Spallative Nucleosynthesis in Supernova Remnants

I. Analytical Estimates

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Abstract. Spallative nucleosynthesis is thought to be the only process capable of producing significant amount of Beryllium (Be) in the universe. Therefore, both energetic particles (EPs) and nuclei to be spalled (most efficiently C, N and O nuclei in this case) are required, which indicates that supernova (SNe) may be directly involved in the synthesis of the Be nuclei observed in the halo stars of the Galaxy. We apply current knowledge relating to supernova remnant (SNR) evolution and particle shock acceleration to calculate the total Be yield associated with a SN explosion in the interstellar medium, focusing on the first stages of Galactic chemical evolution (i.e. when metallicity 0 \leq Z \leq 10^{-2}). We show that dynamical aspects must be taken into account carefully, and present analytical calculations of the spallation reactions induced by the EPs accelerated at both the forward and the reverse shocks following the SN explosion. Our results show that the production of Be in the early Galaxy is still poorly understood, and probably implies either selective acceleration processes (greatly favouring CNO acceleration), reconsideration of the observational data (notably the O vs Fe correlation), or even new energy sources.

Key words: Acceleration of particles; Nuclear reactions, nucleosynthesis; ISM: supernova remnants; Galaxy: abundances

1. Introduction

There has recently been considerable interest in the Galactic evolution of the abundances of the light elements Li, Be and B (Feltzing & Gustafsson 1994; Reeves 1994; Cassé et al. 1995; Fields et al. 1994,1995; Ramaty et al. 1996,1997; Vangioni-Flam et al. 1998). The Be abundance is particularly interesting because this element is thought to be produced exclusively by spallation reactions involving collisions between nuclei of the CNO group of elements and protons or alpha particles at energies greater than about 30 MeV per nucleon (MeV/n). Thus the evolution of the Be abundance contains information about the particle acceleration and cosmic ray history of the Galaxy.

The evolution of the Be abundance, and indeed the evolution of all elemental abundances, has to be deduced from observations of the fossil abundances preserved in the oldest halo stars. Advances in spectroscopy over the last decade have greatly improved the quality of the data available (Duncan et al. 1992,1997; Edvardsson et al. 1994; Gilmore et al. 1992; Kiselman & Carlsson 1996; Molaro et al. 1997; Ryan et al. 1994) and the main result is easily summarized: in old halo stars of low metallicity, the ratio of the Be abundance to the Iron (Fe) abundance appears constant, that is to say the Be abundance rises linearly with the Fe abundance.

This has been a surprising result. Naively one had expected that, because Be is a secondary product produced from the primary CNO nuclei, its abundance should vary quadratically as a function of the primary abundances at low metallicities. Indeed, considering that the cosmic rays (CRs) responsible for the Be production are somehow related to the explosion of supernovae (SNe) in the Galaxy, it is natural to assume that their flux is proportional to the SN rate, \( dN_{SN}/dt \). Now since the number of CNO nuclei present in the Galaxy at time \( t \) is proportional to the total number of SN having already exploded, \( N_{SN}(t) \), the Be production rate has to be proportional to \( N_{SN}(t) dN_{SN}/dt \). Therefore, the integrated amount of Be grows as \( N_{SN}^2 \), that is quadratically with respect to the ambient metallicity (C,N,O or Fe, assumed to be more or less proportional to one another).

The above reasoning, however, relies on two basic assumptions that need not be fulfilled: i) the CRs recently accelerated interact with all the CNO nuclei already produced and dispersed in the entire Galaxy and ii) the CRs are made of the ambient material, dominated by H and He nuclei. Instead, it might be i) that the proton rich CRs recently accelerated interact predominantly with the freshly synthesized CNO nuclei near the explosion site and ii) that a significant fraction of the CNO rich SN ejecta are also accelerated. In both cases, a linear growth of the Be abundance with respect to Fe or O would arise very naturally, since the number of Be-producing spallation reactions induced by each individual supernova would be
directly linked to its local, individual CNO supply, independently of the accumulated amount of CNO in the Galaxy.

In fact, as emphasised by Ramaty et al. (1997), the simplest explanation of the observational data is to assume that each core-collapse supernova produces on average 0.1 \(M_\odot\) of Fe, one to few \(M_\odot\) of the CNO elements and \(2.8 \times 10^{-8} M_\odot\) of Be, with no metallicity dependence. Clearly if this is the case and the production of Be is directly linked to that of the main primary elements, the observed linear relation between Be and Fe will be reproduced whatever the complications of infall, mixing and outflow required by the Galactic evolution models. On the other hand, although the simplest explanation of the data is clearly to suppose a primary behaviour for the Be production, it is possible that this could be an artifact of the evolutionary models (as argued, e.g., by Casuso & Beckman 1997).

Most work in this area has attempted to deduce information about cosmic ray (or other accelerated particle populations) in the early galaxy by working backwards from the abundance observations. While perfectly legitimate, our feeling is that the observational errors and the uncertainties relating to Galactic evolution in general make this a very difficult task. We have chosen to approach the problem from the other direction and ask what currently favoured models for particle acceleration in supernova remnants (SNRs) imply for light element production. This is in the general spirit of recent calculations of the \(\pi^0\)-decay gamma-ray luminosity of SNRs (Drury et al. 1994) and the detailed chemical composition of SN shock accelerated particles (Ellison, Drury & Meyer, 1997) where we look for potentially observable consequences of theoretical models for cosmic ray production in SNRs.

Interestingly enough, the study of particle acceleration in SNRs suggests that both alternatives to the naive scenario mentioned above do occur in practice, as demonstrated qualitatively in Sect. 2. The first of these alternatives, namely the local interaction of newly accelerated cosmic rays in the vicinity of SN explosion sites, has already been called upon by Feltzing & Gustaffson (1994), as well as the second, the acceleration of enriched ejecta through a SN reverse shock, by Ramaty et al. (1997). However, no careful calculations have yet been done, taking the dynamics of the process into account, notably the dilution of SN ejecta and the adiabatic losses. Yet we show below that they have a crucial influence on the total amount of Be produced, and that a time-dependent treatment is required. Indeed, the evolution of a SNR is essentially a dynamical problem in which the acceleration rate as well as the chemical composition inside the remnant are functions of time. The results of the full calculation of both processes and the discussion of their implications for the chemical evolution of the Galaxy will be found in an associated paper (Parizot & Drury 1999). Here we present simple analytical calculations which provide an accurate understanding of the dynamics of light element production in SNRs and elucidates the role and influence of the different parameters, notably the ambient density.

