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GAMMA-RAY BURSTS AND COSMIC RAYS FROM ACCRETION-INDUCED COLLAPSE[1]

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ABSTRACT

Accretion of matter onto the surface of a white dwarf in a binary system can push it over the Chandrasekhar mass limit and may cause it to collapse into a naked or nearly naked neutron star without detectable optical emission. Such an optically quiet neutron star birth should be accompanied by a neutrino burst which could be detected with underground neutrino detectors only if the collapse took place in our own Galaxy or in very close nearby galaxies. However, neutrino-antineutrino annihilation outside the neutron star into electron-positron pairs will produce a gamma-ray burst that can be observed out to distances of at least 300 Mpc, if the mass surrounding the newly formed neutron star is less than about $3 \times 10^{-4}\ M_\odot$. If the surrounding mass is between $\sim 3 \times 10^{-4}\ M_\odot$ and $\sim 0.1\ M_\odot$, it will be injected into the interstellar space with energy above 10 meV per nucleon. Such nonrelativistic nuclei can be further accelerated to cosmic-ray energies in the interstellar space before they slow down by collisions. Thus, accretion-induced collapse may be an important source of cosmic rays and of cosmological gamma-ray bursts. Conversely, the observed rate of gamma-ray bursts and cosmic-ray data can be used to limit the birthrate of naked, or nearly naked, neutron stars to less than one per 10^3 yr in galaxies similar to ours. This rate is too small to contribute significantly to the birthrate of pulsars, and it implies that it is very unlikely that a neutrino burst unaccompanied by optical emission will be detected in the near future by the underground neutrino detectors.

Subject headings: accretion, accretion disks — cosmic rays — elementary particles — gamma rays: bursts — stars: neutron

1. INTRODUCTION

When the mass of an accreting white dwarf in a binary system crosses the Chandrasekhar limit, the core may collapse into a neutron star (see, for instance, Taylor & Stinebring 1986; Baron et al. 1987; Mayle & Wilson 1988). The collapse occurs on a typical free-fall time scale of $t_c \sim 1/(G\bar{\rho})^{1/2} \sim 100$ ms, where $\bar{\rho}$ is the average core density. The gravitational energy which is released in the collapse is about $GM^2/R \sim 3 \times 10^{53}$ ergs, where $M \approx 1.4\ M_\odot$ and $R \approx 1.5 \times 10^6$ cm are the mass and radius, respectively, of the neutron star. Most of this energy is radiated away during a cooling time of $\Delta t \sim 10$ s, by emission of neutrinos of the three known flavors from its surface (Colgate & White 1966). The neutrino temperatures are ~ 5 MeV for electron neutrinos and ~ 8 MeV for muon and tau neutrinos (e.g., Wilson et al. 1986; Bruenn 1987). The neutrino cooling time and temperature can also be estimated directly from very general considerations (see, for instance, Dar 1987).

