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CONF-8406147-4

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TITLE: THE SUPERNOVA ENVELOPE SHOCK ORIGIN OF COSMIC RAYS -
A REVIEW

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LA-UR--84-2517

DE84 016479

SUBMITTED TO XXV COSPAR, Graz, Austria
June 24 - June 30, 1984

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THE SUPERNOVA ENVELOPE SHOCK ORIGIN OF COSMIC RAYS - A REVIEW

IAU/COSPAR
Nucleosynthesis and Acceleration of Cosmic Rays
(6 2.1)

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ABSTRACT

The hydrodynamic shock origin of cosmic rays in the envelope of a Type I presupernova star is reviewed. The possibility of accelerating ultrahigh energy particles to $\gtrsim 10^{18}$ eV is unique to the shock mechanism and currently no other suggested galactic or extragalactic site is likely. Since ultrahigh energy is the most difficult part of any acceleration mechanism, the associated lower energy particles are a priori more likely to dominate all other low energy acceleration mechanisms. The nonrelativistic hydrodynamic supernova explosion shock becomes relativistic at an external mass fraction of $(1-F) = 3 \times 10^{-6}$ of the star that is composed primarily of helium plus heavier nuclei. The resulting ejected relativistic energy, $(1-F) M_{\odot} c^2 \cong 6 \times 10^{48}$ ergs per SNI is adequate to explain the Galactic cosmic ray energy. The resulting spectrum becomes, $N(>E) \propto (1-F) \propto E^{-2.5}$, in agreement with observations. The heavy nuclei are partially spalled in the shock transition and partially resynthesized in the postshock expansion for $E \lesssim 10^{15}$ eV dependent upon the large number of pairs in the post-shock fluid. Above this energy the shock progresses in the magnetized photosphere. The high energy limit is $\cong 10^{21}$ eV due to the coronal density of the presupernova star. The objection to SN shock accelerated cosmic rays by adiabatic deceleration is questioned on the

basis of the Alfvén wave scattering conditions. Ultrahigh energy particles escape because the wave excitation energy density is too low in the dimension of many Larmor radii necessary for scattering back to the SN remnant. Others escape if the energy density is too high. For all others between these two limits the immediately following matter of lower velocity and greater mass compresses and energizes previously trapped higher energy particles, allowing them to escape at energies still higher than originally shock ejected from the supernova. The so-called piston that drives the envelope shock is the same, i.e. the SN bulk ejected matter, or total kinetic energy of the ejecta of ($\cong 2 \times 10^{51}$ ergs) as drives the ISM plasma shock. The efficiency for the envelope shock is (CR energy)/(bulk energy) $\cong 1/300$. For the Alfvén wave ISM shock to have this same efficiency requires that the spectrum of nonrelativistic particles, $E > 100$ keV, $v_{\text{shock}} \cong 3 \times 10^8$ cm s⁻¹, is flatter than $N(>E) \propto E^{-1.62}$. If there are losses in the ISM from the plasma shock such as electron and ion heating and bulk kinetic energy or a steeper nonrelativistic spectrum as in current theories, the hydrodynamic envelope ejecta should dominate the origin of galactic cosmic rays.

I. Introduction

In this paper a review of the work leading to a renewed commitment to the origin of cosmic rays in the shock ejected envelope of supernova is given. The degree to which this interpretation applies to the origin of all cosmic rays is certainly uncertain and does not exclude the possibility of a fraction of the lower energy cosmic rays being accelerated in collisionless plasma shocks in the interstellar medium. On the other hand unless the Alfvén wave shock somehow avoids the investment in energy of a

typical nonrelativistic spectrum, the envelope shock should produce 10^3 to 10^6 more energy in cosmic rays $> 10^9$ eV than the ISM Alfvén wave shock for the same supernova. In addition the SN relativistic envelope shock is currently the only apparently feasible site of ultrahigh energy, $\geq 10^{18}$ eV, particles and so the low-energy spectrum attached to this source must be assumed to exist. Adiabatic deceleration by trapping in the SN expanding bubble from Alfvén wave scattering of particles parallel to the magnetic field assumes that the scattering mechanism is unperturbed except by the particles being scattered. This is clearly not the case for the ejected matter distribution, and so a fraction of the particles will escape with their energy increased as well as decreased.

The plan of this paper is to discuss briefly in section II the uniqueness of the SN envelope acceleration mechanism to produce the ultrahigh energy cosmic rays. This argument requires the likely exclusion of many classical sites for acceleration such as quasars, BL Lac objects, active galactic nuclei, radio sources, intergalactic plasma shocks, pulsars, and the galactic medium shocks. Section III is the main purpose of the paper to review the shock mechanism itself and the work of many people in developing the analytical and numerical simulation of relativistic shock hydrodynamics. The physical structure of such shocks is discussed in IV leading to estimates of the lack of nuclear spallation because of lepton number as well as the magneto relativistic shocks for the high energy tail. The astrophysical setting of the supernova, Type I, with the all important adiabatic deceleration is discussed in V with an estimate of expansion losses as well as gains as a function of cosmic ray energy. In addition a comparison of the efficiency of the acceleration by the envelope shock vs. the ISM plasma scattering shock from the same source of energy, the SN, will be made in Section VI.

II. Alternate Sources of Ultrahigh Energy Cosmic Rays, Photon and Synchrotron Damping

The most obvious sources of cosmic rays more energetic than can be confined by the magnetic field of the Galaxy, $\gtrsim 10^{18}$ to 10^{19} eV, are extra-galactic phenomena, e.g., quasars, active galactic nuclei and radio sources. The extreme emitted energies $\gtrsim 10^{47}$ erg s⁻¹ and short fluctuation times $\lesssim 1$ year of quasars have long implied the necessary conditions for such acceleration. The problem (Colgate, 1983a, 1984) is the energy density in photons necessary to give rise to the observability of the object in the first place. If an energetic proton 10^{19} eV, $\Gamma = 1/\sqrt{(1-\beta^2)} = 10^{10}$ makes one traversal of the light emitting radius to escape the object, it will in general lose by Compton scattering 10^3 to 10^4 times its energy in scattered photons. The photon damping for less energetic objects, AGN and Seyfert galaxies is still greater because of the smaller radius derived from shorter fluctuation times, and hence greater photon density. Acceleration in the magnetic field of a nonobservable AGN with a presumed black hole is still less likely because of the extreme synchrotron losses. Protons of $\Gamma \gtrsim 10^{10}$ must be accelerated in fields $\lesssim 1$ to 10 gauss at radii $\gtrsim 10^{17}$ cm to avoid excessive synchrotron losses per orbit period. (For example, a proton of $\Gamma = 10^{10}$ is stopped by radiation emission in a distance of a micron in a pulsar field of 10^{12} gauss. Curvature from $B_{||}$ acceleration is only less by the ratio of the Larmor to the curvature radius or $\cong 100$.)

