Exploring the Gamma-Ray Emissivity of Young Supernova Remnants I: Hadronic Emission

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ABSTRACT

Using a simplified model for the hadronic emission from young supernova remnants (SNRs), we derive an expression to calculate the hadronic luminosity with time, depending on the SN ejecta density profile and the density structure of the surrounding medium. Our analysis shows that the hadronic emission will decrease with time for core-collapse SNe expanding in the winds of their progenitor stars, but increase with time for SNe expanding into a constant density medium, typical of Type Ia SNe. Using our expressions, we can compute the time-dependent hadronic flux from some well-known young SNe and SNRs with time, and where applicable reproduce previous results in the appropriate parameter regime. Using our calculations, we also emphasize the exciting possibility that SN 1987A may become a visible gamma-ray source in the next decade.

Key words: Acceleration of particles; Shock waves; supernovae: individual: SN 1987A; cosmic rays; supernova remnants; gamma-rays: ISM

1 INTRODUCTION

Supernova remnants (SNRs) are considered as a source of very high energy accelerated particles, at least up to the knee of the cosmic-ray spectrum. Details of the process by which the acceleration happens are not completely understood, but it is assumed to be related to Diffusive Shock Acceleration (DSA; Drury 1983; Malkov & O'C Drury 2001) and its nonlinear modification (Ellison et al. 1997). The discovery of X-ray synchrotron emission from SN 1006 provided the first convincing evidence that electrons can be accelerated to TeV energies. Recent discoveries of SNRs at γ-ray wavelengths have further supported the notion that SNRs can accelerate particles, both electrons and protons, to GeV and even TeV energies. Most recently, direct detection of the pion decay signature in 2 SNRs conclusively shows the presence of accelerated protons in SNRs (Ackermann et al. 2013).

The success of the Fermi telescope, coupled with ground-based Cerenkov telescope arrays, has considerably increased the number of supernova remnants (SNRs) observed in γ-rays. At present the number is large enough that one can begin to determine statistical properties of SNRs, and correlations between their multiwavelength properties. The upcoming release of the Fermi-LAT Supernova Catalog will allow for a comprehensive study of SNRs in gamma-rays, and comparison with multi-wavelength data. The purpose of the catalog is to systematically investigate the properties of observed gamma-ray SNRs in a multi-wavelength context (Brandt et al. 2012). It will characterize GeV emission from SNRs, examine multi-wavelength correlations, determine statistically significant SNR correlations, and calculate a spectral model. The existence of a large number, the continual detection of many more SNRs, and the very nice cataloging of their high energy and multi-wavelength properties, requires a commensurate investment of directed theoretical effort to fully understand the multi-wavelength and high energy properties of SNRs.

In the past decade, there has been considerable interest in understanding the very high energy emission from, and broadband spectrum of, SNRs. There exist very good papers by many groups that carry out sophisticated models for understanding the broadband multiwavelength emission from SNRs, and in particular determining whether the very-high energy (GeV–TeV) emission is due to hadronic or leptonic processes (Berezhko & Völk 2010; Edmon et al. 2011; Caprioli 2011; Ellison et al. 2012; Castro et al. 2012; Atoyan & Dermer 2012; Morlino & Caprioli 2012; Berezhko et al. 2013). Such modelling often requires very complicated codes containing a variety of different physics, which have been developed mainly in the last decade. However, even with sophisticated modelling, it has been hard to understand the origin of the very high energy emission in most SNRs, and to delineate the leptonic and hadronic contributions. Unfortunately, complicated models that require several weeks to months of work per SNR are not conducive to understanding a statistical ensemble of SNRs. What is needed is a simpler, but still...
reasonable accurate way to understand the statistical properties of SNRs, that can reasonably reproduce the spread in overall properties even if it cannot accurately fit the emission from individual SNRs. Such models must necessarily be (semi)-analytic (although informed by the results of numerical calculations), where the effect of varying one or more parameters, keeping all else constant, and studying the results on a statistical sample of SNe, can be achieved in a short time.

In anticipation of the need to understand the general properties of an ensemble of SNRs, and the release of the Fermi SNR catalog, we have embarked on a project to explore the properties of gamma-ray SNRs, especially young SNRs that have not reached the Sedov-Taylor or adiabatic stage. Using semi-analytic arguments coupled with realistic approximations of SNR evolution, we plan to study the gamma-ray emissivity of SNRs, and investigate the time evolution of the gamma-ray luminosity due to pion decay and leptonic processes (non-thermal bremsstrahlung and Inverse Compton emission). In this first paper we concentrate on the solutions for the luminosity in different environments. In §2, we compute an analytic expression for the luminosity of young SNRs due to pion decay, that illustrates the various factors that affect the luminosity. We also explore expressions for the luminosity in different environments. In §4 we apply the model to individual SNe and SNRs such as Cas A, SN 1993J and SN 1987A. §5 puts our work in the context of previously published literature on the subject, and discusses its successes, shortcomings and implications. Finally, §6 summarizes our work, and outlines future directions.

2 EVOLUTION OF YOUNG SUPERNOVA REMNANTS

In the past, the young SNR phase, which we alternately refer to here as the ejecta-dominated phase, has often been neglected in considerations of the particle-acceleration process. This is slowly changing in recent times (Ptuskin et al. 2010; Caprioli 2011; Ellison & Bykov 2011; Ellison et al. 2012; Telezhinsky et al. 2012; Zirakashvili & Ptuskin 2012; Telezhinsky et al. 2013). The reasons generally given for neglecting it are the phase only lasts for a short time, until the swept-up mass equals the ejecta mass; the shock velocity is constant (e.g. Ptuskin & Zirakashvili 2003; Gabici & Aharonian 2007; Prantzos 2011); and that the maximum energy of the particles increases up until the end of the ejecta-dominated stage and the beginning of the Sedov-Taylor stage (Helder et al. 2012), when it is that particles start escaping. However, all of these assertions are false (Dwarkadas 2011). As has been shown by many authors, since at least Gull (1972), the time taken to reach the Sedov-Taylor stage is much larger than assumed in many papers. It requires the swept-up mass to significantly exceed the ejecta mass, by a factor of 20-30, depending on the ejecta profile (Dwarkadas & Chevalier 1998). This considerably increases the time taken to reach the Sedov stage, with the result that a 1000 year old remnant such as SN 1006 is still very far from the Sedov stage. Similarly, SNRs expanding in the low-density interiors of wind-blown bubbles from massive stars may take a long time to sweep-up enough mass to reach the Sedov stage, and then in some cases may bypass the Sedov stage altogether (Dwarkadas 2007). The second assertion, that the velocity is approximately constant, is also incorrect. The velocity is continually decreasing in the ejecta dominated stage. Although this is shown by several calculations (Chevalier 1982), the definitive proof is in the observations of young SNRs such as Cas A, Tycho and SN 1006, where the observed shock velocities are much lower than the fiducial value of $> 10^5$ cm s$^{-1}$ that has been conveniently assumed as the “free-expansion” velocity in many papers. The last assertion, that the maximum energy of accelerated particles is only reached at the beginning of the Sedov stage, has been disproved by several arguments (Tatischeff 2009; Dwarkadas et al. 2013), as well as excellent numerical simulations (Bell et al. 2013). They show that the maximum energy is reached early in the young SNR stage in most cases, after which it begins to decrease. To summarize, the young ejecta-dominated SNR stage lasts for a long time, potentially up to several thousand years for SNRs exploding in a low density region or wind-blown cavity. The maximum energies to which particles are accelerated is reached generally early in this stage, the velocity continually decreases, and the escape of particles from the SNR happens in this phase. Given these considerations, we feel that it is very important to study SNRs in this stage.