In SN II explosions, all the spectacular visual display originates from the mass ejected in the explosion. The scenario of the explosion is essentially the following: When the central core of the proto–neutron star reaches supranuclear density, the repulsive short-range nuclear forces stop the collapse. The core then bounces and drives a shock wave that climbs outside through the infalling layers. This shock, perhaps assisted by neutrino energy deposition behind it, is thought to reverse the infall and eject the mass outside the proton–neutron star. To achieve this, the shock must survive energy losses and be strong enough to reverse the infall velocity of the layers, overcome their gravitational binding and propel them to the observed SN II expansion velocities of more than 3000 km s^{-1} (for a typical 10 $M_\odot$ star, this amounts to a total kinetic energy of about 10^{51} ergs). However, strong shocks also cool quickly by neutrino emission and tend to stall (Colgate 1989). If the shock stalls, the stellar collapse produces a black hole rather than a neutron star and SN II explosion. Until recently numerical simulations of gravitational collapse of evolved cores have not yielded clear-cut evidence for the formation of a neutron star and for a supernova explosion. Whether the prompt shock (Brown et al. 1982) is capable of ejecting the material outside the core, or whether it stalls and is reenergized by the long-term emission of neutrinos from the protoneutron star (Bethe & Wilson 1985) has until recently been the major problem in the theory of SN II explosions (Colgate 1989). The prompt-shock mechanism worked only for artificially soft nuclear equations of state, while the delayed-shock mechanism produced much weaker explosions than observed. However, until recently the numerical calculations have neglected the long-term energy deposition beyond the neutrinosphere by the reactions $\nu_i + \bar{\nu}_i \to e^+ + e^-$, whose rate was estimated by Goodman, Dar, & Nussinov (1987) to exceed 10^{50} ergs s^{-1}. These reactions beyond the neutrinosphere have been included in the numerical simulations of SN II explosions by Bowers & Wilson (1982), by Wilson et al. (1986), and by Wilson & Mayle (1989a, b). They seem to produce a "hot bubble" with high entropy right above the neutrinosphere, which pushes out the ejected matter, and consistently generates SN II explosions (Colgate 1989).

Different numerical simulations of accretion-induced collapse (AIC) also yield contradicting results on whether there is mass ejection by the shock in AIC formation of neutron stars (see, for instance, Baron et al. 1987; Mayle & Wilson 1988).

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Therefore, it is not clear what are the direct observational signatures from the birth of a neutron star in AIC, and consequently what is the birthrate of neutron stars in AIC. In this paper we estimate the effects of neutrino-antineutrino annihilation beyond the neutrinosphere (Goodman et al. 1987; Berezinsky & Prilutsky 1987) in AIC neutron star formation. Following the numerical results of Baron et al. (1987) and Mayle & Wilson (1988), we consider two idealized scenarios: In the first scenario a hot neutron star cools by the long-term neutrino emission from its sharp edge into empty space (naked neutron star). In the second scenario the neutrino emission from the neutron star is into empty space surrounded by shock ejected material or material accreted from the companion star. We show that such a naked neutron star formation produces a gamma ray burst that could be seen from a Hubble distance (Ramaty et al. 1990). Neutrino-antineutrino annihalation behind the shock ejected material can still produce a gamma-ray burst that can be seen at least from a distance of 300 Mpc if the total overlaying mass is less than $\sim 5 \times 10^{-4}\ M_{\odot}$. Such masses will be accelerated and injected into the interstellar medium with relativistic velocities (see also Paczyński 1990; Shemi & Piran 1990). Larger masses will be injected into the interstellar medium with smaller velocities. Thus, AIC can be a novel source of cosmic rays and cosmological gamma-ray bursts. Moreover, cosmic rays and gamma-ray bursts can be used to derive limits on the birthrate of neutron stars in AIC.

Our paper is organized as follows: In § 2 we derive the neutrino Eddington luminosity for a star, and we find that it is larger than the neutrino luminosity from the birth of a neutron star in gravitational collapse by more than two orders of magnitude. In § 3 we discuss the formation and free expansion of the fireball produced by neutrino-antineutrino annihilation near the surface of a naked protonneutron star. In § 4 we discuss the acceleration of a single "atom" by the radiation emitted from the fireball. In § 5 we consider the acceleration of an overlaying shell of matter by the fireball. In § 6 we discuss the gamma-ray bursts expected from naked and nonnaked formation of neutron stars in AIC and their observational consequences. In § 7 we discuss the injection of cosmic rays into the interstellar medium by neutron star formation in AIC. Final conclusions are drawn in § 8.