Extragalactic Plasma Shock Acceleration

One therefore looks to collisionless plasma shock acceleration in the intergalactic medium. Here as Hillas (1984) has pointed out, the time of

acceleration must be less than the time of Hubble expansion. The plasma shock acceleration depends upon a random walk process. This process takes too long in the weak fields of intergalactic space and finite Alfvén speed. A random walk scattering step is roughly 10 Larmor radii and roughly $4 \times (\omega/v_{\text{Alfvén}})^2$ steps are required to double the energy of a relativistic particle. The requirements of space, time, field, and Alfvén speed far exceeds most observed galactic circumstances, with the possible exception of radio source knots. These objects of sufficient size are so rare and the acceleration circumstances must be so singular, that they seem an unlikely source of ultrahigh energy cosmic rays (Hillas 1984).

Diffusion Into the Galaxy

If one considers particles of less than 10^{18} eV, then diffusion into our Galaxy from an extra galactic source requires that the spectrum of any extragalactic source must be one power of E steeper to create our observed Galactic spectrum, $N(>E) \propto E^{-3}$ for $10^{13} \leq E \leq 10^{18}$ eV. Therefore this implies an extragalactic spectrum of $F(>E) \propto E^{-4}$. This places severe constraints on the total available energy. Plasma shocks within our Galaxy could supply cosmic rays at energies less than 10^{13} to 10^{15} eV (Bell 1978, Blanford and Ostriker 1978, Axford et al. 1977), Lagage and Cesarsky 1983). A different site for ultrahigh energy cosmic rays must be found. The supernova envelope shock still appears the most likely source.

III. Supernova Envelope Shock

The concept that cosmic rays might originate in the supernova explosion has been suggested by many, but most extensively reviewed by

Ginzburg and Syrovatskii (1964). However, the possibility that high-energy matter might be ejected from the outer layers of a supernova in magnitude and energy distribution consistent with what is observed in cosmic rays was originally put forward in Colgate and Johnson (1960). The general concept in that paper was that matter should be ejected from the outer layers of an explosion into the relative vacuum of space with energies far greater than the average energy of the explosion. In particular, a shock was assumed of a strength reasonably consistent with the average velocity for a supernova explosion and then the behavior of such a shock wave in the envelope of a presupernova star was analytically continued in the decreasing density of the stellar envelope. The greatest uncertainty in this was the behavior of the supernova itself and this led to the work of Colgate and White (1966) on the mechanisms and explosion history of supernovae.

SN and the Formation of the Envelope Shock

When a star explodes, the energy that is released in the matter is due to either neutrino deposition, core reflected bounce, or thermonuclear reactions. In general it occurs in a time that is short. The traversal time of sound of the core is short compared to the traversal time of sound of the lower density exterior layers of the star. This is particularly true of massive stars, $\gtrsim 10 M_{\odot}$ characteristic of Type II supernova, and is also true for Type I supernova in a thermonuclear detonation (Arnett 1977), or core collapse. This single criterion of sudden release of energy means that the pressure wave that is expanding and ejecting the matter of the star would become a shock wave in the lower density, lower temperature, outer layers of the star. Indeed in the core collapse

models, the shock wave starts at the boundary of a neutron star core and continues throughout the star. All numerical calculations of supernova that eject mass exhibit the formation of a strong shock. It is the behavior of this shock wave in the relatively static and decreasing density of the envelope that forms the basis of presuming that matter could be ejected from a supernova with the relativistic energies characteristic of cosmic rays.

Nonrelativistic Shocks

The nonrelativistic behavior of shocks in density gradients has been analyzed in detail by many hydrodynamicists and the general property of a strong shock speeding up or becoming stronger inversely proportional to a power of the density is well recognized in the literature (Chisnel 1955, Rogers 1957, Kopal 1954).

The numerical confirmation and investigation of some of these solutions was initiated in Colgate and White (1966) and furthered by Grover and Hardy (1966). The excellent agreement between numerical hydrodynamics and analytic solutions in both these papers is strong confirmation of the numerical hydrodynamic procedures used to calculate the behavior of a strong shock in the envelope of a supernova star. Although the peculiar property of a shock speeding up in the envelope has been repeated in other hydrodynamic codes (Imshennik and Nadezhin 1970), there has not been published a calculation of a SN shock in a stellar envelope in conjunction with a known analytical test problem since Colgate and White. We must then rely on these earlier calculations as the single best evidence for the shock velocity as a function of mass fraction of the presupernova star.

Type I SN and the Envelope Nonrelativistic Shock

The most likely source of high-energy cosmic rays from supernova are from supernova of Type I. The behavior of the light curves (Woosley, Axelrod and Weaver 1980, Colgate, Petschek and Kriese 1980) suggest supernova models that give the best agreement with the analysis of the light curves are those with an exterior velocity of $\cong 1.5 \cdot 10^9$ cm/s. This is the velocity of the outer boundary of the simplified model of a uniform density sphere that is used to explain the light curve. In Woosley, Axelrod and Weaver, the speedup or higher velocity for the external mass fractions in a model with density gradients was qualitatively reconfirmed. The velocity distribution leaves a frozen-in density distribution that is invoked to explain the spectra from the radioactive

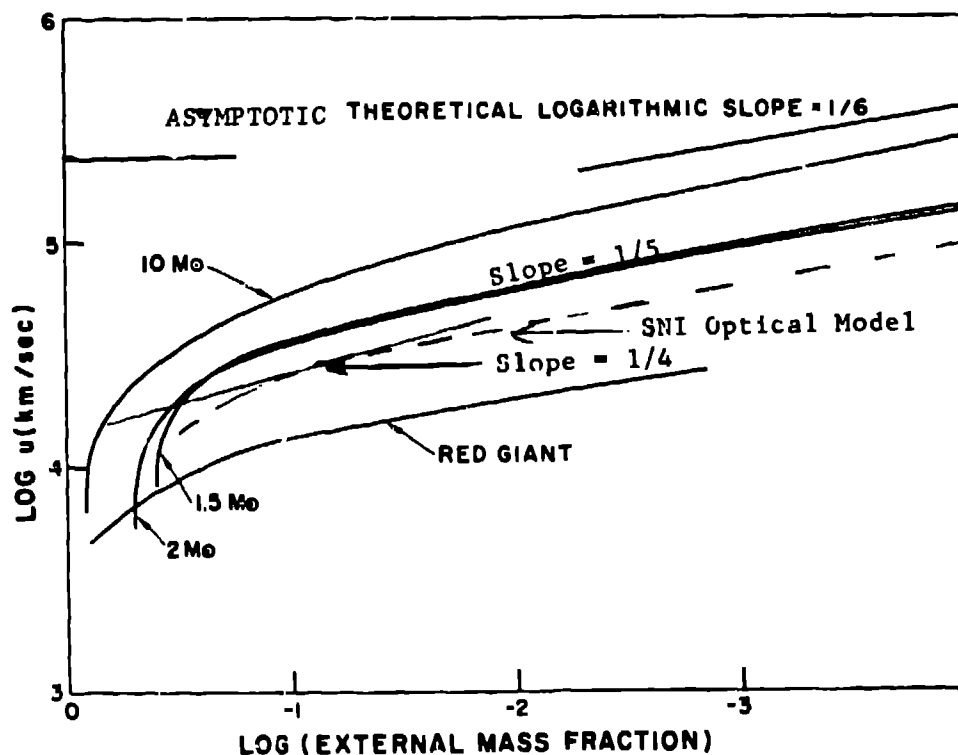


FIG. 1 — Log expansion velocity versus log mass fraction for supernova envelopes of 10, 2, and 1.5 M_{\odot} polytrope 3 initial stars and a red giant. The logarithmic slope corresponds to the theoretical value of $\frac{1}{6}$. The dashed curve corresponds to current

