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The contribution of supernova remnants to the galactic cosmic ray spectrum

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ABSTRACT

The supernova paradigm for the origin of galactic cosmic rays has been deeply affected by the development of the non-linear theory of particle acceleration at shock waves. Here we discuss the implications of applying such theory to the calculation of the spectrum of cosmic rays at Earth as accelerated in supernova remnants and propagating in the Galaxy. The spectrum is calculated taking into account the dynamical reaction of the accelerated particles on the shock, the generation of magnetic turbulence which enhances the scattering near the shock, and the dynamical reaction of the amplified field on the plasma. Most important, the spectrum of cosmic rays at Earth is calculated taking into account the flux of particles escaping from upstream during the Sedov–Taylor phase and the adiabatically decompressed particles confined in the expanding shell and escaping at later times. We show how the spectrum obtained in this way is well described by a power law in momentum with spectral index close to -4, despite the concave shape of the instantaneous spectra of accelerated particles. On the other hand we also show how the shape of the spectrum is sensible to details of the acceleration process and environment which are and will probably remain very poorly known.

1. Introduction

In its original form [16], the supernova remnant (SNR) paradigm for the origin of Galactic cosmic rays (CRs) is based on a purely energetic ground: if ~ 10–20% of the kinetic motion of the expanding shell of a supernova gets converted into accelerated particles, and one accounts for the energy dependent escape time from the Galaxy, SNRs can be the sources of the bulk of Galactic CRs. After the pioneering works on diffusive shock acceleration [DSA, 5,21,9,6]), it became clear that this mechanism is the most promising acceleration process that can be responsible for energy conversion from bulk kinetic motion of a plasma to kinetic energy of charged particles. The DSA naturally leads to spectra of accelerated particles $N(E) \propto E^{-2}$ for strong shocks, not too dissimilar from the ones needed to describe data after accounting for the energy dependent escape time from the Galaxy, with a residence time that scales as $\tau_{esc}(E) \propto E^{-0.6}$.

There are however two main concerns with this simple picture: first, the required acceleration efficiency is not so small that the dynamical reaction of the accelerated particles on the shock can be neglected. Second, if particle scattering is guaranteed by normal interstellar magnetic turbulence alone, the maximum energy of accelerated particles is exceedingly small and the mechanism cannot account for cosmic rays with energies up to the knee ($\sim 3 \times 10^6$ GeV). It was soon understood that this second problem could be mitigated only by requiring CRs to generate the turbulence necessary for their scattering though streaming instability [6,22], a mechanism similar to that discussed by [35] in the context of CR propagation in the Galaxy. This latter point intrinsically makes the acceleration process even more non-linear.

The modern non-linear (NL) theory of DSA allows us to describe particle acceleration at SNR shocks by taking into account (1) the dynamical reaction of the accelerated particles on the system, (2) the magnetic field amplification due to streaming instability, and (3) the dynamical reaction of the amplified magnetic field on the plasma. These effects are interconnected in a rather complex way, so that reaching the knee and having enough energy channelled into CRs are no longer two independent problems. The situation is in fact even more complex given that the evolution of the SNR in time depends on the environment.

A generic prediction of NLDSA is that the spectra of accelerated particles are no longer power laws but rather concave spectra. In the case of extremely modified shocks, the asymptotic shape of the spectrum for $E \gg 1$ GeV is $N(E) \propto E^{-1.2}$ (see e.g. [19,23] for reviews on CR modified shocks) to be compared with the standard E^{-2} spectrum usually associated to DSA. Instead of clarifying the situation, this bit of information made it more puzzling in that so flat spectra are hard to reconcile with the CR spectrum observed at Earth.

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In this paper we show how the application of NLDSA to SNRs leads to time-integrated spectra that are very close to power laws at energies below 10–100 TeV where most measurements of CR spectra are performed with high statistical significance. The crucial piece of physics to connect the acceleration process inside the sources to the spectrum observed at Earth is the escape flux: during the Sedov–Taylor phase of the evolution (and to a lesser amount also during the ejecta dominated phase) particles can escape from a SNR in the form of a spectrum peaked at the maximum momentum reached at any given time. Particles which do not escape are advected downstream, lose energy adiabatically and eventually escape at later times. We calculate the spectrum injected by a single SNR as the superposition of these two components under different assumptions. Indeed, the semi-analytical method adopted here not only allows for a complete treatment of NLDSA but, being computationally very cheap, also allows for a very wide scan of the parameter space and an unprecedented investigation of the poorly known pieces of physics that enter the problem.

For simplicity, here we focus on type I supernovae, which occur in the typical interstellar medium (ISM), while qualitative differences between these and type II supernovae are only discussed by considering expansion in a more rarefied, hotter ISM, but totally ignoring any spatial stratification of the circumstellar region, which might be characterized by winds, bubbles and other complex structures.

We also limit our attention to the proton component, while the results on nuclei will be presented in an upcoming paper since the additional issues that appear in that case deserve a detailed discussion. The introduction of nuclei is a fundamental step in the field and is essential to explain the CR spectrum above the knee (see e.g. [4,8] for a review).