In order to describe the evolution of a young SNR, we use the formulation suggested by Chevalier (1982; Chevalier & Fransson 1994). In brief, the expansion of SN ejecta into the surrounding medium leads to the formation of a double-shocked structure, consisting of a reverse shock that travels back into the ejecta, and a forward shock that expands into the ambient medium. From analytical models as well as hydrodynamic simulations of exploding core-collapse stars, it has been found that SN ejecta can be effectively described by $\rho_{ej} = A e^{-n t^{-3}}$, where the coefficient $A$ depends on the explosion energy and the mass of the ejecta, and can be evaluated with the information given in Chevalier & Fransson (1994). The surrounding medium can generally be ascribed as having a power-law profile in density, $\rho_{amb} = B r^{-s}$, where $s = 0$ denotes a constant density medium, and $s = 2$ denotes a wind medium with constant mass-loss rate $\dot{M}$ and constant wind velocity $v_w$, with the coefficient $B = \dot{M} / (4 \pi v_w^2)$. The resulting expansion of the contact discontinuity happens in a self-similar fashion, with the self-similar solutions given by (Chevalier 1982):

$$R_{CD} = \left(\frac{\delta A}{B}\right)^{1/(n-s)} t^{(n-3)/(n-s)}$$

where the value of the parameter $\delta$ is given in Chevalier (1982). We can write the expansion of the contact discontinuity as $R_{CD} = C t^m$, where $m = (n-3)/(n-s)$ is referred to as the expansion parameter. Note that since the solutions require $n > 5$, and $s < 3$, we have $m < 1$. Since the expansion is self-similar, the forward and reverse
shocks will expand in the same manner. In the self-similar case the ratio of the shocks to the contact discontinuity, and to each other, will be fixed. We can write the radius of the forward shock as \( R_{sh} = \kappa R_{CD} = \kappa C t^m \). The velocity \( v_{sh} = dR_{sh}/dt = m\kappa C t^{m-1} \), and is therefore always decreasing with time.

### 3 HADRONIC EMISSION FROM YOUNG SUPERNova REMNANTS

The interaction of protons accelerated at the SNR shock front with protons in the interstellar medium gives rise to neutral pions (among other species) which subsequently decay to give gamma-rays. The gamma-ray flux of SNRs due to hadronic emission can be written as \( F_{\gamma}(>E_\gamma,t) = \frac{q_\gamma}{4\pi d^2} \frac{M(t)}{\mu m_p} \left[ \frac{\epsilon_{CR}}{V} \right] \) \( (2) \)

where \( q_\gamma \) is the \( \gamma \)-ray emissivity normalized to the cosmic-ray energy density, and tabulated in Drury et al. 1994. Note that all the information regarding the nuclear interactions and spectrum of accelerated particles is contained in the parameter \( q_\gamma, \epsilon_{CR} \) is the energy in cosmic rays, \( V \) is the emitting volume and \( d \) is the distance to the source. \( M(t) \) is the mass of material with which the accelerated protons are interacting, \( \mu \) is the mean molecular weight, and \( m_p \) the proton mass.

We consider a SNR expanding in a medium with surrounding density \( B r^{-s} \), where \( B \) is assumed to be a constant. As the SN shock expands outward at high velocity, it sweeps up the material ahead of it. The mass swept-up by the shock, which forms the “target” mass \( M(t) \) with which accelerated protons will interact to give hadronic emission, is given by:

\[
M_{sw} = \int_0^R 4\pi r^2 B r^{-s} dr = \frac{4\pi B}{3-s} R^{3-s} \tag{3}
\]

The energy density of cosmic-rays from SNRs is a highly-debated question. Here we take the approach that a fraction of the energy available at the shock front is used to accelerate cosmic-rays. The maximum energy that can be extracted from the shock front at any given time in the ejecta-dominated phase is less than the total kinetic energy of the explosion, and is \( 2\pi R_{sh}^2 \rho_{sh} V_{sh}^3 \). The total energy expended in cosmic ray acceleration up to a given time \( t \) will be some fraction of the integral of this quantity over time \( t \):

\[
\epsilon_{CR} = 2\pi \int_0^t \epsilon_{CR} R_{sh}^2 \rho_{sh} V_{sh}^3 dt \tag{4}
\]

In principle the hadronic flux needs to be written as \( \frac{q_\gamma}{4\pi d^2} \int n \left[ \frac{\epsilon_{CR}}{V} \right] 4\pi r^2 dr \)

where \( n \) is the density of the surrounding medium. This gives the same time-dependence as found here, with some variation on order unity in the normalization factor. We prefer to write it this way so that we can easily differentiate the contribution of various shocks, and delineate various contributing factors, as shown in this paper.

\[
= \frac{2\pi B \xi (\kappa C_1)^{5-2s} m_3^3 \int_0^t t^{2m-3s+3m-3} dt}{5m-3s-2} \tag{5}
\]

\[
= \frac{2\pi B \xi (\kappa C_1)^{5-2s} m_3^3 \int_0^t t^{2m-3s-2} dt}{5m-3s-2} \tag{6}
\]

where \( \xi \) denotes the fraction of the shock energy that is converted to cosmic rays. For convenience we assume here that \( \xi \) is a constant, but there is no reason why it should be so, and it is quite possible that it could be a function of time. We note that for a constant density medium \( (s = 0) \), the available energy at the shock front increases with time as long as \( m > 2/3 \), i.e. as long as the SNR has not entered the Sedov-Taylor stage, when all the kinetic energy is available. Similarly, for a wind medium, the available energy goes as \( t^{1/2} \) and thus will continually increase until \( m = 2/3 \), which is when the remnant enters the Sedov-Taylor phase in a wind \( (r^{-3}) \) medium with constant wind parameters.

