nu} \right) \frac{dE_\pi}{\sqrt{E_\pi^2 - m_\pi^2}},$$

(10)

where $E_\pi$ is the pion energy and the pion rest mass $m_\pi \approx 135$ MeV for gamma-ray production and $m_\pi \approx 140$ MeV for neutrino production. $E_{\text{min}} = E_i / \zeta_i + \zeta_\pi m_\pi^2 / 4E_i$ with $\zeta_\gamma = 1$ and $\zeta_\nu = 1 - m_\nu^2 / m_\pi^2 = 0.427$, $K_\pi = 0.17$, and $n$ is a free parameter that is determined by the continuity of the flux of secondaries at 100 GeV. At lower energies one can use a more accurate approximation for the inelastic $pp$ interaction cross section instead, i.e., $\sigma_{pp}(E_p) = (34.3 + 1.88L + 0.25L^2)[1 - (E_\pi/E_p)^4]^2$ mb with $E_{\text{th}} = 1.22$ GeV.

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