Figure 1

energy of the $\text{Ni}^{56} \rightarrow \text{Co}^{56} \rightarrow \text{Fe}^{56}$ decay. In Fig 1 we show the velocity distributions calculated in Colgate and White for various mass stars and total energy. The curve marked SNI is the best fit to the current models needed to explain the optical light curve and spectra. The abscissa in this case is no longer the familiar Lagrange coordinate of the internal mass fraction, F , but instead the external mass fraction, $(1-F)$, of the star. The reason for using this coordinate is that it becomes a measure of all the matter external to the point in question. Since the shock is speeding up as it proceeds outward through the exterior mass distribution of the star, such an external mass coordinate means that all matter exterior to this coordinate is moving with a higher velocity and hence higher energy. It is a coordinate analogous to the familiar cosmic ray coordinate $N(>E)$ for the cosmic ray energy distribution. If we express the external mass coordinate, $(1-F)$, in terms of the final energy of ejected matter, E , then $(1-F)$ is proportional to $N(>E)$. We will use this transformation to calculate an expected cosmic ray spectrum. The exponent of the power law in velocity behavior that averages the nonrelativistic behavior $0.1 \leq (1-F) \leq 3 \times 10^{-6}$ of the numerical models of Colgate and White is $-1/4$ and approaches $-1/5$ for $100 \text{ MeV} \leq E \leq 1 \text{ GeV}$. This final stage is close to the analytical solution which is $-1/6$ for a plane parallel shock. For a uniform density expanding sphere, $v \propto R \propto F^{1/3}$. A SN explosion is spherical inside and planar outside. Hence for $(1-F) < 1/10$ we obtain:

$$v_{ej} \propto (1-F)^{-1/5} \quad (1)$$

or

$$N(>E) \propto (1-F) \propto E^{-2.5} .$$

Since we start at the boundary of the uniform expansion region with velocities that are 1/20 that of light, a factor of ten increase in ejected velocity will become close to relativistic. It was estimated in Colgate and White, and scaled to current SN Type I models that the mass fraction corresponding to relativistic ejected matter would be roughly 3×10^{-6} of the ejected mass. It should be emphasized that this does not mean that the shock has this velocity or strength in the envelope at this mass fraction, but instead ejected matter reaches this velocity after it has expanded in the general expansion that occurs following the shock.

Observational Evidence for Shock Velocity Distribution

Before discussing this expansion, it is useful to point out that there is one observation that confirms this velocity distribution. Branch (1982) has pointed out that the early spectra near maximum of SNI's can be explained with excellent agreement with observations as a near blackbody with absorption lines formed by an overlying lower density photosphere. The density distribution of this photosphere that gives the best fit to the observations is one with a density $\rho \propto R^{-7}$.

If we model the ejected velocity distribution with

$$v_{ej} = v_o (1-F)^{-\alpha}$$

$$\text{or } (1-F) \propto (v_{ej}/v_o)^{-1/\alpha}$$

$$\text{or } d(1-F) = 4\pi R^2 \rho dR = (-1/\alpha) (v_{ej}/v_o)^{-1-1/\alpha} \left(\frac{v_{ej}}{v_o}\right) \text{ but } R = v_{ej} t$$

Therefore at a given time

$$\rho \propto R^{-3-1/\alpha}, \text{ or if } \alpha = \frac{1}{2}, \rho \propto R^{-7} \quad (2)$$

In Fig. 1 we have drawn a line tangent to the SNI best fit curve at $(1-F)=1/20$, the approximate location of the photosphere at light maximum with the slope $\alpha = 1/4$. The agreement is strong evidence of the shock wave speedup in the supernova envelope.

Expansion Velocity Increase

The post-expansion increase in velocity is roughly a factor of two nonrelativistically, greater than would be expected by purely local energy conservation. (In a strong nonrelativistic shock the internal energy equals the kinetic energy of the motion of the matter behind the shock. Therefore, expansion would be expected to lead to an increase in velocity by $\sqrt{2}$.) The larger velocity increase is due to the pressure gradient that exists in the matter during expansion. When the shock becomes near relativistic at a value $\Gamma = 1/\sqrt{1 - \beta^2} = 2$, the increase in kinetic energy following postshock expansion is close to tenfold. The mass fraction of the relativistic boundary of 3×10^{-6} is a reasonable estimate for any compact presupernova star. If there is an extended envelope, as might be the case for a red giant, or the high-density stellar winds recently inferred for explaining the ultraviolet emission from Type II supernova (Fransson 1983), such a shock would not speedup in such an extended mass distribution. The particular velocity versus mass fraction law, Eq. 1, applies to a monotonically decreasing density in a stellar envelope that follows the classic "radiative zero" solution for stellar envelopes, Schwarzschild (1958). In such an envelope the density,

$$\rho \propto (1-F)^{3/4} ; \quad T \propto \rho^{1/3} ; \quad h \cong \left(\frac{1}{\rho} \frac{d\rho}{dr} \right)^{-1} \propto T ; \quad h\rho \propto 1-F \quad . \quad (3)$$

Total Energy of Relativistic Matter

Once we have established the mass fraction of matter whose final kinetic energy has an energy factor $\Gamma > 2$, it is simple to estimate the total energy ejected in such matter. This becomes $M_{ej} (1-F) \tau^2 = 6 \times 10^{48}$ erg, close to what is required ($\cong 10^{49}$ erg) needed to fill the galaxy with cosmic rays with one SNI per 30 years and a cosmic ray life time of 2×10^7 years.

The Two Objections

There are two often repeated problems with the association of this relativistic matter with the cosmic rays observed on earth. The first is the objection raised by Ginzburg and Syrovatskii (1964) that all heavy nuclei would be spalled in the shock transition. This objection will be discussed in terms of shock structure in the presence of a high lepton density. The second objection, adiabatic deceleration (Kulsrud and Zweibel 1975), will be discussed in the last section of this review.