2. THE NEUTRINO EDDINGTON LUMINOSITY

During stellar collapse leading to the formation of a neutron star, most of the gravitational binding energy is first converted into thermal energy and then released by the emission of ν_e, ν_μ, and ν_τ neutrinos. The duration of the neutrino burst, determined by the diffusion of the neutrinos to the surface of the neutron star, is ~ 10 s. The spectrum of the emitted neutrinos is approximately thermal (Fermi-Dirac) with $kT_\nu \approx 5$ MeV for electron neutrinos and $kT_\nu \approx 8$ MeV for muon and tau neutrinos (Wilson et al. 1986; Bruenn 1987; Dar 1987). The peak neutrino luminosity per flavor is $L_{\nu_i} \approx 10^{52}$ ergs s^{-1}. The above picture has been beautifully confirmed by the detailed observations of SN 1987A (see, for instance, Arnett et al. 1989 and references therein).

We first consider whether the direct interaction of the neutrinos with matter beyond the neutrinosphere can overcome its gravitational binding and eject it. Neutrinos from the hot protoneutron star transfer momentum to the surrounding matter via inelastic and elastic scattering off free nucleons, nuclei, and electrons. For the relevant neutrino energies, the neutrino-electron cross sections are very small compared with the neutrino-nucleon and the neutrino-nucleus cross sections and can be neglected. The neutrino-nucleus elastic cross section, $\sigma(\nu A \to \nu A) \approx G_F^2 N^2 E_\nu^2/4\pi$, where G_F is the Fermi weak interaction coupling constant and N is the number of neutrons in the target nucleus, is important as long as most of the matter nuclei have not dissociated into nucleons (Freedman 1974; Freedman, Schramm, & Tubbs 1977). But, even for nondissociated nuclei as heavy as iron nuclei, the averaged cross section for momentum transfer *per nucleon* in neutrino-nucleus coherent elastic scattering,

$$\bar{\sigma} \equiv \frac{1}{2A} \int \frac{d\sigma}{d \cos\theta} (1 - \cos\theta) d\cos\theta$$
$$\approx \frac{\sigma(\nu A \to \nu A)}{3A} 2.2 \times 10^{-44} \left(\frac{E_\nu}{\text{MeV}}\right)^2 \text{cm}^2 , \quad (1)$$

is much smaller than the averaged cross section for momentum transfer to free nucleons in electron neutrino-nucleon charge exchange reactions,

$$\bar{\sigma}(\bar{\nu}_e p \to e^+ n) \approx \bar{\sigma}(\nu_e n \to e^- p)$$
$$\approx 9.1 \times 10^{-44} (E_\nu/\text{MeV})^2 \text{ cm}^2 . \quad (2)$$

Balancing the gravitational force with the rate of momentum deposition by the neutrinos, we obtain that the neutrino luminosity required in order to eject a hot baryonic matter composed of free nucleons, electrons, positrons, and photons, is given by

$$L_E \approx \frac{4\pi G M m c}{\bar{\sigma}} , \quad (3)$$

where G is Newton's gravity constant, M is the mass of the neutron star, m is the nucleon mass, and $\bar{\sigma}$ is given by equation (2). Expression (3) is similar to the expression for the Eddington luminosity for photons, except that the Thomson cross section is replaced by an average neutrino-nucleon cross section. Similarly, the neutrino Eddington luminosity for ejecting matter that consists of nondissociated nuclei, electrons, positrons, and photons is given by equation (3) with $\bar{\sigma}$ given by equation (1). For a 1.4 $M_{\odot}$ neutron star radiating electron neutrinos of average energy $\bar{E}_{\nu_e} \approx \bar{E}_{\bar{\nu}_e} \approx 16$ MeV ($\bar{E}_\nu \approx 3.15T$ for a Fermi-Dirac distribution), equations (1) and (3) yield $L_E \approx 5 \times 10^{54}$ ergs s^{-1}. Such a luminosity exceeds the peak luminosity of electron neutrinos from the protoneutron star (Wilson et al. 1986; Bruenn 1987) by about a factor of 500. Therefore, the ejection of overlaying hot baryonic (dissociated nuclei) matter by direct neutrino momentum deposition is not possible. For nondissociated matter which consists, say, of iron nuclei, the Eddington luminosity for 8 MeV neutrinos obtained from equations (2) and (3) is $L_E \approx 10^{55}$ ergs s^{-1}. Such a neutrino luminosity exceeds the peak neutrino luminosity in all flavors from the proto–neutron star by about a factor of 300. Therefore, the ejection of nondissociated nuclear matter by direct neutrino momentum deposition is not possible either.