Although Li and B are also produced in the processes under study, we shall choose here Be as our ‘typical’ light element, because nuclear spallation of CNO is thought to be its only production mechanism, while Li is also (and actually mainly) produced through \(\alpha + \alpha\) reactions, \(^7\)Li may be produced partly in AGB stars (Abia et al. 1993), and \(^{11}\)B neutrino spallation may be important as seems to be required by chemical evolution analysis (Vangioni-Flam et al. 1996). In order to compare our results with the observations, we simply note that, as emphasized in Ramaty et al. (1997), the data relating to the Galactic Be evolution as a function of [Fe/H] indicate that \(\sim 1.610^{-6}\) nuclei of Be must be produced in the early Galaxy for each Fe nucleus. Therefore, if Be production is indeed induced, directly or indirectly, by SNe explosions, and since the average SN yield in Fe is thought to be \(\sim 0.11 M_\odot\), each supernova must lead to an average production of \(\sim 3.810^{38}\) nuclei (or \(\sim 2.810^{-8} M_\odot\)) of Be, with an uncertainty of about a factor of 2 (Ramaty et al. 1997). We adopt this value as the ‘standard needed number’ of Be per supernova explosion. To state this again in a different way, for an average SN yield in CNO of, say, \(\sim 1 M_\odot\), the required spallation rate per CNO atom is \(\sim 310^{-8}\).

2. Particle acceleration in SNRs

It is generally believed that cosmic ray production in SNRs occurs through the process of diffusive shock acceleration operating at the strong shock waves generated by the interaction between the ejecta from the supernova explosion and the surrounding medium. Significant effort has been put into developing dynamical models of SNR evolution which incorporate, at varying levels of detail, this basic acceleration and injection process (one of the major advantages of shock acceleration is that it does not require a separate injection process). Qualitatively the main features can be crudely summarised as follows.

In a core collapse SN the collapse releases roughly the gravitational binding energy of a neutron star, some \(10^{53}\) erg, but most of this is radiated away in neutrinos. About \(E_{\text{SN}} = 10^{51}\) erg is transferred, by processes which are still somewhat obscure, to the outer layers of the progenitor star which are then ejected at velocities of a few percent of the speed of light. Initially the explosion energy is almost entirely in the form of kinetic energy of these fast-moving ejecta. As the ejecta interact with the surrounding circumstellar and interstellar material they drive a strong shock ahead into the surrounding medium. The region of very hot high pressure shocked material behind this forward shock also drives a weaker shock backwards into the ejecta giving rise to a characteristic forward reverse shock pair separated by a rather unstable contact discontinuity.
This initial phase of the remnant evolution lasts until the amount of ambient matter swept up by the remnant is roughly equal to the original ejecta mass. At this so-called sweep-up time, $t_{SW}$, the energy flux through the shocks is at its highest, the expansion of the remnant begins to slow down, and a significant part of the explosion energy has been converted from kinetic energy associated with the bulk expansion to thermal (and non-thermal) energy associated with microscopic degrees of freedom of the system. The remnant now enters the second, and main, phase of its evolution in which there is rough equipartition between the microscopic and macroscopic energy densities. The evolution in this phase is approximately self-similar and resembles the exact solution obtained by Sedov for a strong point explosion in a cold gas.

It is important to realise that the approximate equality of the energy associated with the macroscopic and microscopic degrees of freedom in the Sedov-like phase is not a static equilibrium but is generated dynamically by two competing processes. As long as the remnant is compact the energy density, and thus pressure, of the microscopic degrees of freedom is very much greater than that of the external medium. This strong pressure gradient drives an expansion of the remnant which adiabatically reduces the microscopic degrees of freedom of the medium inside the remnant and converts the energy back into bulk kinetic energy of expansion. At the same time the strong shock which marks the boundary of the remnant converts this macroscopic kinetic energy of expansion back into microscopic internal form. Thus there is a continuous recycling of the original explosion energy between the micro and macro scales. This continues until either the external pressure is no longer negligible compared to the internal, or the time-scales become so long that radiative cooling becomes important. The time scales for the conversion of kinetic energy to internal energy and vice versa are roughly equal and of order the dynamical time scale of the remnant which is of order the age of the remnant, hence the approximately self-similar evolution.

In terms of particle acceleration the theory assumes that strong collisionless shocks in a tenuous plasma automatically and inevitably generate an approximately power law distribution of accelerated particles which connects smoothly to the shock-heated particle distribution at ‘thermal’ energies and extends up to a maximum energy constrained by the shock size, speed, age and magnetic field. The acceleration mechanism is a variant of Fermi acceleration based on scattering from magnetic field structures on both sides of the shock. A key point is that these scattering structures are not those responsible for general scattering on the ISM, but strongly amplified local structures generated in a non-linear bootstrap process by the accelerated particles themselves. As long as the shock is strong it will be associated with strong magnetic turbulence which drives the effective local diffusion coefficient down to values close to the Bohm value. As pointed out by Achterberg et al. (1994) the extreme sharpness of the radio rims of some shell type SNRs can be interpreted as observational evidence for this type of effect. The source of free energy for the wave excitation is of course the strong gradient in the energetic particle distribution at the edge. Thus in the interior of the remnant, where the gradients are absent or much weaker, we do not expect such low values of the diffusion coefficient.

The net effect is that the edge of the remnant, as far as the accelerated particles are concerned, is both a self-generated diffusion barrier and a source of freshly accelerated particles. Except at the very highest energies the particles produced at the shock are convected with the post-shock flow into the interior of the remnant and effectively trapped there until the shock weakens to the point where the self-generated wave field around the shock can no longer be sustained. At this point the diffusion barrier collapses and the trapped particle population is free to diffuse out into the general ISM.

In terms of bulk energetics, the total energy of the accelerated particle population is low during the first ballistic phase of the expansion (because little of the explosion energy has been processed through the shocks) but rises rapidly as $t \approx t_{SW}$. During the sedov-like phase the total energy in accelerated particles is roughly constant at a significant fraction of the explosion energy (0.1 to 0.5 typically). However, this is because of the dynamic recycling described above. Any individual particle is subject to adiabatic losses on the dynamical time-scale of the remnant, while the energy lost this way goes into driving the shock and thus generating new particles, distributed over the whole energy spectrum.