The volume of the shocked region (the most likely to provide target protons) can be written approximately as \( V = \beta 4\pi R_{sh}^2/3 \). Here \( \beta \sim 0.3 \) or 0.5 accounts for the fact that the volume of the shocked region from which the emission arises, between the forward and reverse shocks, is smaller than the volume of the entire SNR.

Putting all the above back in eqn \( 7 \) we get:

\[
F_{\gamma}(>E_\gamma,t) = \frac{3q_\gamma B^2 \xi (\kappa C_1)^{5-2s} m_3^3}{2(3-s)(5m-3s-2)} \beta 4\pi m_p d^2 \int_0^t t^{5m-3s-2} dt \tag{7}
\]

where \( \mu \) is the mean molecular weight and \( m_p \) the proton mass. This formula relates the hadronic emission from the SNR to the SNR properties. Since we do not generally expect that young SNRs will be interacting with dense clouds, the target mass used here refers only to the mass of the material swept-up by the SNR shock wave.

### 3.1 Core-Collapse SNRs

Core-collapse SNe arise from the explosion of massive stars, with initial mass generally greater than 8 \( M_\odot \). These stars lose a considerable amount of mass, modifying the medium around them, and forming wind-blown bubbles (Weaver et al. 1977). The structure of these bubbles, and the evolution of SNe within them, has been discussed by many authors (Chevalier & Liang 1989; Tenorio-Tagle et al. 1990, 1991; Dwarkadas 2005, 2007b, 2008). The basic structure of a wind-blown bubble, going outwards in radius from the star, consists of a freely expanding wind region ending in a wind-termination shock, a shocked wind region, a contact discontinuity separating the shocked wind from the shocked ambient medium, an outer shock and the unshocked ambient medium (van Marle et al. 2006; Toal´ a & Arthur 2011; Dwarkadas & Rosenberg 2013). The crucial point here is that close in to the star, the SNR will expand in a wind region. In the simplest approximation, if the wind parameters are constant, the wind density will decrease outwards in radius as \( r^{-2} \), i.e. with \( s=2 \). The parameter \( B = M/(4\pi w v) \) where \( M \) is the wind mass-loss rate and \( w \) is the wind velocity. The values of \( M \) and \( w \) can vary widely depending on the progenitor star, and the phase of evolution.

If we put \( s = 2 \) in eqn \( 7 \) we get, for a SNR evolving in a wind medium:

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Thus for a SNR evolving in a wind medium, the hadronic emission is always decreasing in time (after an initial period where the maximum energy is reached (see for e.g. Tatischeff 2009, Dwarkadas et al. 2012)). The reason for this can be identified from the fact that the swept-up mass is increasing only as a factor of $R \propto t^t$, whereas the energy density of cosmic rays is here taken to be decreasing as $t^{-2}$ (assuming $\xi$ to be a constant).

In general the ejecta profile $n$ is expected to vary between 9 and 12 (Matzner & McKee 1999). Thus the expansion parameter $m$ will lie between 0.85 and 0.9, and therefore the emission decreases at a rate between $t^{-1.15}$ and $t^{-1.1}$.

As can be seen from equation 8 the emission due to pion decay is a function of the wind mass-loss rate divided by the wind velocity. This quantity can vary by over two orders of magnitude depending on the SNR progenitor star, and thus at the same age the hadronic emission can vary by up to 4 orders of magnitude depending on the progenitor star. It will be highest for stars that have a high mass-loss rate and low wind velocity. Stars that fall into this category would be Red Supergiant (RSG) stars, the progenitors of Type IIP, and perhaps Ib, SNe. Since they can have mass-loss rates as high as $10^{-4} M_\odot \, yr^{-1}$, and wind velocities of order 10 km s$^{-1}$, the value of $M/v_w = 6.35 \times 10^{15}$. On the other hand, Wolf-Rayet (W-R) stars can have slightly lower mass-loss rates of about $10^{-5} M_\odot \, yr^{-1}$ and wind velocities of order 2000 km s$^{-1}$, leading to $M/v_w = 3.17 \times 10^{12}$. Although there will be some differences in the other parameters, this potentially could lead to more than 5 orders of magnitude difference in the emission. Thus we would expect that Type IIP SNe, which arise from RSG progenitors, would have the largest hadronic luminosity. The caveat though is that since RSG winds have a small velocity, the RSG wind region cannot extend far out from the star. For a RSG stage that last about $2 \times 10^5$ years, and a wind velocity of 10 km s$^{-1}$, the region would extend only just over 2 pc.

The results indicate that, as the emission is decreasing with time, the best time to observe young core-collapse SNe evolving in winds would be early on. Very early observations in the TeV range can be attenuated by the pair production process (Tatischeff 2009), but according to that paper it would only be important in the first year or so.

### 3.2 Type Ia SNRs

Many of the better known young SNRs that have been observed in gamma-rays appear to be of Type Ia. These include Tycho, SN 1006, Kepler and RCW 86. Type Ia SNRs arise from the thermonuclear deflagration of white dwarfs, and are not expected to considerably modify their medium (although see Williams et al. 2011 for the case of SNR RCW 86). Thus as a first approximation we may assume that they evolve in a constant density medium, and $s = 0$.

Unfortunately, although the power-law density profile assumed for core-collapse SNe has also been used to describe Type Ia’s, it is not a good representation. Dwarkadas & Chevalier (1998) showed that an exponential density profile much better represents profiles of Type Ia SNe calculated from Type Ia explosion models. A power-law solution with $n = 7$, which has often been used, was shown by Dwarkadas & Chevalier (1998) to not correctly represent the density and temperature distribution of the material in the shocked region. Use of an exponential adds one more variable, and therefore the resultant solution is not self-similar, and not analytically tractable. Thus it is not possible to employ an exponential profile and derive a solution analogous to the one above.

As shown in Dwarkadas & Chevalier (1998), while there is some difference in the radius and velocity of the outer shock calculated with a power-law as compared to an exponential density ejecta, the difference is much larger for the reverse shock parameters. If we consider only the outer shock and wish to calculate the radius, energy and mass of shocked material as we need here, the power-law profile may be used to give an answer possibly correct to order of magnitude or somewhat better, keeping the above caveats in mind.