IV. Initial Envelope Conditions for the Relativistic Shock

The shock conditions in the presupernova star, at the point in the envelope where relativistic matter is expected to arise is determined by the external mass fraction 3×10^{-6} of the star. Recently an accretion model of a white dwarf at $1.38 M_{\odot}$ has been calculated in detail by Starrfield, Sparks and Truran (1984). This mass is reached by accretion just before the initiation of collapse by electron capture, with roughly a factor of 2 increase in density yet to occur before the outer layers are ejected in the SN explosion. The conditions they calculate for this model after reaching equilibrium are: radius = 1.6×10^8 cm and for the

envelope mass fraction $1-F = 3 \times 10^{-6}$, $5 \times 10^4 \text{ gm cm}^{-3}$ and $T = 2.4 \times 10^7$ degrees. These conditions are well beyond hydrogen burning which occurs at a mass fraction of $\cong 10^{-11}$. Consequently the ejected matter will be helium with heavier nuclei characteristic of the companion star. The proton fraction of cosmic rays will be the residue of the accumulation of spallation in the galaxy. We therefore choose the initial conditions for the envelope shock at the mass fraction $(1-F) = 3 \times 10^{-6}$, $\rho = 10^5 \text{ g cm}^{-3}$, scale height = 5 km, thickness $5 \times 10^{10} \text{ g cm}^{-2}$ and radius $1.5 \times 10^8 \text{ cm}$. Later we will show that the minimum thickness necessary for the propagation of the relativistic shock corresponding to the dynamic friction of the lepton fluid is 1 g/cm^2 . Then the relativistic shock can propagate through nearly eleven orders of magnitude change of mass fraction before breaking out of the surface of the star. It is this eleven orders of magnitude of change in mass fraction or density change of 10^8 that gives rise to a spectrum of relativistic ejected matter that closely follows that observed for cosmic rays.

The Relativistic Spectrum

Relativistic hydrodynamics is, in many ways, very similar to the standard nonrelativistic hydrodynamics. Numerical codes can be written that reproduce analytical solutions and therefore can be extended to cases that are nonanalytical. This has been the case with the relativistic hydrodynamics associated with the supernova shock in the envelope. The analytical approximations of Colgate and Johnson showed that the method of characteristics developed by Chisnell (1955) could be applied to the relativistic case with certain simplifications inherent to relativistic hydrodynamics. In particular the sound speed in a relativistic gas is a

constant, $c/\sqrt{3}$. The matter behind the shock is connected to the shock by sound characteristics of constant value. This leads to analytical solutions. This has proven to be the case in relativistic shocks in density gradients. The solution given in Colgate and Johnson (1960) and later in significantly greater detail by Johnson and McKee (1971), is that the shock strength, $\Gamma_s = 1/\sqrt{1-\beta_s^2}$, $\beta_s = (\text{fluid velocity})/c$ behind the shock increases as the rest mass density as:

$$(\Gamma_s - 1) = (\rho/\rho_0)^\alpha$$

$$\alpha = \sqrt{3}/[2(2+\sqrt{3})] = 0.232 \quad (4)$$

The shock strength is therefore a relatively weak power of the rest mass density, $\rho^{-.232}$, and so naively we might expect a relatively steep and uninteresting spectrum of ejected matter.

Relativistic Postshock Expansion

However, the expansion of the post shock fluid is different in relativistic hydrodynamics, because the internal energy (specific energy density = $\rho_s \Gamma_s c^2$) is comoving with the rest mass, which is also moving at Γ_s behind the shock. Hence the mass density of the internal energy is greater by $\times \Gamma_s$ than the rest mass of the comoving matter. This means that the total energy per unit rest mass measured in the star frame behind the shock is far greater, Γ_s^2 than the kinetic energy, Γ_s , of the rest mass alone, and hence subsequent expansion of the comoving relativistic fluid leads to a significantly greater velocity or kinetic energy, Γ_{final} of the fluid than might have been thought possible from the rather weak and uninteresting behavior of the shock itself.

Similar to nonrelativistic hydrodynamics, the internal energy in the comoving frame per unit of rest mass is exactly equal to the kinetic energy of that same rest mass and hence the internal energy measured in the laboratory frame has the total energy of Γ_s^2 . One would therefore expect a final energy of the rest mass after expansion, Γ_{final} , to be some factor greater than this, and indeed, analogous to the nonrelativistic expansion, the multiplication in energy of the rest mass energy factor becomes $\Gamma_s^{(1+\sqrt{3})}$. When the above expression is combined with the shock strength as a function of density, one obtains the final energy, $\Gamma_{\text{final}} \propto \rho^{-0.634}$. For the "radiative zero" density distribution of the envelope, the external mass fraction (1-F), is proportional to $\rho^{4/3}$. Therefore one obtains:

$$\Gamma_{\text{final}} = \Gamma_{\text{shock}}^{(1+\sqrt{3})} \alpha (1-F)^{-.48} \quad (5)$$

This then leads to the cosmic ray spectrum:

$$N (>E) \propto (1-F) \alpha \Gamma_{\text{final}}^{-2.10}$$

The shock strength as a function of mass fraction has been derived by Johnson and (Chris) McKee and confirmed in calculations by (Chester) McKee and Colgate (1973), Fig. 2. Eltgroth considered the purely spherical case, but the supernova case is initially planar and partially spherical in the expansion. Chester McKee developed a shock following code using the sound characteristics to determine the shock position and verified in great detail the shock behavior in the density gradient of the star. In a separate calculation using the velocity and internal energy associated with such a shock and the density gradient, he performed the expansion calculation with the appropriate planar and then spherical expansion

EARLY RADIATION SPECTRUM

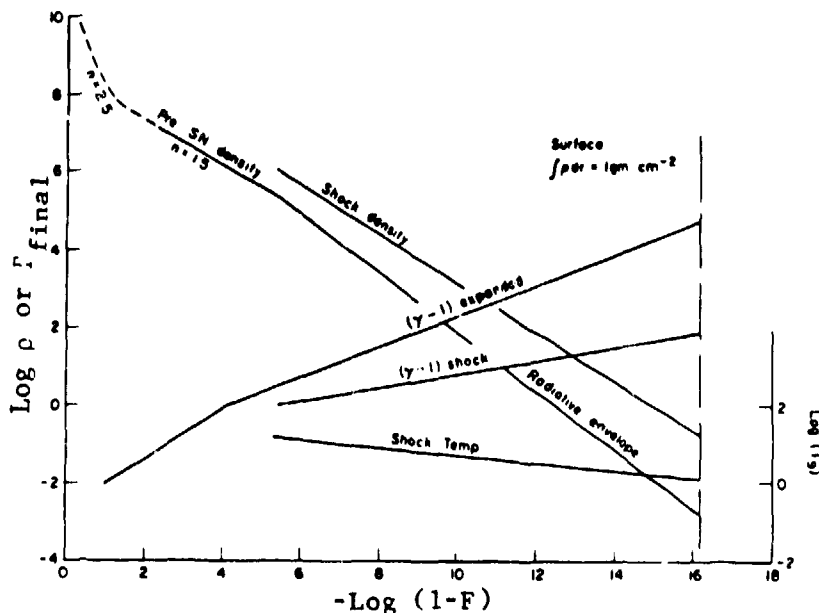


FIG. 2—The presupernova density is shown as a function of mass fraction. The inner core is approximated by a polytrope of index 2.5 and mass $1.5 M_{\odot}$ with a radiative envelope and radius of 10^8 cm. The energy factor of the fluid velocity immediately behind the shock is shown, as well as the corresponding energy factor after expansion. The temperature in units of 10^7 K of the fluid behind the shock is shown with the scale at the right.