3. Spallation reactions within SNRs

3.1. Qualitative overview

We now turn to the production of Li, Be and B (LiBeB) by spallation reactions within a SN. As emphasized above, there are two obvious mechanisms. One is the irradiation of the CNO ejecta by accelerated protons and alphas. It is clear that the fresh CNO nuclei produced by the SN will, for the lifetime of the SNR, be exposed to a flux of energetic particles (EPs) very much higher than the average interstellar flux, and this must lead to some spallation production of light elements. This process starts at about $t_{SW}$ with a very intense radiation field and continues with an intensity decreasing roughly as $R^{-3} \propto t^{-6/5}$ (where $R$ is the radius of the SNR) until the remnant dies.

The second process is that some of the CNO nuclei from the ejecta are accelerated, either by the reverse shock in its brief powerful phase at $t \approx t_{SW}$ or by some of this material managing to get ahead of the forward shock. This later possibility is not impossible, but seems unlikely to be as important as acceleration by the reverse shock. Calculations of the Raleigh-Taylor instability of the contact discontinuity do suggest that some fast-moving blobs
of ejecta can punch through the forward shock at about $t_{SW}$ (Jun & Norman 1996), and in addition Ramaty and coworkers have suggested that fast moving dust grains could condense in the ejecta at $t < t_{SW}$ and then penetrate through into the region ahead of the main shock. In all these pictures acceleration of CNO nuclei takes place only at about $t_{SW}$ and the energy deposited in these accelerated particles is certainly less than the explosion energy $E_{SN}$, although it might optimistically reach some significant fraction of that value (say $\leq 10\%$). Crucially, the accelerated CNO nuclei are then confined to the interior of the SNR and will thus be adiabatically cooled on a rather rapid time-scale, initially of order $t_{SW}$.

3.2. Evaluation of the first process (forward shock)

From the above arguments, it is clear that SNe do induce some Be production. Now the question is : how much? Let us first consider the irradiation of the ejecta by particles (H and He nuclei) accelerated at the forward shock during the Sedov-like phase – process 1. We have already indicated that detailed studies of acceleration in SNRs show that the fraction of the explosion energy given to the EPs is roughly constant during the Sedov-like phase and of order 0.1 to 0.5 or so. Let $\theta_1$ be that fraction. Since the EPs are distributed more or less uniformly throughout the interior of the remnant, the energy density can be estimated as

$$\mathcal{E}_{CR} \approx \frac{3\theta_1 E_{SN}}{4\pi R^3}$$

where $R$ is the radius of the remnant and $E_{SN}$ is the explosion energy.

To derive a spallation rate from this we need to assume some form for the spectrum of the accelerated particles. Shock acceleration suggests that the distribution function should be close to the test-particle form $f(p) \propto p^{-4}$ and extend from an injection momentum close to ‘thermal’ values to a cut-off momentum at about $p_{max} = 10^5$ GeV/c. The spallation rate per target CNO atom to produce a Be atom is then obtained by integrating the cross sections

$$\nu_{spall} = \int_{p_{th}}^{p_{max}} \sigma vf(p) dp,$$

with the normalisation $\int E(p) f(p) 4\pi p^2 dp = \mathcal{E}_{CR}$. Looking at graphs of the spallation cross-sections for Be (as given, e.g., in Ramaty et al. 1997), it is clear that these cross-sections can be well approximated as zero below a threshold at about 30–40 MeV/n and a constant value $\sigma_0 \approx 5 \times 10^{-27}$ cm$^2$ above it. One then obtains roughly :

$$\nu_{spall} \approx \frac{\sigma_0}{mc} \mathcal{E}_{CR} \frac{1 - \ln(p_{th}/mc)}{1 + \ln(p_{max}/mc)} \approx 0.2\sigma_0 c \mathcal{E}_{CR} \frac{mc}{m^2},$$

where $p_{th} \approx mc/5$ is the momentum corresponding to the spallation threshold and $m$ refers to the proton mass. Fortunately, for this form of the spectrum the upper cut-off and the spallation threshold only enter logarithmically. A softer spectrum would lead to higher spallation yields and a stronger dependence on the spallation threshold.

Using Eq. (1) and the adiabatic expansion law for the forward shock radius, $R = R_{SW}(t/t_{SW})^{2/5}$, we can now estimate the total fraction of the CNO nuclei which will be converted to Be during the Sedov-like phase as :

$$\phi_1 = \int_{t_{SW}}^{t_{end}} 0.2\sigma_0 c \frac{\theta_1 E_{SN}}{mc^2} \frac{3}{4\pi R^3} dt$$

$$= 0.2\sigma_0 c \frac{\theta_1 E_{SN}}{mc^2} \frac{3}{4\pi R_{SW}^3} \int_{t_{SW}}^{t_{end}} \left( \frac{t}{t_{SW}} \right)^{-6/5} dt,$$

or

$$\phi_1 = \sigma_0 c \frac{\theta_1 E_{SN}}{mc^2} \frac{\rho_0}{M_{ej}} t_{SW} \left[ 1 - \left( \frac{t_{SW}}{t_{end}} \right)^{1/5} \right]$$

where as usual $\rho_0$ denotes the density of the ambient medium into which the SNR is expanding and $M_{ej}$ is the total mass of the SNR ejecta. We now recall that the sweep-up time is given in terms of the SN parameters and the ambient number density, $n_0 \approx \rho_0/m$, as

$$t_{SW} = \frac{n_0^{-1/3}}{v_\infty} \left( \frac{3 M_{ej}}{4\pi m} \right)^{1/3},$$

where $v_\infty \approx (2E_{SN}/M_{ej})^{1/2}$ is the velocity of the ejecta, or numerically :

$$t_{SW} = (1.4 \times 10^4 \text{yr}) \left( \frac{M_{ej}}{10M_\odot} \right)^{2/3} \left( \frac{E_{SN}}{10^{51}\text{erg}} \right)^{-1/3} \left( \frac{n_0}{1\text{cm}^{-3}} \right)^{-1/3}.$$