If we set $s = 0$ in equation 7 we get:

$$F_\gamma (> E_o, t) = \frac{3q_p B^2 (\xi \kappa C) m^3}{2(5m - 2) \beta \mu m_d^2 t^{m-2}}$$

Thus for a SNR evolving in a wind medium, the hadronic emission is always decreasing in time (after an initial period where the maximum energy is reached (see for e.g. Tatischeff 2009, Dwarkadas et al. 2012)). The reason for this can be identified from the fact that the swept-up mass is increasing only as a factor of $R \propto t^t$, whereas the energy density of cosmic rays is here taken to be decreasing as $t^{-2}$ (assuming $\xi$ to be a constant).

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The results indicate that, as the emission is decreasing with time, the best time to observe young core-collapse SNe evolving in winds would be early on. Very early observations in the TeV range can be attenuated by the pair production process (Tatischeff 2009), but according to that paper it would only be important in the first year or so.

### 3.3 SNRs in Wind Bubbles

As mentioned above, the surroundings of SNRs are much more complicated. In general core-collapse SNRs would first expand in a wind medium, then encounter a wind termination shock followed by a more or less constant density medium. In the case of a Type Ic SN the wind density at the termination shock would likely be a factor of 4 lower than that of the shocked constant density wind medium; thus the SNR is subsequently evolving into a higher density medium. In the case of a HP/Ib SN, the density in the wind could be much higher than that in the constant density medium, and the SN is evolving from a high density to a lower density medium.

It would be tempting to assume that while the SNR is expanding in either the freely expanding wind medium or the shocked wind medium, the above expressions could be used to describe it. The difficulty is that for a time which can be several doubling times of the radius, as the SN shock crosses the wind termination shock or a discontinuity, the SNR is in a transition stage, and no longer able to
be described by the self-similar solutions given above. This is much more accentuated in the case where the jump in density from the wind to the main-sequence shocked bubble is large. This can happen in Type IIP SNe, when crossing from the RSG wind to the main-sequence bubble. The shocked structure then differs from the self-similar solution (see Telezhinsky et al. 2013), limiting the applicability of the above formulae. However, we would expect that although the actual radius evolution with time cannot be easily described, the general trend regarding the increase or decrease of hadronic emission would still hold.

In cases where the density jump, and transition time, may be smaller, say Type Ic SNe, the solution may reach the self-similar structure quickly after the transition, and the above results may be more easily attained. Therefore we would expect the hadronic emission to first decrease with time as the SNR evolves in the wind region, then undergo a period of transition, and then gradually start to increase in time. This trend is also seen in the numerical calculations of Caprioli (2011). If the SNR shock goes on to collide with the dense shell surrounding the wind-bubble, or comes close enough to it, the accelerated particles may impact the dense shell, giving a further increase in the hadronic emission.

3.4 SNRs in Time-Dependent Winds

In §3.1 we assume that the wind which the SNR was evolving in had constant wind parameters, namely the mass-loss rate and wind velocity. There is of course no good reason why the wind parameters should be constant. As stars evolve on the main-sequence and beyond, their mass-loss properties change with time. Especially towards the end of their lifetime, the surroundings of many observed SNe have indicated the presence of time varying winds (McCray 1993; Chugai et al. 2004; McCray 2005; Smith et al. 2008; Dwarkadas et al. 2010). The X-ray light curves of young SNe suggest that many of them evolve in winds that do not appear to have constant wind parameters (Dwarkadas & Gruszko 2012). The further out in radius one goes, the more likely it is that the wind parameters may not be constant. In practice this means that the circumstellar density will not decrease as r⁻².

In principle, if the winds are time-dependent, then we have to go back to equation 2 and insert the time-dependent part into the integral and re-calculate. To estimate the effect of time-dependent winds, one could assume that the time dependence for instance somehow conspired to give a value of s intermediate between the wind and constant density case. For a sample case s = 1, we can write:

\[
F_γ(>E_\gamma,t) = \frac{3q_γB^2(\kappa C_1^3 m_3^3}{4(4m - 2)βμmdt^{3m-2}} \times t^{3m-2} \tag{11}
\]

Note that in this case the emission is increasing for m > 2/3 and decreasing for lower values of m as the remnant decelerates. The important point is that it may no longer be monotonically increasing or decreasing throughout the ejecta-dominated stage, but reaches a maximum and then begins to decrease.

3.5 Reverse Shocked Emission:

The above analysis mainly considered particles accelerated at the forward shock. In young SNRs that have not reached the Sedov or adiabatic stage, a reverse shock exists that expands back into the ejecta. Whether particles can be accelerated at the reverse shock or not is an ongoing question. There have been several suggestions of non-thermal emission arising from the reverse shock (DeLaney et al. 2002; Rho et al. 2002), the most convincing one being for the case of Cas A (Helder & Vinh 2008). Since we know that the reverse shock is also a collisionless shock and must be mediated by magnetic fields, one cannot deny the existence of a magnetic field, however small, at the reverse shock. The question then is whether the field can be amplified enough to accelerate particles. In a series of papers, Telezhinsky et al. (Telezhinsky et al. 2012a,b, 2013) have computed the acceleration of particles in SNRs, assuming acceleration at both forward and reverse shock. They show that due to the larger number of particles but lower maximum energy at the reverse shock, the resultant integrated spectrum from the SNR will be steeper (softer) at TeV energies than at GeV energies, which is consistent with the observations.

If particles are accelerated at the reverse shock, it may be especially important in situations where the reverse and forward shocks are expanding in regions of different densities. This may occur for instance when the SNR crosses a discontinuity or wind termination shock, such as in wind bubbles. The reverse shock may be in a region with different density compared to the forward shock, and may even dominate the emission.

In the above calculations it is possible to take the reverse shocked emission into effect, or to assume emission from both forward and reverse shocked plasma. The calculations are simplified in the thin shell approximation, where the region between the forward and reverse shocks is assumed to be small compared to the radius of the remnant. This requires a few changes in equation 2. For κ we substitute κᵣ < 1 for the reverse shock. In the thin shell approximation, the mass swept up by the reverse shock can be written in terms of that swept-up by the forward shock as (Nynmark et al. 2008):

\[
M_{cuv} = \frac{n - 4}{4 - s} M_{cs} \tag{12}
\]

Similarly, the density behind the reverse shock can be written in terms of the density behind the circumstellar shock:

\[
ρ_{cuv} = \frac{(n - 3)(n - 4)}{(3 - s)(4 - s)} ρ_{cs} \tag{13}
\]

With the appropriate changes, we get the hadronic emission from the reverse shock to be:

\[
F_γR(>E_γ,t) = \frac{(n - 4)^3(n - 3)}{(4 - s)^2(3 - s)^2} \times \frac{3q_γB^2(\kappa_5C_1^3(1 - m)^3)}{2(5m - ms - 2)βμmdt^{5m - 2ms - 2}} \times t^{5m - 2ms - 2} \tag{14}
\]

If the thin-shell solution is not applicable, then one needs to go back to equation 2 and work through, while inputting the actual reverse shock parameters.