Figure 2

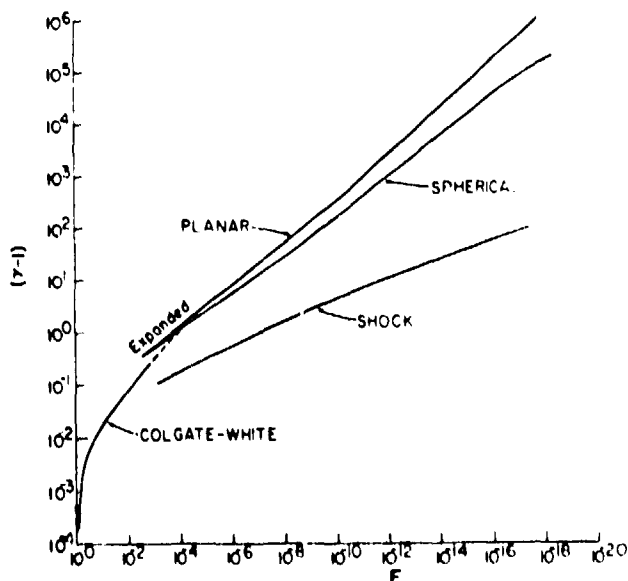


FIG. 3—The kinetic energy factor $(\gamma - 1)$ of the envelope material versus external mass fraction for the nonrelativistic calculations of Colgate and White matched onto the recent calculations. The lowest curve is the energy factor immediately behind the shock before any expansion has taken place. Since the outer polytropic layers are thin compared with the stellar radius, the plane-parallel approximation is valid and the slope derived from the numerical calculations of 0.175 compares favorably with the theoretical one of 0.178. The upper curve corresponds to a numerical calculation of a plane parallel expansion, and the slope 0.46 compares with a theoretical value of 0.48. The actual polytropic expansion which initially appears as plane parallel ultimately involves a relatively large change in radius ($\times 10$); and results in a slope of 0.40. A value lower than 0.48 is anticipated by Eltgroth (1971).

Colgate, Blevins and McKee (1972)

Figure 3

inherent to the changing geometry of the expansion, Fig. 3, (Colgate, McKee and Blevins, 1972). In this case $\Gamma_{\text{final}} \propto (1-F)^{-0.40}$ and the cosmic ray spectrum becomes: $N > E \propto E^{-2.5}$ as observed.

Both these calculations confirm the analytical work and showed the small correction necessary for the late spherical expansion of the ejecta. Additional relativistic shock analyses have been performed by Shapiro (1979), Blanford, and (Chris) McKee (1976). Figure 2 shows the density, shock strength, and final expansion energy factor for a typical envelope over the region in which the hydrodynamics and the shock structure should be reasonably "normal," i.e. depending only upon collisional dynamic friction. Figure 3 shows the comparison of planar and partially spherical expansion.

In these calculations the thickness of the matter everywhere is assumed great enough such that the dynamic friction insures equilibrium fluid properties as opposed to the free streaming of matter. With this definition of normal hydrodynamics the surface of the star is then roughly 1 g/cm^2 thick. This thickness means that we expect a proton whose kinetic energy relative to the fluid behind the shock is $\Gamma_s m_p c^2$ will have a range in the hot, high-lepton density post-shock matter of 1 g/cm^2 of the original rest mass.

V. Shock Structure

Shock structure is important for defining the surface of the star where the normal shock hydrodynamics breaks down as well as the shock structure in the region of $\Gamma_{\text{final}} = 1$ to 1000 where the composition of cosmic rays is normally measured. A related problem of diffusion occurs in the release of radiant energy from the comoving matter during the large

expansion (expansion ratio $\cong 100 \Gamma_{\text{g}}^3$) that must take place before all the internal energy is converted to kinetic energy. During this expansion nuclear, radiation and lepton pair processes are taking place. Also during the time of expansion one might expect radiation transport to possibly alter the velocity trajectories calculated on the basis of adiabatic processes.

Radiation Diffusion and Uncovering During Expansion

Colgate and Petschek (1979) considered this problem and showed that in the comoving frame of the fluid the internal energy was uncovered by what one might call a radiation transparency wave. In other words the transparency occurred in the comoving frame so suddenly that negligible diffusion occurred before the sudden transition to transparency. This is because the surface corresponding to $\tau = 2/3$ moves relative to the fluid at close to the velocity of light. Diffusion in more than a few mean free path occurs significantly slower than the velocity of light, $c/3\tau$. Therefore the uncovering is a sudden event, and radiation diffusion does not redistribute the energy in such a fashion as to alter the velocity versus mass distribution. Instead the expansion process is truncated at the point of transparency and a somewhat steeper velocity distribution results. The equations of coupled radiation and hydrodynamics dealing with this problem are discussed in Glaviano and Raymond (1981).

Shock State Conditions at the Relativistic Mass Fraction

As stated earlier the expanded matter following the explosion shock traversing the external mass fraction 3×10^{-6} of the presupernova star becomes relativistic after speeding up in the expansion by a factor of

$\cong \times 10$ in energy. This mass fraction corresponds to a layer of the pre-supernova star initially roughly 5 km thick, density 10^5 g cm^{-3} and $5 \times 10^{10} \text{ g/cm}^2$ thick. The energy density behind the shock in this matter will therefore be $\rho_0 \eta_{\text{shock}} c^2/10$ or $6 \times 10^{25} \text{ erg/cm}^{-3}$. The compression ratio, η_{shock} behind a strong shock is seven-fold nonrelativistically for $\gamma = 4/3$ and is $4 \Gamma_{\text{shock}}$ for a relativistic shock ($\Gamma_s \gg 1$) as measured in the comoving frame. In the relativistic limit in the laboratory frame we would see a layer of matter passing us compressed by $\cong 4 \Gamma_{\text{shock}}^2$. We are concerned here with conditions in the comoving frame for a shock at strength, $c^2/10$ with a compression ratio of seven. Hence $\rho_{\text{shock}} \cong 7 \times 10^5 \text{ g cm}^{-3}$ in the comoving frame. The resultant energy density of $7 \times 10^{25} \text{ erg/cm}^3 \cong 2aT^4$ results in a temperature of 8×10^9 degrees or $kT = 1.4 mc^2$. (The coefficient, a , is increased by $\cong \times 2$ because of pairs.) It may be that a smaller mass fraction or different stellar structure reduces the density to 10^5 g cm^{-3} and hence the temperature to $kT = mc^2$, $T_g = 6$. The conditions assumed in Colgate 1975.

Nuclear Spallation From Thermal Photons

These state conditions are near the point of rapid thermal decomposition of most nuclear species. However, the time during which this state condition lasts is short. The time for expansion in the comoving frame is roughly the compressed local scale height $1.4 \times 10^4 \text{ cm}$ divided by the sound speed ($c/\sqrt{3}$) or $\cong 10^{-6}$. This time is too short for total thermal decomposition to take place by the thermally generated gamma rays by $\gamma - N$ and $\gamma - P$ reactions. The proton and neutron capture will also be important. A rough estimate for alpha particle nuclei is a threshold of $\cong 10 \text{ MeV}$ and the resonance at 15 MeV or at roughly 21 kT for

$T_g = 8.3$. The cross sections are $\cong 10^{-25} \text{ cm}^{-2}$. The number density of pairs and photons at this temperature is $\cong 10^{31} \text{ cm}^{-3}$. At the resonance for spallation, 15 MeV, $n_\gamma \cong 10^{24} \text{ cm}^{-3}$. (The pair density is 100 times the alpha particle density.) Then $n_\gamma \sigma \cong 3 \times 10^9 \text{ s}^{-1}$. In order for the time for thermal decomposition to become longer than the expansion time, the temperature would have to be less than $t_g = 6$. (Multiple state excitations and the proton and neutron capture have not been calculated.) Therefore we expect a temperature and density in the range $6 \leq T_g \leq 8$, $10^5 \leq \rho \leq 7 \times 10^5$ where thermal destruction of high atomic number species, is within the range of error of current cosmic ray composition measurements. Spallation of high atomic number nuclear species can also be caused by a collision density in the nonthermalized region of the shock. This depends upon shock structure.