Replacing in Eq. (6) and using canonical values of $E_{SN} = 10^{51}$ erg and $M_{ej} = 10M_\odot$, we finally get :

$$\phi_1 \approx 4 \times 10^{-10} \theta_1 \left( \frac{n_0}{1\text{cm}^{-3}} \right)^{2/3} \left[ 1 - \left( \frac{t_{SW}}{t_{end}} \right)^{1/5} \right].$$

Clearly this falls short of the value of order $10^{-8}$ required to explain the observations, even for values of $\theta_1$ as high as 0.5. It might seem from Eq. (7) that very high ambient densities could help to make the spallation yields closer to the needed value. This is however not the case. First, the above estimate does not take energy losses into account, while both ionisation and adiabatic losses act to lower the genuine production rates. Second, and more significantly, the ratio $t_{SW}/t_{end}$ (and a fortiori its fifth root) becomes very close to 1 in dense environments, lowering $\phi_1$ quite notably (see Fig. 3). In fact, it turns out that there is no Sedov-like phase at all in media with densities of order $10^4$ cm$^{-3}$, the physical reason being that the radiative losses then act on a much shorter time-scale, eventually shorter than the sweep-up time.
3.3. Evaluation of the second process (reverse shock)

Let us now turn to the second process, namely the spallation of energetic CNO nuclei accelerated at the reverse shock from the SN ejecta and interacting within the SNR with swept-up ambient material. We have argued above that this reverse shock acceleration is only plausible at times around $t_{SW}$ and certainly the amount of energy transferred to CNO nuclei cannot be more than a fraction of $E_{SN}$. Let $\theta_d$ be the fraction of the explosion energy that goes into accelerating the ejecta at or around $t_{SW}$, and $\theta_{CNO}$ the fraction of that energy that is indeed transferred to CNO nuclei. These particles are then confined to the interior of the remnant where they undergo spallation reactions as well as adiabatic losses. Let us again assume that the spectrum is of the form $f(p) \propto p^{-4}$. Then the production rate of Be atoms per unit volume is approximately

$$0.2 \frac{n_0 \sigma_0 E_{CR}}{mc^2 \pi^3}$$

(10)

where $E_{CR}$ now refers to the accelerated CNO nuclei, the factor 14 comes from the mean number of nucleons per CNO nucleus and the factor 0.2, as before, from the $f(p) \propto p^{-4}$ spectral shape (assuming the same upper cutoff position, but this only enters logarithmically). Integrating over the remnant volume, we obtain the spallation rate at $t_{SW}$:

$$\frac{dN_{Be}}{dt} \approx 0.2 \sigma_0 \frac{\theta_{CNO} \theta_d E_{SN}}{14mc^2}$$

(11)

Now the adiabatic losses need to be evaluated rather carefully. It is generally argued that they act so that the momentum of the particles scales as the inverse of the linear dimensions of the volume occupied. Accordingly, in the expanding spherical SNR the EPs should lose momentum at a rate $\dot{p}/p = -\dot{R}/R$, reminiscent, incidentally, of the way photons behave in the expanding universe. In our case, however, the situation is complicated by the fact that the EPs do not push directly against the ‘walls’ limiting the volume of confinement, which move at the expansion velocity, $V = \dot{R}$, but are reflected off the diffusion barrier consisting of magnetic waves and turbulence at rest with respect to the downstream flow, and thus expanding at velocity $\frac{3}{4} \dot{R}$.

To see how this influences the actual adiabatic loss rate, it is safer to go back to basic physical laws. Adiabatic losses must arise because the EPs are more or less isotropised within the SNR and therefore participate to the pressure. Now this pressure, $P$, works positively while the remnant expands, implying an energy loss rate equal to the power contributed, given by:

$$\frac{dU}{dt} = - \int_S F \cdot v = - \int_S P dS \times \frac{3}{4} \dot{R} = -3\pi R^2 \dot{R} P$$

(12)

where $U = \frac{4}{3} \pi R^3 c$ is the total kinetic energy of the particles. Considering that $P = \frac{2}{3} \dot{c}$ in the non-relativistic limit (NR) and $P = \frac{4}{3} \dot{c}$ in the ultra-relativistic limit (UR), Eq. (12) can be re-written as:

$$\frac{dE}{dt} = - \frac{9}{4} \frac{\dot{R}}{R} = -\frac{3 \dot{R}}{2 R}$$

NR

$$= \frac{3 \dot{R}}{4 R}$$

UR

(13)

Finally, dividing both sides by the space density of the EPs and noting that $E = p^2/2m$ in the NR limit, and $E = pc$ in the UR limit, we obtain the momentum loss rate for individual particles, valid in any velocity range:

$$\frac{\dot{p}}{p} = - \frac{3 \dot{R}}{4 R}$$

(14)

From this one deduces that at the time when the remnant has expanded to radius $R$, only those particles whose initial momenta at $t_{SW}$ were more than $(R/R_{SW})^{3/4} p_{th}$ are still above the spallation threshold. For a $p^{-4}$ distribution function the integral number spectrum decreases as $p^{-4}$ and thus the number of accelerated nuclei still capable of spallation reactions decreases as $R^{-3/4} \propto t^{-3/10}$. For a softer accelerated spectrum the effect would be even stronger because there are proportionally fewer particles at high initial momenta.

This being established, we can integrate Eq. (11) over time, to obtain the total production of Be atoms:

$$N_{Be} = \frac{0.2 n_0 \sigma_0 c \theta_{CNO} \theta_d E_{SN}}{14mc^2} \int_{t_{SW}}^{t_{end}} \left( \frac{t}{t_{SW}} \right)^{-3/10} dt$$

(15)

that is:

$$N_{Be} = \frac{2}{3} n_0 \sigma_0 c \theta_{CNO} \theta_d E_{SN} t_{SW} \left[ \left( \frac{t_{end}}{t_{SW}} \right)^{7/10} - 1 \right]$$

(16)

Dividing by the total number of CNO nuclei in the ejecta, $N_{ej,CNO} = \phi_{CNO} N_{ej,tot}$, where:

$$\phi_{CNO} = \frac{\theta_{CNO} E_{SN}}{14mc^2} \rho_0 \frac{2}{7} \left( \frac{t_{end}}{t_{SW}} \right)^{7/10} \left[ 1 - \left( \frac{t_{SW}}{t_{end}} \right)^{7/10} \right]$$