4 APPLICATION TO INDIVIDUAL SNRS

The validity of these solutions can be gauged by computing the emission for observed SNRs. Unfortunately, there are
not many young core-collapse SNRs that have been observed at GeV and TeV energies. As pointed out above, many of the observable ones appear to be of Type Ia, although SN surveys suggest that core-collapse ones should be more common. This could be interpreted in the context of our above results as saying that since the emission from core-collapse SNe decreases with time, one would be less likely to see them after several hundreds of years, unless they are expanding in a dense wind.

4.1 Cas A

One of the SNRs that is assumed to be expanding in a dense wind, and has been detected at high energies, is Cas A. Chevalier & Oishi (2003) suggest that the SNR is expanding in a RSG wind, with a wind mass-loss rate of $2 \times 10^{-9} M_\odot yr^{-1}$ for a wind velocity of 10 km s$^{-1}$. They infer the energy of the explosion to be about 4 times the standard energy $(4 \times 10^{51}$ ergs), and the ejecta to have a density profile that goes as $r^{-10.12}$, in accordance with the work of Matzner & McKee (1999) for a massive star with a radiative envelope. This gives $m \sim 0.88$. Using these parameters (with $n = 10$), we can calculate the value of $C_1$ to be $1.165 \times 10^{10}$, and take $\kappa \sim 1.2$, and $\beta = 0.5$. We assume the fraction of energy transferred to cosmic rays to be 0.1. The value of $q_e > 100$ MeV is obtained from DAV94 to be $0.5 \times 10^{-13}$, almost independent of the spectral index $\alpha$ at low energies. We take $\mu = 1.4$. Using these values we get from equation (10) that

$$F_{\gamma \text{Cas A}}(> 100 \text{ MeV}) = 5.5 \times 10^{-8} \text{ cm}^{-2} \text{ s}^{-1}$$

This is about 6 times larger than the flux reported by Fermi (Abdo et al. 2010), which is the flux above 500 MeV, and so may be expected to be slightly lower. However, the Fermi best fit suggests that at most 2% of the total energy has gone into cosmic rays. If we assume $\xi = 0.02$ then we get a decrease of a factor of 5, giving a flux $F_{\gamma \text{Cas A}}(> 100 \text{ MeV}) = 1.1 \times 10^{-10}$ that is comparable to the Fermi result. Furthermore, the Fermi result assumes a total energy half that of Chevalier & Oishi (2003), which would reduce the flux further, although not exactly by a factor of 2 since other parameters would also change. However this exercise suggests that within the error bars, the flux computed via this method is close to that detected by Fermi, and therefore suggests that the $\gamma$-ray emission from Cas A could be due to hadronic processes. A similar inference was made in Abdo et al. (2010).

4.2 SN 1993J

SN 1993J is one of the closest SNe to have exploded in the past couple of decades, and one of the most observed in X-ray and radio. It is not surprising that several papers have attempted to model the $\gamma$-ray emission from SN 1993J.

\footnote{In actual fact, Chevalier & Oishi (2003) use the density profile for a red supergiant star with a radiative envelope given in Matzner & McKee (1999), which is somewhat different from that assumed here. That profile has a density power-law density that starts off as steep and slowly becomes less steep. This will cause some difference in the results.}

Kirk et al. (1995), Tatischeff (2009), Tatischeff (2009) made a very detailed computation of the particle acceleration, radio emission and hadronic emission from SN 1993J. He finds (eqn. 53) that the emission goes as $t^{-1}$. This is primarily because he assumes that $\xi_{CR}$, the ratio of the cosmic-ray pressure to gas pressure, goes as $V_{sh}^{-1} \propto t^{1-m}$, and therefore the cosmic ray energy density decreases as $t^{-(m+1)}$. As shown above, in the case of core collapse SNs, we have assumed that the cosmic ray energy density decreases as $t^{-2}$, the same as the internal energy density and therefore the gas pressure. Thus the difference between the time evolution in the two cases is due to the fraction of energy going into cosmic rays, and its constancy with time.

We can compute the $\gamma$-ray flux from SN 1993J, using the parameters assumed by Tatischeff (2009), namely $m=0.83$, $B=1.9 \times 10^{14} \text{ g cm}^{-1}$, $d=3.63 \text{ Mpc}$, and $\kappa C_1 = 2.79 \times 10^{10}$. We then get from equation (10), assuming $\mu = 1.4$, that the flux from SN 1993J in the TeV range goes as

$$F_{\gamma \text{93J}}(> 1 \text{ TeV}) = 2.14 \times 10^{-11} t^{-1.17}_{\text{days}} \text{ cm}^{-2} \text{ s}^{-1}$$

The constant in front is about an order of magnitude larger than in Tatischeff (2009), but the flux is decreasing somewhat faster, as $t^{-1.17}$ rather than $t^{-1}$.

The flux greater than 100 MeV is then:

$$F_{\gamma \text{93J}}(> 100 \text{ MeV}) = 1.07 \times 10^{-7} t^{-1.17}_{\text{days}} \text{ cm}^{-2} \text{ s}^{-1}$$

After 20 years the flux in the Fermi range is $3.23 \times 10^{-12} \text{ cm}^{-2} \text{ s}^{-1}$, which is below the level that can be detected by Fermi.

It is not clear however that the density of the medium around SN 1993J does decrease as $r^{-2}$, and that a constant mass-loss rate is appropriate. Early reports (van Dyk et al. 1994) inferred a density decline of $r^{-1.5}$ from the radio emission, and $r^{-1.7}$ from the X-ray emission (Suzuki & Nomoto 1993). A circumstellar medium density decreasing as $r^{-1.5}$ was used by Kirk et al. (1993) to compute the emission from SN 1993J. Inspired by these results, Mioduszewski et al. (2001) modelled the hydrodynamic interaction, and found that the radio emission was better explained by a complicated density profile that diverged from $r^{-2}$. On the other hand, detailed calculations of the radio emission, encompassing much more physics, were carried out by Pransson & Björnsson (2003), who suggested that the radio could be explained by the self-similar solution embodying shock expansion into a steady wind, a result that was echoed by Tatischeff (2009). However, Nymark et al. (2009) then found that the self-similar solution was not adequate to explain the X-ray emission, and used the ejecta profile in Suzuki & Nomoto (1993). Interpretation of the data does not lead to conclusive results (Bietenholz et al. 2010), further complicating the situation.