2. A back-of-the-envelope calculation of the escape flux from a SNR

The escape of cosmic rays from a SNR is a very difficult problem to tackle, both from the physical and mathematical point of view. One can envision that at some distance upstream of the shock the particle density (or current) gets sufficiently small that the particles are no longer able to generate the waves that may scatter them and lead to their return to the shock front. These are escaping particles. However, the location of this free escape boundary is not easily calculated from first principles and it is usually assumed to be a given fraction of the radius of the shock. An additional uncertainty is introduced by the fact that the shock dynamics changes in time. The evolution of a SNR is characterized by three phases: an ejecta-dominated phase, in which the mass of material accumulated behind the blast wave is less than the mass of the ejecta; a Sedov–Taylor (ST) phase, that starts when the accumulated mass equals the mass of the ejecta; a radiative phase, when the shock dissipates energy through radiation. The SNR is expected to spend most of the time over which it is active as a CR factory in the ST phase, that typically starts 500–1000 years after the initial explosion.

The maximum momentum of accelerated particles during the ejecta-dominated phase is expected to increase with time [22]. As discussed by Caprioli et al. [13], this is due to the fact that magnetic field amplification is rather efficient and the shock speed stays almost constant during this stage. After the beginning of the ST phase, the shock velocity, and thus also the efficiency of magnetic field amplification, decrease with time: as a consequence, the maximum momentum, p_{max}, is expected to drop with time as well [13]. The process of particle escape from the upstream region becomes important. At any given time, the system is no longer able to confine the particles that were accelerated to the highest energies at earlier times, so these particles escape from the shock. The instantaneous spectrum of the escaping particles at any given time is very much peaked around $p_{max}(t)$ [14]. This qualitative picture of particle escape is the one that we mimic by assuming the existence of a free escape boundary, but as stressed above, the escape phenomenon is likely to be much more complex than suggested by this simple picture.

Before embarking in a detailed calculation including the nonlinear effects, it is useful to illustrate the results of a back-of-the-envelope calculation, based on a test-particle approach. Let us consider a SNR shell with a time dependent radius $R_{sh}(t)$ expanding with velocity $V_{sh}(t)$ in a uniform medium with density ρ_0 and suppose that escaping particles have momentum $p_{max}(t)$ and carry away a fraction F_{esc} of the bulk energy flux $\frac{1}{2}\rho_0 V_{sh}(t)^3$. Let $N_{esc}(p)$ be the spectrum of cosmic rays inside the remnant, so that the energy contained in a range dp around p is

$$d\mathscr{E}(p) = 4\pi p^2 N_{esc}(p) pc\, dp. \tag{1}$$

The energy carried away by particles escaping in a time interval dt at time t is

$$d\mathscr{E}(t) = F_{esc}(t)\frac{1}{2}\rho V_{sh}^3(t) 4\pi R_{sh}(t)^2\, dt. \tag{2}$$

In a general way we can write $R_{sh}(t) \propto t^{\nu}$, and thus $V_{sh}(t) \propto t^{\nu-1}$. Using these time-dependencies, and equating the two expression for d$\mathscr{E}$, one obtains

$$N_{esc}(p) \propto t^{5\nu-3} F_{esc}(t) p^{-3}\frac{dt}{dp}. \tag{3}$$

During the ST stage, p_{max} is determined by the finite size of the accelerator, therefore we require that the diffusion length $\lambda(p)$ at $p_{max}(t)$ is a fraction χ of the SNR radius (free escape boundary):

$$\lambda(p_{max}) \simeq D(p_{max})/V_{sh} = \chi R_{sh}. \tag{4}$$

Assuming for the diffusion coefficient the generic form $D(p) \propto p^{\alpha}/\delta B^{\gamma}$ and, for a magnetic field scaling as $\delta B(t) \propto t^{-\mu}$, we obtain

$$p(t)^{\alpha} \propto R_{sh}(t) V_{sh}(t) \delta B(t)^{\gamma} \propto t^{2\nu-1-\gamma\mu}, \tag{5}$$

which implies

$$\frac{dt}{dp} \propto \frac{t}{p}. \tag{6}$$

Substituting Eq. (6) into Eq. (3) one obtains:

$$N_{esc}(p) \propto p^{-4} t^{5\nu-2} F_{esc}(t); \quad t = t(p). \tag{7}$$

This relation illustrates a striking result: if the fraction of the bulk energy going into escaping particles is roughly constant in time, and if the SNR evolution during the ST stage is adiabatic and self-similar (i.e. $\nu = 2/5$), the global spectrum of particles escaping the system from the upstream boundary is exactly p^{-4}. This means that the diffuse CR spectrum, usually explained by invoking the quasi-universal slope predicted by Fermi's mechanism at strong shocks, may be as well due the equally general evolution of a SNR during the ST stage.

Possible corrections to Eq. (7) might lead to a slightly different spectrum for the escape flux. For instance, if the SNR evolution were not perfectly adiabatic, e.g. as a consequence of the energy carried away by escaping particles ($\nu \leqslant 2/5$) or if F_{esc} decreased with time (corresponding to a reduction of the shock modification), the spectrum of the escaping particles could be as flat as $\sim p^{-3.5}$. Reasonable modifications to the basic prediction for the escaping particle spectrum generally lead to spectra that are somewhat flatter than p^{-4}.