Shock Structure and Pair Density

The thickness of the shock wave is characterized by a sequence of relaxation processes that terminates in thermodynamic equilibrium at some distance behind the precursor of the shock wave. The various relaxation processes include thermal diffusion as well as the dynamic friction between the mass containing species and the fluid behind the shock. Most frequently we think of radiation diffusion as spreading a radiation dominated shock leading to the longest characteristic length (Weaver, 1976 and Chapline and Weaver, 1979), but since our shock velocities are close to the speed of light, thermal diffusion lengths will be small, a few radiation mean free paths, compared to the dynamic friction lengths of the slowing down of the ions that constitute the rest mass of the fluid. Here we are concerned with the slowing down length governed by the fluid

properties which in turn are determined by the heating from the dynamic friction of the various ion species. The detailed modeling of this coupled radiation diffusion hydrodynamic and viscosity problem has not been performed, but one can understand the order of magnitude of the result as the slowing down length of the mass containing atomic number species. If this is hydrogen or helium, other high atomic number species will have a velocity relative to the mass containing ones because of a different dynamic friction proportional to Z^2/A .

We then ask for the collision density of the heavy ions resulting in spallation in this relative slowing down process. We assume one component, protons and alphas, define the center of mass frame and therefore are at rest and the other component, heavy ions, slow down within their classical dE/dx range. The slowing down length $E/(dE/dx)$, is determined by the dynamic friction between a near relativistic heavy ion and a high temperature relativistic fluid characterized by the post shock fluid density and temperature.

We then wish to know the effective spallation of a heavy ion undergoing the slowing down process. The slowing down length in normal matter would be long enough such that almost complete spallation of the heavy nuclei would take place. This has been cited (Ginzburg and Syrovatskii 1964), as the strongest argument against the stellar envelope shock hydrodynamic origin of cosmic rays. This misconception was addressed by Colgate (1980, 1981a,b), where it was pointed out that the lepton number density due to pairs in the comoving fluid is so high that the dynamic friction dE/dx is much greater than normal matter so that the heavy nuclei come to rest in a pair dominated fluid and hence undergo less spallation.

We quantify this by calculating spallation in the nonrelativistic shock in the external mass fraction 3×10^{-6} of the stellar envelope where relativistic, $\Gamma_{\text{final}} = 2$, cosmic rays first originate. Here the relative energy across the shock is roughly 1/10 the final value, or $m_p c^2/10 = 10^8$ eV/nucleon kinetic energy. The relative energy between protons or α 's and heavies might be half of this. Let us consider a carbon nucleus slowing down in such matter.

Stepney (1983) has considered two-body relaxation in relativistic thermal plasmas. The differences in the relaxation time for proton energies equal to the electron temperature and proton velocity equal to the electron thermal velocity is within a factor of two up to several mc^2 electron temperature. This temperature exceeds our range of interest and so we approximate relaxation of the alpha particles as midway between the two limits and at the temperature of $1.4 mc^2$. For this temperature and $\ln \Lambda = 5$, the relaxation becomes $\tau_r = 10^3 / (n_e \sigma_T c)$ seconds. Then the slowing down length at 50 MeV per nucleon or $v = c/3$ becomes $300 / (n_e \sigma_T)$. The spallation cross section for α 's on heavier nuclei is roughly twice that of protons so that $\sigma_s \cong 0.6 \times 10^{-24} \text{ cm}^2 \cong \sigma_T$ so that the fractional spallation becomes $s = n_\alpha \sigma_s \ell = 300 (\sigma_s / \sigma_T) (n_\alpha / n_e)$. We have already calculated the ratio of pairs to α 's behind the shock as $n_\alpha / n_e \cong 1/100$, so that $s = 3$ nucleons removed. This would be a significant spallation if burn-back by the free neutrons and protons was not very rapid. The neutron capture in neutron depleted nuclei is more rapid than the expansion time (8×10^{-7} s) even at unity density so that at $\rho_s = 7 \times 10^5 \text{ g cm}^{-3}$, capture will be near instantaneous. Similarly proton burn-back for some nuclei is also very rapid. A detailed calculation has yet to be performed.

Spallation and Higher Energies

Thereafter as the shock wave speeds up, the lepton to alpha ratio increases as $T_{\text{shock}}^3 / \rho_{\text{shock}}$, the path length decreases as $T^{3/2}$, and the fractional spallations become $\propto T_{\text{shock}}^{3+3/2} / \rho_{\text{shock}} \propto (1-F)^{0.1}$ and hence decreases with increasing energy. Hence the fractional spallation of heavy nuclear species in the outer layers of a presupernova star should be less than that occurring in the propagation in the ISM.

The cosmic ray composition ejected in a supernova should reflect the composition of the matter accreted from the assumed binary companion before the supernova Type I event and then processed through hydrogen burning. This composition will be partially spalled and the neutrons and some protons will be recaptured during expansion and cooling. The primary composition would be almost all helium. The protons of cosmic rays are formed in subsequent spallation in the ISM. The propagation, spallation and final storage of protons in the galaxy has not been calculated, but should be similar to the problem calculated by Peters and Rasmusen (19??) for injection of pure iron.

High Energy Spallation

The acceleration of very high cosmic ray energies comes from small external mass fraction and hence low density. Then the lepton number density decreases more rapidly than T^3 when the temperature immediately behind the shock drops below roughly $mc^2/2$, which occurs when $\Gamma_s \cong 12$, and $\Gamma_{\text{final}} \cong 10^3$. Spallation then increases and we expect it to become important at about the time when the shock breaks out the surface layer of a star at 1 g/cm^2 thick and a mass fraction of $3 \cdot 10^{-17}$, or 10^{-11} of the original mass fraction where relativistic matter was first created. This

corresponds to the shock strength, $\Gamma_s \cong 100$ and a final expanded energy of Γ_{final} of 10^6 or 10^{15} eV per nucleon.

Limiting Hydrodynamic Shock

This outer layer where the shock conditions break down is defined where the range of the proton in the shock fluid conditions at a $\Gamma_s = 100$ is just the thickness of the layer itself, namely 1 g/cm^2 . This corresponds to a pair density of roughly $3 \times 10^4 \times$ the baryon density at a temperature of $T_9 \cong 1$, $mc^2/5$, in a rest mass density of 1 g/cm^3 . Here we use the classical relativistic slowing down dynamic friction because $v_{\text{proton}} \gg v_{\text{electrons}}$. This is 4 MeV/g cm^{-2} (hydrogen) increased by the lepton density, or $1.2 \times 10^{11} \text{ eV/g cm}^{-2}$ at the surface. For the shock to propagate further requires a magnetic field strong enough such that the particles remain local in the accelerated fluid frame.