In this context it is interesting that the hadronic emission, and its time evolution, differ considerably depending on whether the medium goes as $r^{-2}$ or $r^{-1.5}$. One may speculate that in future, if nearby SNe could be detected and followed with advanced instruments at very high energies, they may be combined with other multi-wavelength results to help infer the nature of the medium that SNRs are expanding in, and thereby the nature of mass-loss from the progenitor star.
4.3 SN 1987A

Even though SN 1993J happens to be one of the closer core-collapse SNe of the modern era, the large distance to the SN renders the flux small enough to be practically undetectable with current telescopes. It is therefore opportune to turn our attention to the closest SN in modern times, SN 1987A, which exploded in the LMC at a distance of about 50 kpc. The ambient medium into which the SN shock is expanding is distinctly non-uniform, as revealed by X-ray, optical, and radio observations, and quite unlike other core-collapse SNe that we have discussed. In order to explain the increasing radio and X-ray emission, Chevalier & Dwarkadas (1992) suggested that, after about 3 years of expansion in a wind medium, the SN shock impacted a dense HII region formed by the progenitor star and ionized by the pre-SN star. Finally (after about 25 years or so), the SN shock will begin to interact with the dense equatorial ring of material that is seen so beautifully in optical Hubble Space Telescope images. The density profile into which the shock is expanding is thus quite inhomogeneous. However, and since we are mainly interested in the mass of material that has been shocked, we can obtain a rough estimate of the flux by neglecting the mass of the swept-up wind, and assuming that up to about 26 years the SN shock is interacting with a constant density HII region with mean hydrogen density of about 200, which gives approximately the correct mass. We use equation (10) to compute the flux from the SN, assuming the parameter $n = 9$, and ejecta parameters as used in Dewey et al. (2012). The flux greater than 100 MeV, at an age of 26 years, is then:

$$F_{1987A} (> 100 \text{ MeV}) = 4.04 \times 10^{-10} \text{ cm}^{-2} \text{ s}^{-1}$$  (18)

The observed synchrotron power-law index of SN 1987A, as measured in the radio, is very soft (Zanardo et al. 2013) and varies over the remnant, with values 2.4 to 2.8. At TeV energies, as shown by Drury et al. (1993), the emissivity $q_\gamma$ can vary by over two orders of magnitude as the spectrum steepens from a power-law index of 2.1 to 2.6. However, at the shock front, within the assumption of Bohm diffusion, we would expect the spectrum at very high energies to be close to 2. In fact non-linear DSA predicts a harder spectrum. The observed spectrum may be softened due to other effects such as Alfvénic drift (Zirakashvili & Ptuskin 2008; Caprioli 2011). Therefore, we assume a spectral index at the shock close to 2, and a gamma-ray emissivity of $q_\gamma \approx 1 \times 10^{-17}$, which gives:

$$F_{1987A} (> 1 \text{ TeV}) = 8.1 \times 10^{-14} \text{ cm}^{-2} \text{ s}^{-1}$$  (19)

This is a few times smaller than that calculated by Berezhko et al. (2011). This is to be expected given the difference in their assumptions about the surrounding medium, and the nature of our estimate. In fact it is reassuring that they are comparable. One of the problems with our approach is in fact that since the SN is first evolving in a wind and then in a circumstellar medium, as pointed out earlier, the shock structure is going to be in a transition state for the first few years after impact with the HII region (Dwarkadas 2007; Dewey et al. 2012). Nevertheless, this does provide us with an order of magnitude estimate of the hadronic flux. Note that in our case the flux is increasing with time as $t^{1.33}$.

The hadronic flux level of SN 1987A is at present undetectable. However, it is exciting to note that the SN will soon, if it has not already, start interacting with the dense equatorial ring formed by the progenitor star, with density $n_H \sim 10^4$ (Lundqvist 1999). Then, depending on the thickness of the ring, the $\gamma$-ray flux is expected to increase significantly over the next decade. We can approximate the luminosity at 36 years assuming an average density of about $n_H \sim 2 \times 10^3$ into which the SN shock has been expanding, again keeping in mind that the self-similar solution may not be entirely appropriate. We then get for the flux at 36 years to be:

$$F_{1987A} (> 100 \text{ MeV}) = 1.9 \times 10^{-8} \text{ cm}^{-2} \text{ s}^{-1}$$  (20)

and

$$F_{1987A} (> 1 \text{ TeV}) = 3.8 \times 10^{-12} \text{ cm}^{-2} \text{ s}^{-1}$$  (21)

The latter is detectable with an instrument such as the current HESS telescope array in less than an hour. Even if the number is off by a factor of 5, which is quite possible given the inhomogeneous and complex nature of the surrounding medium, it should still be detectable within a few hours. In fact it is possible that a 25 hour observation could detect it a few years earlier. All in all, even these order-of-magnitude estimates suggest that SN 1987A should be kept in mind as a potential gamma-ray source in the next decade, whose detection would strongly support the model for cosmic-ray acceleration in SNRs.

SN 1987A is unusual in that it is a core-collapse SN whose hadronic luminosity is expected to increase with time early on, whereas for most SNe it is decreasing with time. This had previously been pointed out, albeit with different time dependence due to different assumptions of velocity and density, by Kirk et al. (1995). The increasing flux in the radio and X-ray regimes has been monitored for the past two decades, and detection at very high GeV-TeV energies would serve to corroborate our understanding of both the hadronic emission and the details of the SN environment. Unfortunately, the Fermi space telescope will probably not be active by the time the SN brightens sufficiently, but ground based ACTs such as HESS should be able to detect it in the TeV range. It also promises to be an exciting target for the upcoming Cherenkov Telescope Array, which, if our estimates are good to even an order of magnitude, should definitely detect it.

5 DISCUSSION:

The above results indicate that if we take a statistical ensemble of young core-collapse SNRs with age, one would see a gradual decrease in the hadronic emission with age, varying between $t^{-1.1}$ and $t^{-1.5}$, but with a scatter that extended to about 2 orders of magnitude in each direction. Hp’s arise from progenitor stars that are much lower in mass than 1b/c’s, and are larger in number, given the weighting of the initial mass function. In theory they would be expected to dominate, thus leading to a high luminosity that decreases more or less as $t^{-1.1}$. Unfortunately, because of the very few young SNRs that are observed in our galaxy, the small-number statistics do not present a true picture. Furthermore, even among those that are seen, many of them are...
Type Ia, which should show hadronic emission that increases with time. We would expect therefore that a plot of luminosity v/s age for SNRs would initially decrease (due to the predominance of core-collapse SNe) but then begin to increase as the Type Ia’s become brighter, and the core-collapse ones begin to expand in a constant density medium. However, unless the statistics can be improved by future telescopes, perhaps detecting young SNRs such as SN 1987A in galaxies outside our own, it will be hard to get a good statistical description.

