DIFFUSE GALACTIC ANNIHILATION RADIATION FROM SUPERNOVA NUCLEOSYNTHESIS

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1. Introduction. A primary source of nucleosynthesis in the Galaxy is type I supernovae. The observation of gamma-ray lines produced by the radioactive decay of unstable elements created in these explosive events would be a direct test of models of galactic nucleosynthesis (Clayton, Colgate, and Fishman 1969). However the primary unstable elements produced in these explosions, ^{56}Ni and ^{56}Co (e.g. Weaver, Axelrod, and Woosley 1980) are short lived with respective mean lives of 8.8 and 113 days. These mean lives are much less than the mean time between type I supernova explosions in the Galaxy, ~50 years (Tammann 1974). This difference in time scales suggests that the observation of such short-lived phenomena from galactic type I supernovae is unlikely in the near future. Clayton (1973) and Ramaty and Lingenfelter (1979) have considered, however, another potentially observable signature of supernova type I nucleosynthesis. They conjectured that, if a small, ~10%, fraction of the MeV positrons generated by the decay of ^{56}Co in remnant ejecta escape into the diffuse interstellar medium, where positron lifetimes are significantly longer than the time between supernova events, diffuse gamma ray radiation would be produced by the annihilation of positrons accumulated from a large number of supernova explosions. This mechanism should be the dominant source of diffuse annihilation radiation (Ramaty and Lingenfelter 1979). The measurement of such annihilation radiation would place critical constraints on the galactic production rate of iron.

The escape of MeV positrons from the ejecta of type I supernovae was first suggested by Colgate (1970). Subsequent analyses of light curves of extragalactic supernovae at times >40 days after their explosion support Colgate's hypothesis that a fraction of the positrons escape from the dense ejecta material (Arnett 1979; Colgate, Petschek, and Kriese 1980). In view of the fact that such ejecta is expanding at speeds in excess of 10^4 km/s (e.g. Weaver, Axelrod, and Woosley 1980), escape seems to require that the electrons, when uncoupled by Coulomb collisions, stream along field lines without any significant pitch-angle scattering. Such escape appears to be contrary to a basic principle of plasma physics that relativistic particles streaming faster than the Alfven wave speed excite hydromagnetic waves, which in turn efficiently scatter in pitch angle the relativistic particles, thus inhibiting streaming (e.g. Kulsrud and Zweibel 1975). As will be discussed Alfven speeds in the dense ejecta are low, significantly less than light speed.

2. Alfven-Wave Generation. The escape of relativistic cosmic rays from supernova explosions was investigated by Kulsrud and Zweibel (1975). They found that the escaping relativistic cosmic-ray nuclei excite hydromagnetic waves by resonant interactions. This study will be employed here to investigate the escape of relativistic electrons. Kulsrud and Zweibel (1975) calculated the growth rate, Γ, for generating hydromagnetic waves by particles of number density, n_{rel}, streaming at velocity, V_s,

$$\Gamma=4\Gamma_o K^2 V_a(V_s-V_a)/[\Gamma_o^2 + (\Sigma-2KV_a)^2] \quad (1)$$

$$\Gamma_o=C_1\omega(n_{nel}/n)f(>E), \quad \Sigma=C_2\omega(n_{rel}/n)f(>E)$$

where K is the wavenumber of the Alfven wave pulse excited by relativistic particle streaming; V_a is the Alfven velocity; C_1 and C_2 are constants of order unity; ω is the nonrelativistic gyrofrequency; n_{rel}/n is the ratio of the relativistic particle density to the ambient ion density; and f(>E) is the fraction of relativistic particles more energetic than the minimum energy, E, required to excite the instability. When plasmas are nearly fully ionized, as would be found in supernovae at times $<10^3$ days, the waves generated by streaming dissipate by nonlinear wave-wave interactions (e.g. Kulsrud 1982).

3. Conditions in Ejecta. Ejecta is decelerated by a reverse shock wave propagating inward from the contact discontinuity between the ejecta and the swept up interstellar gas; another shock wave propagates outward from the contact discontinuity heating the swept up interstellar gas (Kahn 1973). After the passage of a reverse shock the ejecta is subject to Rayleigh-Taylor instabilities, which can efficiently amplify ambient magnetic fields (Gull 1973). McKee (1974) has shown that at early times the fraction of the ejecta decelerated is $(M_{sw}/M_{ej})^{0.5}$, where M_{sw} is the mass of the swept up interstellar gas and M_{ej} is the mass of the ejecta. At early times when the escape of positrons is expected to occur the reverse shock has propagated only a short distance from the contact discontinuity and the great bulk of the ejecta has not yet decelerated. The results of McKee (1974) show that, at ^{56}Co lifetime, 113 days, the reverse shock has decelerated only 10^{-4} of the ejecta for a remnant expanding in the tenuous interstellar medium of 4×10^{-3} cm^{-3} when a maximum ejecta velocity of 8×10^4 km/s (e.g. Woosley, Weaver, and Tamm 1980) is employed. At these early times the bulk of the ejecta, that creates the positrons, is unaffected by interaction with the interstellar medium.