Shock in a Magnetic Field and the Ultrahigh Energy Matter

For an orthogonal shock the compression ratio of the magnetic field will be the same as the compression ratio of the fluid. Therefore the ratio of Larmor orbit to scale height remains constant for a constant Γ_s . Since the scale height of the outer layer, 1 g/cm^2 thick, is that of the optical surface or roughly 10 meters, the magnetic field strength necessary to confine a proton to this dimension is 10^5 gauss at $\Gamma_s = 100$. The magnetic flux of such a field in a compact white dwarf at 10^8 cm radius is relatively small, 10^{-3} , of the canonical value that makes a neutron star with 10^{12} gauss. This is also a relatively typical magnetic field of white dwarfs. In such a field the relativistic shock will propagate primarily as a pair fluid with an occasional baryon that remains

local to the fluid element because of the presence of the magnetic field. The compression of the magnetic field by a factor approximately $4\Gamma_g$ means that the magnetic energy soon exceeds other forms of internal energy behind the shock so that the expansion ratio required after the shock to recover the internal energy becomes much smaller. The effective adiabatic index γ is now 2 rather than $4/3$. In addition the scale height of the photosphere slowly increases from 10 m at a surface temperature of 10 eV to 50 m at 80 eV, corresponding to a typical corona. The energy spectrum should then be at least as flat as a shock in a constant scale height exponential density gradient or $\Gamma_{\text{final}} \propto (1-F)^{-0.48}$, or $N(>E) \propto E^{-2.08}$. This is a flatter spectrum than at lower energy from the polytropic envelope density distribution.

Upper Energy Limit

The highest energy CR's should correspond to where the exponential density distribution flattens out to a typical stellar wind. If we set this at a relatively high value like $n_e \cong 10^9 \text{ cm}^{-3}$, $\rho_{\text{corona}} = 10^{-15} \text{ g cm}^{-3}$, then the magnetospheric shock will propagate in a density ratio of $\rho_{\text{surface}} = 10^{-3} \text{ g cm}^{-3}$ to $\rho_{\text{corona}} = 10^{-15}$, or 10^{12} change in density, or $10^{12/2.08} = 10^{5.77}$ increase in energy. Γ_{final} at the surface is $10^{5.4}$ so that the maximum energy is $\Gamma_{\text{final}} = 10^{11}$ or 10^{20} eV. The thickness behind the shock of the mass fraction of the maximum energy associated with the initial scale height of 50 meters will be 10^{-11} g/cm^2 or roughly $4 \times 10^{-15} \text{ g/cm}^{-3}$ density in the comoving frame. The shock energy factor $\Gamma_B = 2.5 \times 10^4$. The initial rest mass density of this mass fraction is considerably less than the rest mass of the magnetic field, $B^2/8\pi c^2 = 4 \times 10^{-13} \text{ g/cm}^2$, so that the shock Hugoniot relations will be different.

The shock energy factor, $\Gamma_s = 2.5 \times 10^4$, is so large that the compressed magnetic field will be nearly the total comoving internal energy. However since the internal energy will include some pairs, and photons, such a shock should be opaque to photons, and therefore similar to the nonmagnetic case. The limitation to this process and the upper energy limit of the cosmic rays will most likely be determined by the breakdown in locality conditions and the expansion necessary to convert the internal energy back to kinetic energy of the ions.

Magnetic Shock Expansion

This expansion is initially plane parallel and later will become spherical. The expansion of magnetic field energy density is far more efficient and requires a smaller expansion ratio ($\gamma = 2$ plane parallel). The magnetic energy will be converted to kinetic energy for a volume expansion ratio of $\Gamma_s^{\gamma-1} = \Gamma_s$ rather than proportional to Γ_s^3 for $\gamma = 4/3$. The magnetic shock expansion will therefore be planar rather than spherical. The advantage of this high energy acceleration mechanism is that the particles are accelerated in a two-step mechanism, (1) the shock where the energy change per particle is relatively modest, Γ_s of 100 to 10^4 , and (2) a subsequent slower expansion where the dynamic friction in the comoving pair fluid or magnetic field assures that the particles stay in step with the fluid. In addition finite Larmor orbit effects are minimal as well as radiation damping like synchrotron and Compton scattering. The natural upper limit of this acceleration will be determined by the chromosphere structure where a wind of constant density of the order of 10^9 particles per cm^3 or 10^{-15} g/cm^3 dominates. Then the

density gradient no longer exists and the shock instead of speeding up starts to decelerate.

This then becomes the point where the formation of the so-called "bubble" of Kulsrud and Zweibel becomes important. Hence we now discussed the possible trapping and escape from the expansion of such a bubble and the expected cosmic ray flux which should become part of the interstellar medium.

VI. Adiabatic Deceleration

The concept of adiabatic deceleration as presented by Kulsrud and Zweibel (1975), Kulsrud (1978, 1979 and 1982), and Foote and Kulsrud (1979) is that the ejection of high energy matter from the supernova explosion blows a field-free bubble in the interstellar medium and the pressure boundary of such a bubble removes the energy of the expanding matter by PdV work. In turn this energy shows up as a shock wave in the interstellar medium. This shock is later invoked to produce the cosmic ray spectrum by Alfvén waves scattering in the ISM. One might naively expect that particles whose trajectory is initially parallel to the average magnetic field of the interstellar medium might escape. Instead Zweibel and Kulsrud demonstrate that such particles excite Alfvén waves which in turn lead the scattering and their ultimate diffusion back into the bubble. This picture of adiabatic deceleration by Alfvén wave scattering has received relatively wide acceptance despite objections by Holman, Ionson and Scott (1979) who claim that the velocity distribution of the particles exciting the Alfvén waves makes a major difference to the strength of the waves and hence their scattering properties. Here it is

suggested instead that for the typical SN power law velocity distributions, the excitation of Alfvén waves would be time and spatially dependent such that a higher energy mass fraction is trapped and compressed by the larger mass, lower energy following mass fractions.

For a particle to be scattered by an Alfvén wave, the dimensions of the wave must be significantly greater than a Larmor orbit of the particle to be confined. In strong turbulent cosmic ray shock acceleration, this ratio is roughly 10 and the mean number of scatterings necessary to reverse the direction of a particle will be at least 3 to 4 scatterings. Therefore particles whose kinetic energy is large enough such that their Larmor orbit is greater than $\cong 3\%$ of the dimensions of the Alfvén shock or roughly 100 pc will certainly escape, or $E > 3 \times 10^{16}$ eV. (The Alfvén wave shock acceleration mechanism folds up at a lower energy, 10^{13} to 10^{15} eV (Lagage and Cesarsky, 1983).