Moreover, since the neutrino temperatures are only a few MeV, matter near the neutrinosphere cannot be heated by neutrino energy deposition to temperatures high enough to overcome the gravitational binding (~ 100 MeV per proton near the surface of the neutron star).

On the other hand, the production of electron-positron pairs near the surface of the neutron star by neutrino-antineutrino annihilation can lead to the ejection of gravitationally bound matter. Approximately 0.3% of the neutrino luminosity is converted into electron-positron pairs beyond the neutrinosphere (Goodman et al. 1987). These pairs, and the gamma rays resulting from their annihilation, are coupled to overlaying matter via cross sections of the order of the Thomson cross section, $\sigma_T \approx 0.665 \times 10^{-24}$ cm^2, for which the Eddington luminosity is $\approx 10^{38}$ ergs s^{-1}. Since the luminosity in the pairs and gamma rays ($\approx 10^{50}$ ergs s^{-1}) exceeds this Eddington luminosity by 12 orders of magnitude, overlaying matter will be ejected. The observational consequences of the formation of the electron-positron pairs will depend on the amount of the overlaying mass and its initial distribution.

3. PAIR FORMATION AND FREE EXPANSION OF THE FIREBALL

Goodman et al. (1987) showed that the neutrino burst from the birth of a neutron star produces a large quantity of electron-positron pairs outside the neutron star through neutrino-antineutrino annihilation,

$$\nu_i + \bar{\nu}_i \rightarrow e^+ + e^- : \quad i = e, \mu, \tau . \tag{4}$$

The annihilation cross section, computed in the standard electroweak theory, is

$$\sigma(\nu_i \bar{\nu}_i \rightarrow e^+ e^-) \approx K_{\nu_i} G_F^2 s ; \quad i = e, \mu, \tau , \tag{5}$$

where s is the square of the total energy in the center of mass frame, $G_F^2 \approx 5 \times 10^{-44}$ cm^2 MeV^{-2} is the Fermi weak interaction coupling constant squared,

$$K_{\nu_e} = \frac{1 + 4 \sin^2 \theta_W + 8 \sin^4 \theta_W}{6\pi} , \tag{6}$$

and

$$K_{\nu_\mu} = K_{\nu_\tau} = \frac{1 - 4 \sin^2 \theta_W + 8 \sin^4 \theta_W}{6\pi} \tag{7}$$

are dimensionless constants, and $\sin^2 \theta_W \approx 0.23$ is the experimental value of the Weinberg angle (Aguilar-Bennitez et al. 1988). Expression (5) is valid for $4m_e^2 \ll s \ll M_Z^2$, where $M_Z \approx 91$ GeV is the mass of the Z boson.

The rate per unit volume of e^+e^- pair production, $\dot{n}_\pm$, by neutrino-antineutrino annihilation at a radius r beyond R, the radius of the neutron star, is given by

$$\dot{n}_\pm(r) = \Sigma_i \frac{K_{\nu_i} G_F^2 \langle E_{\nu_i}^2 \rangle \dot{N}_{\nu_i}^2}{12\pi^2 c R^4} (1 - x)^4 (x^2 + 4x + 5) , \tag{8}$$

where $\dot{N}_{\nu_i}$ is the total emission rate of neutrinos of flavor i, $\langle E_{\nu_i}^2 \rangle$ is their mean squared energy and $x \equiv (1 - R^2/r^2)^{1/2}$. Integrating over the volume exterior to the neutron star we find for the temperatures given above that $\dot{N}_\pm \sim 10^{54} e^+e^-$ pairs per second are produced exterior to the neutron star. The total energy deposition rate exterior to the neutron star is $\dot{Q} \sim 10^{50}$ ergs s^{-1}, where all neutrino flavors contribute roughly equally. Most of these pairs are produced within few tenths of R beyond the neutrinosphere at R.