The calculations of Weaver, Axelrod, and Woosley (1980) and Woosley, Weaver, and Taam (1980) are employed here to model conditions in a supernova explosion produced by the detonation of a 0.5 M_o accreting C/O white dwarf. Table 1 shows at the ^{56}Co mean lifetime, $\sim10^7$ seconds, the densities, ρ; velocities, V_{ej}; and distances from center of explosion in the outer undisturbed ejecta, R; for three Lagrangian mass coordinates, 0.1, 0.01, and 0.0001 in units of the ratio of ejecta mass external to the reference position, M_{ex}, to total ejecta mass, M_{ej}. Conditions most favorable for the excitation of waves occur in the outer ejecta. (Note that the last Lagrangian mass coordinate refers to the position of the reverse shock.) Also tabulated in Table 1 are esimates for positron escape times by streaming, $T_{es}=(R_s - R)/c$, where R_s is the position of the reverse shock and c is the light speed.

The magnetic field strength at these reference positions is estimated from magnetic flux conservation employing an initial radius of 5.5×10^8 cm (Weaver, Axelrod, and Woosley 1980). The stellar precursor is assumed to have a magnetic strength of 10^7 gauss, consistent with models of convective field amplification in the carbon-burning core of the white dwarf progenitor (Levy and Rose 1974). Listed in Table 1 are the magnetic field strengths, B, and Alfven speeds, V_a, at the reference positions.

4. Model Calculations. Employing the parameters discussed above the growth rates for the excitation of hydromagnetic waves by relativistic particle streaming, equation (1), are calculated assuming that f(>E) is unity and n_{rel}/n is 0.1. The resulting growth rates, Γ, wavenumbers, K, and wave periods, T_w, for these hydromagnetic waves are listed in Table 1. As can be seen the time scales for wave excitation, $1/\Gamma$, are much longer than the age of the ejecta, $\sim 10^7$ seconds. Thus the streaming of positrons generates negligible fluxes of hydromagnetic waves. More significantly at this time, the gyroperiod of the escaping positrons, T_w, is very much greater than the escape time. To scatter a particle in pitch angle by resonance interactions requires a minimum time scale of a gyroperiod (e.g. Wentzel 1974). Employing longer times does not significantly change these results. At a fixed Lagrangian position in the undisturbed ejecta the wave growth rate scales as $1/t^{0.5}$ where t is the time. The ejecta is decelerated and mixed with the swept up gas long before the time scale for wave generation equals the remnant age.

5. Escape. The electrons stream along field lines in the outer ejecta until they reach the vicinity of the reverse shock. The gyroradii of the relativistic electrons are much greater than the thickness of the collisionless reverse shock front which is of the order of 10 ion inertial lengths, $\sim 10^8$ cm (McKee and Hollenbach 1980). Since the gyroradii of the energetic electrons are signifcantly greater than the thickness of the shock front the electrons pass freely through the shock (Bell, 1978).

Table 1

Ejecta Parameters as Function of Relative Lagrangian Coordinates, M_{ex}/M_{ej}

M_{ex}/M_{ej}	0.1	0.01	0.0001
ρ(gm/cm^3)	1.2×10^{-17}	1.9×10^{-19}	3.4×10^{-22}
V_{ej}(km/s)	2.2×10^{4}	3.1×10^{4}	5.5×10^{4}
R(cm)	2.1×10^{16}	3.0×10^{16}	5.4×10^{16}
T_{esc}(s)	1.1×10^{6}	8.0×10^{5}	0.0
B(gauss)	6.9×10^{-9}	3.4×10^{-9}	1.0×10^{-9}
V_a(cm/s)	0.6	2.2	1.6×10^{1}
K(1/cm)	7.5×10^{-12}	3.7×10^{-12}	1.14×10^{-12}
T_w(s)	1.5×10^{12}	7.7×10^{11}	3.4×10^{11}
Γ(1/s)	1.5×10^{-11}	3.0×10^{-10}	6.8×10^{-10}

6. Conclusion. The propagation of MeV positrons in the outer ejecta of type I supernovae was investigated. It was found that the positrons created at times $\sim 10^2$ days propagated along magnetic field lines in the outer ejecta without any appreciable pitch-angle scattering or excitation of hydromagnetic waves. The lack of significant pitch-angle scattering is well consistent with models of wave excitation and scattering by resonant interactions. This occurs because time periods to scatter the particles or to excite waves are significantly longer than escape times. Thus it is expected that, when positrons are not coupled to the ejecta by Coulomb collisions, they escape from the relatively cold, dense ejecta and reside predominantly in the tenuous, hotter, shock-heated

interstellar gas. In the tenuous shock-heated gas the positron lifetime against annihilation is much greater than lifetimes in the dense ejecta. Thus the production of steady-state diffuse annihilation radiation by some fraction of these escaped positrons seems probable.

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