We start by considering the escape of the highest energy particles and then the problems of confinement of progressively lower energy ones. For a particle to be returned to the field-free debris bubble requires:

(1) The Larmor orbit R_L must be $\cong 1/10 \lambda_{\text{Alfvén}}$ and $h_{\text{galaxy}} \geq 3 \lambda_{\text{Alfvén}}$. Therefore $h_{\text{galaxy}} > 30 R_L$ where h_{galaxy} is the scale height or the thickness $\cong 1/2$ kpc. Therefore for $\Gamma > 1.7 \times 10^8$, or $E \geq 1.7 \times 10^{17}$ eV, and such particles will escape the galactic field regardless of the excitation of Alfvén waves.

(2) The cosmic rays travel on field lines as a front of width $(1 - \beta)R$ where $\beta = v/c$, or $(1 - \beta) = 1/(2\Gamma^2)$. Presumably this front must broaden until it has a width comparable to that necessary for Alfvén wave scattering $\cong 3\lambda_{\text{Alfvén}} \cong 30 R_L$. When the width of the front is equated with the scattering region thickness, one obtains $\Gamma > 50 R^{1/3}$ where R is the

radial distance traversed in pc. Since the adiabatic expansion loss is negligible for $R > 100$ pc where the subsequent expansion is small, then $\Gamma > 250$ and particles of energy greater than this will not suffer adiabatic deceleration. (The energy density of this front falls below $B^2/8\pi$ only after $R > 100$ pc.)

(3) The lower energy particles, $\Gamma < 250$, have time to be scattered back into the bubble so that one might expect a progressive deceleration to lower energy. Instead it is evident that for any mass fraction that is being scattered, that the immediately following matter of greater energy or momentum flux $\propto E^{-1.5}$, will excite greater Alfvén wave turbulence. This greater turbulence will have the effect of trapping the preceding higher energy fraction by the action of the following lower energy fraction. This will lead to compression of the preceding fraction between the following fraction within and the ISM. This compression of the original mass fraction between the ISM and the following mass fraction is similar to the plasma shock mechanism. Hence the higher energy matter will be heated by the lower energy matter and the subsequent expansion deceleration will be less important - if not reversed. It is clear that a complicated calculation needs to be performed to better understand the balance between compression and scattering heating versus deceleration. Adiabatic deceleration is too simplistic with a velocity distribution.

Orthogonal Shock

Finally we note that the initial expansion of the bubble is driven by matter whose velocity is extremely relativistic and will create a shock in this medium when the velocity vector is orthogonal to the magnetic field. We foresee that the ISM medium adjacent to the SN will then be loaded with

shocked particles whose energy is directly related to the driving particles. The shock wave created by the expanding matter is an inversion of the energy of the driving matter such that the interstellar medium will be loaded with particles whose energy reflects the energy of the particle group that made the shock in the first place. Hence, in this picture the relativistic ejected matter produces a shock across the field which in turn will reproduce the energy spectrum of the matter originally ejected from the supernova. The composition will now be the composition of the matter of the interstellar medium, the same as for the collisionless plasma shock that depends on Alfvén wave scattering. Here, however, the Alfvén wave scattering limit is far exceeded by the shock strength of the ejected matter and the particles attain the energy of the driving piston, that is, original supernova ejected matter on the first shock traversal.

Therefore, in general, we envisage that the bulk of the $B_{||}$ ejected-relativistic matter will enter the magnetic field and a significant fraction will escape without adiabatic deceleration. In addition, we see that a significant fraction of the work of adiabatic deceleration of the high energy matter will be exchanged with energy of the immediately adjacent interstellar medium and hence reproduce the cosmic ray spectrum in new matter. In addition there is the possibility that the cosmic rays that do escape immediate deceleration and that are trapped in the magnetic field of the interstellar medium will be further compressed and heated by the subsequent compression of the slower moving but far more massive and energetic matter following in the supernova ejected mass distribution.

These are first-order processes and do not depend upon the stochastic scattering and acceleration across the subsequent shock front in the interstellar medium. In this picture there is one monotonic source of

high energy particles from the supernova that creates the whole of the cosmic ray spectrum. I believe this is the most likely circumstance for cosmic ray acceleration because of the problem of serendipitous overlapping of several different sources. The highest energy source should usually dominate the particles of the lower energy source.

VII. Efficiency for Producing Cosmic Rays - Envelope versus ISM

The inefficiency of the envelope shock is the kinetic energy of the matter moving more slowly than $0.82c$, or $\Gamma = 2$, i.e., the definition of "cosmic ray" of kinetic energy 1 GeV per nucleon. The source of energy is the bulk mass motion ($v \cong 10^9 \text{ cm s}^{-1}$) of the type I supernova, $\cong 10^{51}$ erg and $.5$ to $1.4 M_{\odot}$ ejected. The energy spectrum of ejected matter between these two velocities is at first steeply rising $N(>E) \propto E^{+3}$ rapidly declining to $N(>E) \propto E^{-2.5}$, so that the nonrelativistic energy fraction is approximately $300 \times$ the relativistic energy or an average spectrum 1 MeV to 1 GeV of $N(>E) \propto E^{-1.83}$. This is then an efficiency of 0.3%. The ISM plasma shock on the other hand is driven by roughly the same total energy, 10^{51} erg, for presumably both type I and type II supernova roughly twice as frequently, but at a lower ISM shock velocity of $\gtrsim 3 \times 10^8 \text{ cm s}^{-1}$. Injection has not yet been fully settled for the ISM plasma shock. In order for the ISM plasma shock cosmic ray acceleration efficiency to be $1/300$ of the input energy, the nonrelativistic spectrum connecting the shock energy (100 keV) to 1 GeV cosmic rays (i.e., 10^4 in energy) would have to have a spectrum flatter than $N(>E) \propto E^{-1.62}$ and in addition experience no losses from heating the ISM. This seems most unlikely. Instead a reasonable acceleration spectrum for the nonrelativistic fraction would be $N(>E) \propto E^{-2.5}$ as shown in the calculations of Ellison,

Jones, and Eichler (1983). In addition $3/4$ of the internal energy of the shock would be left as kinetic energy of the ISM. The efficiency then becomes $(100 \text{ keV}/1 \text{ GeV})^{1.5}/4 = 2.5 \times 10^{-7}$. The suggestion of seeding the shock with CRs just shifts the burden to the origin of the seed. -Repeated ISM shock acceleration reaccelerates spallation products and then the resulting composition disagrees with measurements.

VIII. Conclusions

The hydrodynamic shock wave occurring during the explosion of a Type I supernova speeds up in the decreasing density of the stellar structure producing high energy particles whose spectrum and total energy reproduce the cosmic ray flux observed in the Galaxy. The composition of this matter should reflect the composition recently accreted from a companion star before the supernova explosion, but partially spalled in the shock and partially resynthesized in the postshock expansion. Some of these particles should exchange their energy with others in the interstellar medium and others escape to produce directly the cosmic ray spectrum. Alfvén wave scattering may lead to an increase in the ejected relativistic matter energy rather than an adiabatic deceleration. The efficiency of a supernova envelope shock far exceeds the acceleration efficiency of the ISM Alfvén shock if nonrelativistic particle spectra are similar to present calculations.

Acknowledgements

I am indebted to Albert Petschek for many discussions and collaboration as well as very many others, particularly Montgomery Johnson, Dick White and Chester McKee. This work was supported by the Department of Energy.

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