Pair annihilation, Compton scattering, bremsstrahlung, and pair production subsequently produce an opaque $e^+e^-\gamma$ plasma. Because the time scales for these processes are much shorter than the duration of the neutrino emission, a quasi steady state is achieved. The plasma pressure is many orders of magnitude higher than the magnetic pressure near the neutron star (Dar & Ramaty 1990). Thus, the $e^+e^-\gamma$ plasma will expand freely with expansion velocity close to the speed of light as long as it does not encounter overlaying material. Therefore the initial pair density is given by $n_\pm \approx \dot{N}_\pm / 4\pi R^2 c \approx 2 \times 10^{30}$ cm^{-3}. For such a density the mean free path for both the pairs and the photons is $\sim 10^{-6}$ cm, about 12 orders of magnitude shorter than R, showing that the plasma is indeed highly opaque. At formation, the energies of the electrons and positrons are approximately the same as those of the neutrinos, $\bar{E}_\nu \approx 3.15 T_\nu$, ~ 16 MeV for electron neutrinos and ~ 25 MeV for muon and tau neutrinos. However, while the neutrinos near the neutrinosphere have approximately the densities and energy distributions of neutrino blackbody radiations, the density of the pairs is much lower because of the relative rarity of the reactions $\nu_i \bar{\nu}_i \rightarrow e^+e^-$. Hence, the system will immediately rearrange itself to a blackbody of lower temperature via particle-number changing reactions, such as $\gamma e \rightarrow 2\gamma e$, $e^+e^- \rightarrow 3\gamma$ and $e^+e^- \rightarrow ee\gamma$. Conservation of energy then implies that the temperature of the $e^+e^-\gamma$ plasma near the surface of the neutron star is given by

$$(1 + 2 \times \tfrac{7}{8}) a T^4 \approx n_\pm \bar{E}_\pm \approx 7.4 \times 10^{31} \text{ MeV cm}^{-3} , \tag{9}$$

i.e.,

$$T(R) \approx 1 \text{ MeV} , \tag{10}$$

where aT^4 is the energy density of a blackbody radiation and $\bar{E}_\pm \approx 44$ MeV is the initial average energy of the e^+e^- pairs.

Subsequent evolution of the system consists of expansion and cooling via photon number changing reactions. After a very short time compared with the duration of the neutrino burst, a quasi steady state is achieved where the rate of energy escape from the fireball equals the energy deposition rate in the fireball ($L \approx 10^{50}$ ergs s^{-1}). As long as the plasma is highly opaque, matter and radiation are in thermal equilibrium, and the outward flow of photons is described by the diffusion equation

$$L = -\frac{16\pi r^2 a c T^3}{3\sigma_T 2 n_\pm} \frac{dT}{dr} . \tag{11}$$

Therefore, as long as $n_\pm \sim T^3$ the radiative transport equation yields $T \propto 1/r$. Beyond the radius where the plasma becomes thin to photon number changing reactions, the total sum of electron-positron pairs and photons pairs is conserved. The radiation, however, is no longer in thermal equilibrium with the electron-positron pairs. The temperature of the plasma fluid *in the fluid rest frame* drops then with distance r from the neutron star as $T(r) \propto r^{-1}$. This drop reflects the ordering of the particle trajectories, which at a distance $r \gg R$ diverge by only $\theta \sim R/r$ relative to the local radial direction. Thus the energy available in the center of mass system in a $\gamma\gamma$ collision falls as $E\theta \sim ER/r \sim T(r)$. Consequently, the reverse pair creation reaction $\gamma\gamma \rightarrow e^+e^-$ is slowly phased out leading to the Boltzmann suppression of the number of pairs,

$$\frac{n_\pm}{n_\gamma} \approx 0.52 \left(\frac{m_e c^2}{T}\right)^{3/2} e^{-m_e c^2/T} . \tag{12}$$