(17)

Note that we assumed that the mass fraction of CNO in the ejecta is the same as the energy fraction of CNO in the EPs (which was the original meaning of $\theta_{CNO}$). Considering that all nuclear species have the same spectrum in MeV/n, and thus a total energy proportional to their mass number, this simply means that the acceleration process is not chemically selective, in the sense that the composition of the EPs is just the same as that of the material passing through the shock.
Numerically, again with \( E_{\text{SN}} = 10^{51} \) erg and \( M_{\text{ej}} = 10 M_\odot \), we finally obtain:

\[
\phi_2 \approx 1 \times 10^{-10} \theta_2 \left( \frac{n_0}{1 \text{ cm}^{-3}} \right)^{2/3} \times \left( \frac{t_{\text{end}}}{t_{\text{SW}}} \right)^{7/10} \left[ 1 - \left( \frac{t_{\text{SW}}}{t_{\text{end}}} \right)^{7/10} \right].
\]

### 3.4. Relative contribution of the two processes

It is worth emphasizing the similarity between expressions (30) and (17) that we obtained for the spallation rates per CNO nuclei by the two processes considered here. This formal analogy allows us to write down their relative contributions straightforwardly:

\[
\frac{\phi_2}{\phi_1} = \frac{\theta_2}{\theta_1} \times \frac{2}{7} \left( \frac{t_{\text{end}}}{t_{\text{SW}}} \right)^{7/10} \left[ 1 - \left( \frac{t_{\text{SW}}}{t_{\text{end}}} \right)^{7/10} \right] \left[ 1 - \left( \frac{t_{\text{SW}}}{t_{\text{end}}} \right)^{1/5} \right].
\]

As is often the case, this similarity is not fortuitous and has a physical meaning. The two processes may indeed be regarded as ‘dual’ processes, the first consisting of the irradiation of the SN ejecta by the ambient medium, and the second of the ambient medium by the SN ejecta. The ‘symmetry’ is only broken by the dynamical aspect of the processes. First, of course, the energy imparted to the EPs ‘symmetry’ is only broken by the dynamical aspect of the EPs in both cases needs not be the same, for it depends on the acceleration efficiency as well as the total energy of the shock involved (forward or reverse). This is expressed by the expected ratio \( \theta_2/\theta_1 \). And secondly, in the first process one has to fight against the dilution of the ejecta – integration of \((t/t_{\text{SW}})^{-6/5}\), see Eq. (18) – while in the second process one fights against the adiabatic losses – integration of \((t/t_{\text{SW}})^{-3/10}\), see Eq. (15). This is expressed by the last factor in Eq. (19).

Clearly the latter decrease of the production rates is the least dramatic, and the reverse shock process must dominate the LiBeB production in supernova remnants. However, this conclusion still depends on the genuine efficiency of reverse shock acceleration, and once the relative acceleration efficiency \( \theta_2/\theta_1 \) is given, the weight of the first process relative to the second still depends on the total duration of the Sedov-like phase, appearing numerically in Eq. (18) through the ratio \( t_{\text{end}}/t_{\text{SW}} \), which in turn depends on the ambient density, \( n_0 \). The expression of \( t_{\text{SW}} \) as a function of the parameters has been given in Eq. (15), so we are left with the evaluation of the time, \( t_{\text{end}} \), when the magnetic turbulence collapses and the EPs leave the SNR, putting an end to Be production. We argued above that \( t_{\text{end}} \) should correspond to the end the Sedov-like phase, when the shock induced by the SN explosion becomes radiative, that is when the cooling time of the post-shock gas becomes of the same order as the dynamical time.

In principle, the cooling rate can be derived from the so-called cooling function, \( \Lambda(T) \)(erg cm\(^{-3}\)s\(^{-1}\)), which depends on the physical properties of the post-shock material, notably on its temperature, \( T \), and metallicity, \( Z \):

\[
\tau_{\text{cool}} \approx \frac{3k_B T}{n \Lambda(T)},
\]

where \( n \) is the post-shock density, equal to \( 4n_0 \) if the compression ratio is that of an ideal strong shock (nonlinear effects probably act to increase the compression ratio to values larger than 4). As for the dynamical time, we simply write

\[
\tau_{\text{dyn}} \approx \frac{R}{R} \approx \frac{5}{2} t.
\]

To obtain \( t_{\text{end}} \), we then need to solve the following equation in the variable \( t \), obtained by equating \( \tau_{\text{cool}} \) and \( \tau_{\text{dyn}} \) given above:

\[
t \approx \frac{3k_B T}{20n_0 \Lambda(T)},
\]

where it should be clear that the right hand side also depends on time through the temperature, \( T \), and thus indirectly through the cooling function too. In the non-radiative SNR expansion phase, the function \( T(t) \) is obtained directly from the hydrodynamical jump conditions at the shock discontinuity:

\[
T \approx \frac{3m}{8k_B} V^2,
\]

or numerically:

\[
T \approx (2 \times 10^5 \text{ K}) \left( \frac{E_{\text{SN}}}{10^{51} \text{ erg}} \right)^{\frac{1}{3}} \left( \frac{n_0}{1 \text{ cm}^{-3}} \right)^{-\frac{2}{3}} \left( \frac{t}{10^7 \text{ yr}} \right)^{-\frac{1}{3}}.
\]
it happens to depend significantly on metallicity, with differences up to two orders of magnitude for metallicities from \( Z = 0 \) to \( Z = 2Z_{\odot} \) (Böhringer & Hensler 1989). Because we focus on Be production in the early Galaxy, we adopt the cooling function corresponding to zero metallicity, represented in Fig. 2 (adapted from Böhringer & Hensler 1989), which holds for values of \( Z \) up to \( \sim 10^{-2}Z_{\odot} \).