Kirk et al., (1995) calculated the evolution of the hadronic photon luminosity for two SNe. Their results can easily be understood in terms of the expressions derived herein. They assumed that in the ejecta-dominated phase, the shock velocity is a constant (m = 1). That implies from equation [4] that the flux will decrease as $F_s \propto \rho^{-1.5} v^2$. For a SN expanding in a wind, as they assumed for SN 1987A, with $s = 3/2$ we then get that the flux decreases as $t^{-1}$, as they found. In the case of SN 1993J, they assumed it to be expanding in a medium whose density goes as $\rho \propto r^{-1.5}$, i.e. $s = 3/2$, in which case, with $m = 1$, we get that $F_s \propto t^{3.5}$, i.e. the flux is a constant with time. Thus our results reduce to theirs in the appropriate parameter range, further emphasizing their applicability and versatility.

In this context it must be mentioned that by taking the quantity $\xi$ out of the integral, we are in effect replacing the instantaneous energy loss to cosmic rays with the total energy lost over a finite period of time, thus making $\xi$ a global efficiency parameter, and a proxy for the total energy lost to cosmic rays with the total energy. Thus our estimates are reasonable in calculating the luminosity or flux good to a factor of a few. (Kirk et al., 1995) carried out a highly detailed exploration of the particle acceleration in SN 1993J. The differences between our simpler estimate and his detailed computation can be understood in terms of the differences in the evolution of the cosmic ray pressure, and suggest that our estimates are reasonable in calculating the luminosity or flux good to a factor of a few.

5.1 Energy in Cosmic Rays

A question that naturally arises is whether the self-similar solution used herein would be modified due to the particle acceleration. Many authors have shown that the shocked region will contract further as more and more kinetic energy is diverted to cosmic-ray acceleration (Blondin & Ellison 2003; Ellison et al. 2003; Ferrand et al. 2012), and the density jump across the shock will change. However, as long as the ratio of cosmic-ray to gas pressure does not exceed 10%, the test particle solutions, and the unmodified shock structure, should be reasonably correct (Kang, 2013). This was addressed by Chevalier (1983), who showed that while there will be some effect on the radii and densities, it will be only at a few percent level if 10% of the kinetic energy goes into accelerated particles. (As an aside, we note that although the energy expended in cosmic ray acceleration is calculated differently in Chevalier (1983), the resulting time-dependence, although expressed in terms of $n$ and $s$ rather than $m$, is exactly the same as reported in this paper.) Thus, using $\xi = 0.1$, it is reasonable to assume that the shock structure is reasonably well represented by the self-similar solutions.

The question that remains then is whether the cosmic-ray pressure exceeds the gas pressure by 10%, or equivalently whether a large amount of energy is expended in accelerating cosmic rays. Unfortunately there is no easy answer to that. A simple calculation (for e.g. Longair (2011)) shows that about 10% of the total energy in SNRs is required to fuel the galactic cosmic-ray flux. Therefore it does not seem that a much higher amount is necessary to explain the observed cosmic-ray flux. Analysis of SNR evolution, including particle acceleration, provide conflicting numbers. For Cas A, Patnaude & Fesen (2009) suggest that about 30% of the energy goes into accelerating cosmic rays, while the broad-band fits by the Fermi team find that the energy content in cosmic rays is less than 2% of the total explosion energy. In the case of Tycho’s SNR, Kosenko et al. (2011) find that about 5-20% of the energy is lost in cosmic-ray escape; Morlino & Caprioli (2012) suggest that 5-10% of the explosion energy goes into cosmic-rays. Zhang et al. (2013) suggest that the energy conversion efficiency is only 1% but that molecular-cloud interaction is involved; Berezhko et al. (2013) require a two-phase inhomogeneous medium; while Atoyan & Dermer (2012) require a very small efficiency but show that using a multi-zone model, lepton emission remains a possibility to explain the very high energy emission. Therefore, in the case of Tycho’s SNR, no consensus exists on even the emission mechanism, let alone the energy needed. In the case of SN 1006, Berezhko et al. (2012) find that the energy in accelerated cosmic rays is about 5% of the total explosion energy, although this could double over time, while Kosenko et al. (2011) find that energy losses due to cosmic ray escape are 20-50% of the kinetic energy flux.

Studies of other SNRs are equally conflicting. Tatischeff (2009) found that about 19% of the energy processed by the shock goes into cosmic-ray energy over the first 8.5 years, but yet the shock is only weakly modified. This can perhaps be understood by the fact that the energy processed by the shock is only a small percentage of the total kinetic energy available, so the fraction of the cosmic-ray energy to total energy is much smaller. Similarly, Ellison et al. (2010) find that the instantaneous energy that goes into cosmic-rays can be as high as 25-50% for SNR RX J1713, but that the overall energy is about 13% for the leptonic model that they favor.

Another observation supposedly pointing to efficient cosmic ray acceleration is the ratio of the forward shock (blast wave) to the contact discontinuity. For e.g. in Tycho’s SNR, Warren et al. (2003) find that the ratio of the blast wave to the contact discontinuity is 0.93, which is much larger than that expected from adiabatic hydrodynamic models, and suggests that the blast wave is much closer to the contact discontinuity than expected. They take that as a sign that the region has been compressed due to efficient cosmic-ray acceleration. However, this presumes that the shock wave has evolved over a constant density medium throughout its lifetime, which is not certain. Furthermore, Orlando et al. (2012) clearly show that the smaller separation can be easily reproduced by clumping of the ejecta, and does not require cosmic-ray acceleration. It is significant that the ratio of forward to reverse shock can be reproduced by hydrodynamical models, it is only the radius of the con-
tact discontinuity that seems contrary to expectations, thus further supporting the claim of Orlando et al. (2012).

Do the non-linear DSA models produce spectra that fit the observations better? Since the low-energy particles see a smaller shock jump, and the higher energy particles a larger shock jump, the models actually produce spectra that are steeper at GeV than at TeV energies, which is the inverse of the observations, which are far steeper at TeV energies. However, as shown by Caprioli (2011), this can be modified to match the observations by assuming that the Alfvénic drift of magnetic scattering centers plays a role.