The fireball becomes optically thin to photon production [approximate cross section $\sim (\alpha/\pi)\sigma_T$, where α is the fine-structure constant] when

$$\int_{r_0}^{\infty} 2n_\pm \left(\frac{\alpha}{\pi}\right) \sigma_T \, dr \approx 2 \times 10^{10} \left(\frac{T_0}{m_e c^2}\right)^{3/2} e^{-m_e c^2/T_0} \lesssim 1 . \tag{13}$$

This happens when $T_0 \equiv T(r_0) \approx 18$ keV and

$$r_0 \approx \frac{RT(R)}{T(r_0)} \approx 9 \times 10^7 \text{ cm} . \quad (14)$$

Beyond this radius the gamma rays do not suffer any energy losses. They have approximately the energy spectrum of a blackbody with temperature $T(r_0)$ which is blueshifted to much higher energies due to the radial outflow of the plasma fluid. Thus, the escape of radiation from the fireball into free space produces a gamma ray burst of $\sim 10^{51}$ ergs. Rather than attempting to calculate the energy spectrum of the gamma rays from the Doppler shift, we will estimate the approximate temperature of the gamma-ray burst from an energy conservation consideration. During the "steady state" emission, the energy flow is constant and approximately $L = 10^{50}$ ergs s^{-1} are radiated from the photosphere, whose radius is r_0, with a blackbody spectrum. Energy conservation and the Stefan-Boltzmann law imply that the temperature of the emitted blackbody radiation is $T_\gamma \approx (L/4\pi r_0^2 \sigma)^{1/4} \approx 140$ keV, where σ is the Stefan-Boltzmann constant. Because the fireball expands essentially with the speed of light, the duration of the burst, Δt_γ, is approximately the duration of the neutrino burst, i.e., $\Delta t_\gamma \approx \Delta t \approx$ a few seconds. Hence, a typical naked neutron star formation in AIC produces a $\sim 10^{51}$ ergs gamma-ray burst with a short rise time, a few seconds duration and a spectrum of a blackbody with temperature ~ 140 keV.

Let us next turn to the case where a small amount of stellar material is present beyond the neutrino-antineutrino annihilation region. Recalling that plasma neutrality implies equal accelerations of electrons and nuclei we next treat the acceleration of a "single electron proton system."

4. ACCELERATION OF AN ELECTRON-PROTON PAIR BY A GAMMA-RAY BURST

Let us consider the acceleration of an electron (that carries along with it a proton mass) by a gamma-ray burst of a constant luminosity, which during a finite time Δt_γ greatly exceeds the Eddington luminosity, $L \gg L_{\rm E}$, during a finite time Δt_γ. Let us assume also that the electron-proton system is positioned at an initial radius r_i close enough to the source such that in a very short time it is accelerated to a relativistic energy, $\gamma \equiv E/m_p c^2 \gg 1$. The electron-photon system will be accelerated mainly by Compton scattering of photons on the electron. Because of its high velocity, the electron sees the photons redshifted to energies $h_\nu \ll m_e c^2$. In the electron rest frame the photons scatter with a cross section given by the low-energy limit of the Klein-Nishina formula,

$$\frac{d\sigma}{d\Omega} \approx \frac{3\sigma_T}{16\pi} (1 + \cos^2 \theta) . \quad (15)$$

The scattering angle in the lab and in the electron rest frame are related via

$$\cos \theta_L = \frac{\sqrt{\gamma^2 - 1} - \gamma \cos \theta}{\gamma + \sqrt{\gamma^2 - 1} \cos \theta} . \quad (16)$$