For high enough ambient densities, the shock will become radiative early in the SNR evolution, when the temperature is still very high, say above \( T > 2 \times 10^6 \) K. In this case, the cooling function is dominated by Bremsstrahlung emission and can be written analytically as:

\[
\Lambda_{\text{Br}}(T) \approx (2.4 \times 10^{-23} \text{erg cm}^3\text{s}^{-1})(\frac{T}{10^6 \text{K}})^{1/2}. \tag{25}
\]

Substituting from (24) and (23) in Eq. (22) and solving for \( t \), we find:

\[
t_{\text{end}} = (1.1 \times 10^5 \text{yr}) \left( \frac{E_{\text{SN}}}{10^{51} \text{erg}} \right)^{1/8} \left( \frac{n_0}{1 \text{cm}^{-3}} \right)^{-3/4}. \tag{26}
\]

To check the consistency of our assumption \( \Lambda \approx \Lambda_{\text{Br}} \) (i.e. \( T \gtrsim 2 \times 10^6 \) K), let us now report Eq. (26) in (23) and write down the temperature \( T_{\text{End}} \) at the end of the Sedov-like phase:

\[
T_{\text{End}} \approx (2 \times 10^5 \text{K}) \left( \frac{E_{\text{SN}}}{10^{51} \text{erg}} \right)^{1/4} \left( \frac{n_0}{1 \text{cm}^{-3}} \right)^{1/2}, \tag{27}
\]

which means that the above analytical treatment is valid only for ambient densities greater than about 100 cm\(^{-3}\). For lower densities, we must solve Eq. (23) graphically. First, we invert Eq. (24) to express \( t \) as a function of temperature, then we plot the function \( f(T) = 3k_B T / (20 n_0 t) \) on the same graph as \( \Lambda \) (see Fig. 2 for an example), find the value of \( T \) at intersection, and finally convert this value into the sought time \( t_{\text{end}} \) making use again of Eq. (24). The results, showing \( t_{\text{end}} \) as a function of the ambient density, are shown in Fig. 3.

We now have all the ingredients to plot the efficiency ratio of the two processes calculated above. Figure 3 shows the ratio \( \phi_2/\phi_1 \) given in Eq. (13) as a function of the ambient density, assuming that \( \theta_1 = \theta_2 \). Two different values of the ejected mass have been used, corresponding to different progenitor masses (\( \sim 10 – 40 M_{\odot} \)). It can be seen that low densities are more favourable to the reverse shock acceleration process. This is due to \( t_{\text{end}}/t_{\text{SW}} \) being larger, implying a larger dilution of the ejecta (process 1 less efficient) and smaller adiabatic losses, which indeed decrease as \( t^{-1} \) (process 2 more efficient). The part of the plot corresponding to \( \phi_2/\phi_1 \leq 1 \) is not physical, because it requires \( t_{\text{end}} \leq t_{\text{SW}} \), which simply means that the Sedov-like phase no longer exists and the whole calculation becomes groundless. Note however that in Fig. 3 the energy imparted to the EPs has been assumed equal for both processes, which is most certainly not the case. Actually, if \( \theta_2/\theta_1 = 0.1 \) (e.g. \( \theta_1 = 10\% \) and \( \theta_1 = 1\% \)), then process 1 is found to dominate Be production during the Sedov-like phase, regardless of the ambient density.

4. Spallation reactions after the Sedov-like phase
At the end of the Sedov-like phase, the EPs are no longer confined and leave the SNR to diffuse across the Galaxy. At the stage of chemical evolution we are considering here, there are no or few metals in the interstellar medium (ISM), so that energetic protons and \( \alpha \) particles accelerated at the forward shock will not produce any significant amount of Be after \( t_{\text{end}} \) (although Li production will still be going on through \( \alpha + \alpha \) reactions). In the case of the second process, however, the EPs contain CNO nuclei which just cannot avoid being spalled while interacting with the

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**Fig. 2.** Comparison of \( t_{\text{end}} \) and \( t_{\text{SW}} \) as a function of the ambient density, \( n_0 \), for two different values of the mass ejected by the supernova (10 and 30 \( M_{\odot} \)). The dashed line shows the asymptotic analytic estimate of Eq. (26).

**Fig. 3.** Comparison of \( t_{\text{end}} \) and \( t_{\text{SW}} \) as a function of the ambient density, \( n_0 \), for two different values of the mass ejected by the supernova (10 and 30 \( M_{\odot} \)). The dashed line shows the asymptotic analytic estimate of Eq. (26).
ambient H and He nuclei at rest in the Galaxy. This may be regarded as a third process for Be production, which lasts until either the EPs are slowed down by Coulombian interactions to subnuclear energies (i.e. below the spallation thresholds) or they simply diffuse out of the Galaxy. Since the confinement time of cosmic rays in the early Galaxy is virtually unknown, we shall assume here that the Galaxy acts as a thick target for the EPs leaving the SNR, an assumption which actually provides us with an upper limit on the spallation yields.

Unlike the first two processes evaluated above, this third process is essentially independent of dynamics. Thus, time-dependent calculations are no longer needed and, from this stage on, the calculations made by Ramaty et al. (1997) or any steady-state calculation is perfectly valid. In particular, the ambient density has no influence on light element production, since a greater number of reactions per second, as would result from a greater density, implies an equal increase of both the spallation rates and the energy loss rate. Once integrated over time, both effects cancel out exactly, and in fact, given the energy spectrum of the EPs, the efficiency of Be production (and Li, and B), expressed as the number of nuclei produced per erg injected in the form of EPs, depends only on their chemical composition.

Results are shown in Fig. 4 for different values of the source abundances ratios, H/He and (H + He)/(C + O), allowing one to derive the spallation efficiency for any composition. Two-steps processes (such as $^{12}$C $\rightarrow$ $^{10}$B followed by $^{10}$B $\rightarrow$ H $\rightarrow$ Be) have been taken into account. Test runs show good agreement with the results of Ramaty et al. (1997).

As can be seen on Fig. 4, pure Carbon and Oxygen have a production efficiency of about 0.22 nuclei/erg, while this efficiency decreases by at least a factor of 10 for compositions with hundred times more H and He than metals (or about ten times more by mass). According to models of explosions for SN with low metallicity progenitors, the average (H+He)/(C+O) ratio among the EPs should indeed be expected to be $\gtrsim 200$, unless selective acceleration occurs to enhance the abundance of the metals. As a consequence, efficiencies greater than $\sim 10^{-2}$ nuclei/erg should not be expected, so that a production of $\sim 4\times 10^{48}$ atoms of Be requires an energy of $\sim 4\times 10^{50}$ erg to be imparted to the EPs. This seems very unlikely considering that the total energy available in the reverse shock (the source of the EPs) should be of order one tenth of the SN explosion energy, not to mention the acceleration efficiency. Moreover, a significant fraction of the energy originally imparted to the EPs has been lost during the Sedov-like phase of the SNR evolution through adiabatic losses.