On the other hand, as shown by Telezhinsky et al. (2012, 2013), softer spectra in the TeV regime that match the observations can be obtained even in test-particle mode, by making the simple assumptions that both forward and reverse shocks can accelerate particles. The higher intensity and lower maximum energy of particles accelerated at the reverse shock, when combined with the forward shock accelerated particles, produce integrated particle spectra that are softer at TeV wavelengths than in the GeV region, consistent with the observations.

It seems, at least at present, that there is neither motivation nor compulsion for a large amount of explosion energy (> 10%) to be diverted to cosmic-ray acceleration, nor is there any specific clear observation of the same.

5.2 Circumstellar Environments

As shown above, the hadronic emission goes as $\rho^2$ for young SNRs. Drury et al. (1994) have an expression for hadronic emission with one power of density in it. The second arises from the fact that the energy processed instantaneously at the shock is much smaller than the total SNR energy. This dependence on the square of the density is similar to that for thermal X-ray emission from SN shock waves (see for e.g. Dwarkadas & Gruszko 2012), therefore it is no surprise that hadronic emission is pre-dominant in SNRs which also show a high intensity of thermal x-ray emission.

Caprioli (2011) had modelled the hadronic emission from SNRs interacting with “non-homogeneous circumstellar environments”, basically a wind-blown bubble. That model though has many problems, as is evident by comparing to analytical and numerical models of wind-bubbles (García-Segura et al. 1996; Dwarkadas 2003, 2007b; Toalá & Arthurs 2011; van Marle et al. 2006; Dwarkadas & Rosenberg 2013): (1) Firstly, it assumes a RSG wind close in to the star with a density profile that goes as $r^{-2}$, followed by a W-R bubble with constant density. As shown by many authors, the RSG wind also blows a shell, but due to the much larger velocity and momentum of the W-R wind, the latter completely destroys the RSG shell and pushes the RSG material further out, mixing it up in the process. The RSG and W-R portions do not exist separately as used in Caprioli (2011). In particular, there is no RSG wind at the end of the W-R stage, but a W-R wind, whose density would be orders of magnitude lower.

(2) A RSG wind going directly into a constant density W-R medium without an intervening discontinuity such as a shock is in fact not even a hydrodynamically stable situation, and cannot exist in practice. (3) Caprioli (2011) has a high-density jump at the end of the bubble, after which the density stays constant. In practice W-R bubble shells are bounded by radiative shocks, where the shock jump can be extremely high. Thus there is a dense shell surrounding the bubble, following which the ISM density should be much lower than the shell. The SNR shock may be considerably decelerated at the dense shell. None of these factors are taken into account. (4) The region is arbitrarily divided into ejecta-dominated and Sedov-Taylor. As mentioned above it takes several swept-up masses before the Sedov-Taylor stage is reached. In many cases, the Sedov-Taylor stage will not be reached for evolution in a W-R bubble. It is also very unlikely that the Sedov stage will be reached by a SN expanding in a wind medium, given the decreasing density and the fact that a substantially large mass is needed. (5) The solution given for the evolution of the shock wave in the wind medium by Caprioli (2011) (their equations 3.4 and 3.5) are for a specific value of the parameter “n” (as used here), which is not generally applicable over the entire mass range of progenitors. More specifically, it is not necessarily applicable to W-R stars (Matzner & McKee 1999), to which it was applied. (6) Finally, there is a large variety in wind mass-loss rates and velocities as shown above, up to several orders of magnitude, and therefore a single model with fixed parameters to fit all progenitor stars is a gross oversimplification that will not work for any particular star in practice. We think it is this oversimplification that leads to the bimodal distribution suggested in $\gamma$-ray hadronic emission found in Caprioli (2011).

To be fair, the description of the emission given in this paper is simplified. On the other hand, we have a realistic treatment of the surrounding medium and the SN evolution within this medium, which we believe is a crucial ingredient and essential to understanding the hadronic emission. Unless the SNR evolution in the ambient medium is properly treated, even the most accurate treatment of the particle acceleration and hadronic emission will likely lead to incorrect results, because as pointed out the emission depends crucially on the shock radius, velocity and swept-up mass. Our intention as specified is not to calculate in detail the emission from any given SNR (although it appears to do that reasonably well as shown), but to understand the time-evolution, and the factors it depends on, as well as the role of the circumstellar medium, and to begin to evaluate the statistical properties of an ensemble of SNRs.

6 SUMMARY AND CONCLUSIONS

We have derived a simple expression to compute the hadronic emission from young SNRs, taking the evolution of the SN ejecta into various environments into account, and assuming that the target mass for proton collisions is the material swept-up by the SN shock wave(s). We calculate its dependence on the evolutionary parameters (ejecta energy and mass-loss rate), the density structure of the surrounding medium, and the expansion parameter of the shock wave. Our results show that for core-collapse SNRs expanding into the winds of their progenitor stars, the hadronic emission must decrease in time for winds with constant parameters. On the other hand, for expansion into a medium with a constant density, the hadronic emission is expected to increase with time. In the appropriate parameter regime, these results can reproduce the time evolution of the hadronic lu-

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minosity calculated for many specific SNe in the literature, thus demonstrating their broad applicability.

We have not included the so-called “interacting” SNRs, which are interacting with dense molecular clouds. Most of those are in a much more evolved stage, and not in the ejecta-dominated stage. In principle, if they were to be in the ejecta-dominated stage, it is not difficult to include them. The main difference between the non-interacting remnants discussed here, and the interacting remnants, is the mass of the material with which they are interacting, and its distance from the SNR shock wave. That means that we could start from equation 2 and, instead of the swept-up mass, use the mass of the cloud or interacting material and rederive the equations. A major difference is that, depending on the distance to the cloud, only the highest energy cosmic rays would be able to diffuse towards it (Gabici 2011, Telezhinsky et al. 2012b).

The equations derived here can provide observers with a simple method to compute the hadronic emission from any given young SN or SNR, and thus get a better-than-order-of-magnitude estimate of the hadronic flux with knowledge of a few basic explosion parameters. Conversely, knowing that the emission is hadronic, they could be used to derive the density of the ambient medium or the SN parameters.

In future papers we plan to derive similar expressions for the leptonic emission. Our goal is to have a suite of formulae that will provide an aid to understanding the emission from young SNRs in the age of Fermi and ground-based IACT’s, allow us to estimate the dependence on various parameters, and therefore help to interpret the upcoming Fermi SNR catalog.

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