To evaluate the ‘surviving’ fraction of energy, it suffices to go back to Eq. (4), which indicates that when the radius of the shock is multiplied by a factor $\eta$, the momentum $p$ of all the particles is multiplied by a factor $\eta^{-3/4}$. It is worthwhile noting that, because of their specific momentum dependence, adiabatic losses do not modify the shape of the EP energy spectrum. In our case, $f(p) \propto p^{-4}$, so that when all momenta $p$ are divided by a factor $\zeta$, the distribution function $f(p)$ is divided by the same factor $\zeta$. To see that, the easiest way is to work out the number of particles between momenta $p$ and $p + dp$ after the momentum scaling. This number writes $dN' = f'(p)4\pi p^2 dp$, where $f'(p)$ is the new distribution function. Now $dN'$ must be equal to the number of particles that had momentum between $\zeta p$ and $\zeta(p + dp)$, which is, by definition, $dN = f(\zeta p)4\pi (\zeta p)^2 \zeta dp$. Equating $dN$ and $dN'$ yields the result $f'(p) = \zeta^{-1} f(p)$.

Putting all pieces together, we find that when the shock radius $R$ is multiplied by a factor $\eta$, the distribution function and, thus, the total energy of the EPs are multiplied by $\eta^{-3/4}$. Now considering that $R$ increases as
$t^{2/5}$ during the Sedov-like phase, we find that the total energy of the EPs decreases as $t^{-3/10}$. Note that this is nothing but an other way to work out the decrease of the spallation rates for our second process during the Sedov-like phase (cf. Sect. 3.3). Finally, we find that a fraction $(t_{\text{end}}/t_{\text{SW}})^{-3/10}$ of the initial energy imparted to the EPs is still available for spallation at the end of the Sedov-like phase. This factor is plotted on Fig. 3 as a function of the ambient density. It can be seen that for $n_0 = 1 \text{ cm}^{-3}$, the energy available to power our third process of light element production has been reduced by adiabatic losses to not more than one third of its initial value, and less than one half for densities up to $100 \text{ cm}^{-3}$. Clearly, high densities are favoured (energetically) because they tend to shorten the Sedov-like phase, and therefore merely avoid the adiabatic losses.

5. Discussion

Since light element production in the interstellar medium obviously requires a lot of energy in the form of supernova-decelerated particles (i.e. with energies above the nuclear thresholds) as well as metals (especially C and O), it is quite natural to consider SNe as possible sources of the LiBeB observed in halo stars. We have analysed in detail the spallation nucleosynthesis induced by a SN explosion on the basis of known physics and theoretical results relating to particle shock acceleration. Two major processes can be identified, depending on whether the ISM or the ejecta are accelerated, respectively at the forward and reverse shocks. In the first case, the EPs consist mostly of protons and alpha particles and must therefore interact with C and O nuclei, which are much more numerous within the SNR than in the surrounding medium (especially at early stages of Galactic evolution). The process will thus last as long as the EPs stay confined in the SNR, i.e. approximately during the Sedov-like phase, but not more. In the second case, freshly synthesized CNO nuclei are accelerated, and Be production occurs through interaction with ambient H and He nuclei. The process is then divided into two, one stretching over the Sedov-like phase, with the particles suffering adiabatic losses, and the other one occurring outside the remnant, with only Coulombian losses playing a role.

We have calculated the total Be production in these three processes, taking the dynamics of the SNR evolution into account (dilution of the ejecta by metal-poor material and adiabatic losses). The results are shown in Fig. 6 for processes 1 and 2 (from Eqs. (9) and (18)). We find that with canonical values of $\theta_1 = 0.1$, $\theta_2 = 0.01$, $M_{\text{ej}} = 10 M_\odot$ and a mean ambient density $n_0 = 10 \text{ cm}^{-3}$, the fraction of freshly synthesized CNO nuclei spalled into Be in these processes is $\phi_1 \sim 8 \times 10^{-11}$ and $\phi_2 \sim 3 \times 10^{-11}$, respectively, which is very much less than the value ‘required’ by the observations, discussed in the introduction ($\phi_{\text{obs}} \sim 3 \times 10^{-8}$). Even allowing for unreasonably high values of the acceleration efficiency, $\theta_1 \sim \theta_2 \lesssim 1$, the total Be production by processes 1 and 2 would still be more than one order of magnitude below the observed value.

As suggested by Fig. 3 and our analytical study, higher densities improve the situation. However, even with $n_0 = 10^3 \text{ cm}^{-3}$ and acceleration efficiencies equal to 1, the Be yield is still unsufficient. Moreover, it should be noted that our calculations did not consider Coulombian energy losses (because they are negligible as compared to adiabatic losses for usual densities), which become important as the density increases and therefore make the Be yield smaller. Finally, since we are trying to account for the mean abundance of Be in halo stars, as compared to Fe, we have to evaluate the Be production for an ambient density corresponding to the mean density encountered around explosion sites in the early Galaxy, which is very unlikely to be as high as $10^3 \text{ cm}^{-3}$. It could even be argued that although the gas density might have been higher in the past than it is now (hence our ‘canonical value’ $n_0 = 10 \text{ cm}^{-3}$), the actual mean density about SN explosion sites could be lower than $1 \text{ cm}^{-3}$, because most SNe may explode within superbubble interiors, where the density is much less than in the mean ISM.

Thus, our conclusion is that processes 1 and 2 both fail in accounting for the Be observed in metal-poor stars in the halo of our Galaxy. Concerning the third process, adopting canonical values for the parameters again leads to insufficient Be production, as noted in the previous section. While higher densities improve the situation by avoiding the adiabatic energy losses, one should nevertheless expect at least half of the EP energy to be lost in this way, for any reasonable density. This means that even if 10% of the explosion energy is imparted to EPs accelerated at the reverse shock, which is certainly a generous
upper limit, the required number of $\sim 4 \times 10^{48}$ nuclei of Be per SN implies a spallation efficiency of $\sim 0.1$ nucleus/erg. Now Fig. 3 shows that this requires an EP composition in which at least one particle out of ten is a CNO nucleus. In other words, the ejected mass of CNO must be of the order of that of H and He together. None of the SN explosion models published so far can reproduce such a requirement, and so there is clearly a problem with Be production in the early Galaxy.

The results presented here are in fact interesting in many regards. First, they show that it is definitely very difficult to account for the amount of Be found in halo stars. Consequently, we feel that the main problem to be addressed in this field of research is probably not the chemical evolution of Be (and Li and B) in the Galaxy, as given by the ensemble of the data points in the abundance vs metallicity diagrams (e.g. whether Be is proportional to Fe or to its square) but, to begin with, the position of any of these points. Are we able to describe in some detail one process which could explain the amount of Be (relative to Fe) present in any of the stars in which it is observed? The answer, we are afraid, seems to be no at this stage. It is however instructive to ask why the processes investigated here have failed. Concerning process 1 (acceleration of ISM, interaction with fresh CNO within the SNR), the main reason is that the CNO rich ejecta are ‘too much diluted’ by the swept-up material as the SNR expands, so that the spallation efficiency is too low (or the available energy is too small). However, it seems rather hard to think of any region in the Galaxy where the concentration in CNO is higher than inside a SNR during the Sedov-like phase (especially in the first stages of chemical evolution)! So the conclusion that process 1 cannot work, even with a 100% acceleration efficiency, seems to rule out any other process based on the acceleration of the ISM, initially devoided of metals.

The other solution is then of course to accelerate CNO nuclei themselves, which provides the maximum possible spallation efficiency, independently of the ambient metallicity. Every energetic CNO will lead to the production of as much Be as possible given the spallation cross sections and the energy loss rates. The latter cannot physically be smaller than the Coulombian loss rate in a neutral medium, and this leads to the efficiency plotted in Fig. 3. Unfortunately, a significant amount of the CNO rich ejecta of an isolated SN can only be accelerated at the reverse shock at a time around the sweep-up time, $t_{SW}$. This means that i) the total amount of energy available is smaller than the explosion energy (probably of order 10%, i.e. $\sim 10^{50}$ erg), and ii) the accelerated nuclei will suffer adiabatic losses during the Sedov-like phase, reducing their energy by a factor of 2 or 3. As shown above, this makes process 2-3 incapable of producing enough Be, as long as the EPs have a composition reflecting that of the SN ejecta.

This suggest that a solution to the problem could be that the reverse shock accelerates preferentially CNO nuclei rather than H and He. For example, recent calculations have shown that such a selective acceleration arises naturally if the metals are mostly condensed in grains (Ellison et al. 1997). The proposition by Ramaty et al. (1997) that grains condense in the ejecta before being accelerated could then help to increase the abundance of CNO in the EPs. However, we have to keep in mind that any selective process called upon must be very efficient indeed, since as we indicated above, the data require that the EP composition be as rich as one CNO nuclei out of ten EPs, which means that CNO nuclei must be accelerated at least ten times more efficiently than H and He. This would have to be increased by another factor of ten if the energy initially imparted to the EPs by the acceleration process were only a factor 2 or 3 lower (i.e. $\sim 3 \times 10^{49}$ erg, which is more reasonable from the point of view of particle acceleration theory). Clearly, more work is needed in this field before one can safely invoke a solution in terms of selective acceleration.

As can be seen, playing with the composition to increase the spallation efficiency has its own limits, and in any case, Fig. 3 gives an unescapable upper limit, obtained with pure Carbon and Oxygen (at least for the canonical spectrum considered here - other spectra were also investigated, as in Ramaty et al. (1997), leaving the main conclusions unchanged). This would then suggest that another source of energy should be sought. However, the constraint that it should be more energetic than SNe is rather strong.

Another interesting line of investigations could be the study of the collective effects of SNe. Most of the massive stars and SN progenitors are believed to be born (and indeed observed, Melnik & Efremov 1995) in associations, and their joint explosions lead to the formation of superbubbles which may provide a very favourable environment for particle acceleration (Bykov & Fleishman 1992). Parizot et al. (1998) have proposed that these superbubbles could be the source of most of the CNO-rich EPs, and Parizot & Knoedlseder (1998) further investigated the gamma-ray lines induced by such an energetic component. The most interesting features of a scenario in which Be-producing EPs are accelerated in superbubbles is that i) when a new SN explodes, the CNO nuclei ejected by the previous SNe are accelerated at the forward shock, instead of the reverse shock in the case of an isolated SN, which implies a greater energy, and ii) no significant adiabatic losses occur, because of the dimensions and low expansion velocity of the superbubble. This makes the superbubble scenario very appealing, and it will be investigated in detail in a forthcoming paper.

However that may be, we should also keep in mind that when we say that a process does not produce enough Be, it always means that it does not produce enough Be as compared to Fe. Now it could also be that SN explosion models actually produce too much Fe. The point is
that Be is compared to Fe in the observations, while it has no direct physical link with it. Indeed, Be is not made out of Fe, but of C and O. So to be really conclusive, the studies of spallative nucleosynthesis should compare theoretical Be/O yields to the corresponding abundance ratio in metal-poor stars. Unfortunately, the data are much more patchy for Be as a function of [O/H] than as a function of [Fe/H], especially in very low metallicity stars. The usually assumed proportionality between O and Fe could turn out to be only approximate, as recent observational works possibly indicate (Israelian et al. 1998; Boesgaard et al. 1998; these observations, however, still need to be confirmed by an independent method, all the more that they come into conflict with several theoretical and observational results; cf. Vangioni-Flam et al. 1998b). We shall address this question in greater detail in the attending paper (Parizot and Drury, 1999, Paper II).

Finally, we wish to stress that the calculations presented in this paper rely on a careful account of the dynamics of the problem. More generally, time-dependent calculations are required to properly evaluate the spallation processes in environments where compositions and energy densities are evolving. In particular, as argued in Parizot (1998), no variation with density can be obtained with a stationary model, since an increase in the density induces an equivalent and cancelling increase in the spallation rates and the energy loss rates. By contrast, we have shown that all three of the processes considered here are more efficient at higher density – a result which could not have been found otherwise. Detailed, numerical time-dependent calculations will be presented in paper II, with conclusions similar to those demonstrated here.

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