Consequently, the average momentum transfer to the electrons in Compton scattering of photons of energy E_γ is

$$\langle \Delta p_e \rangle \approx p_\gamma \langle 1 - \cos \theta_L \rangle = \frac{\eta E_\gamma}{c\gamma^2} , \quad (17)$$

where

$$\eta \equiv 1 - \frac{3}{8}\left(\ln\frac{\gamma + \gamma_-}{\gamma - \gamma_-} + \frac{2}{3}\frac{\gamma}{\gamma_-} + \frac{2\gamma^3}{\gamma_-^3} - \frac{\gamma^4}{\gamma_-^4}\ln\frac{\gamma + \gamma_-}{\gamma - \gamma_-}\right)$$

and

$$\gamma_- \equiv \sqrt{\gamma^2 - 1} .$$

The function η changes very slowly between $\eta = 1$ for $\gamma = 1$ and $\eta \approx 4 \ln(2\gamma) - 5$ for $\gamma \gg 1$.

The energy increase of the already highly relativistic electron, due to Compton scattering of photons emitted between times t and $t + dt$, is therefore given by

$$dE = m_p c^2 \, d\gamma = \frac{\eta E_\gamma}{\gamma^2} \frac{L \, dt \sigma_{\rm T}}{E_\gamma 4\pi r^2} , \quad (18)$$

where the first term on the right-hand side is the average energy transferred to the electron per collision by a photon of energy E_γ, and the second term is the number of photons that collide with it at a distance r. For highly relativistic particles,

$$dt = \frac{dr}{v} - \frac{dr}{c} \approx \frac{dr}{2\gamma^2 c} . \quad (19)$$

Equation (18) can be integrated numerically with the aid of equation (19) from $r_i \ll 2ct\gamma^2$ to $r_f \to \infty$ to yield the maximum value of γ, which a single atom can be accelerated to by a gamma-ray burst. An approximate analytic expression is obtained if η is treated as a constant:

$$\gamma_{\rm max} \approx \left(\frac{Lt}{m_p c^2} \frac{5\eta\sigma_{\rm T}}{8\pi r_i ct}\right)^{1/5} . \quad (20)$$

Equations (18) and (20) can be generalized to the case of Z electrons which carry with them a nucleus of atomic weight A by replacing $\sigma_{\rm T}$ by $\bar{\sigma}_{\rm T} = \mu\sigma_{\rm T}$, where $\mu \equiv Z/A$.

5. SHELL ACCELERATION BY THE FIREBALL

Since the free-fall time of the outer layers of the collapsing star is longer than that of the core, the strong shock, which is formed at the early stage of the collapse and propagates outward with a speed close to the speed of light, may reverse their fall. An external shell well separated from the proto–neutron star can thus be ejected, by the shock itself (Mayle & Wilson 1988). Below we shall consider the subsequent acceleration of such a shell by the fireball.

Equation (20) can be used to estimate the maximum energy per nucleon that an optically thin ionized shell around the protoneutron star ($r_i \sim R \approx 1.5 \times 10^6$ cm) can be accelerated to by the gamma rays from the fireball ($L \sim 10^{50}$ ergs s^{-1}). For a hydrogen shell, the maximal energy is about 300 GeV per proton ($\gamma_{\rm max} \approx 330$), while for a shell that consists of heavier nuclei the maximal energy is about 270 GeV per nucleon ($\gamma_{\rm max} \approx 290$). A shell of density ρ and thickness ΔR is optically thin to Compton scattering if

$$\frac{\rho}{m_p} \mu\sigma_{\rm T} \Delta R < 1 . \quad (21)$$

Thus a fully ionized shell of total mass

$$\Delta M \equiv \epsilon \, M_\odot \quad (22)$$

is optically thin so long as it is placed at a radius $r > R_t$ where

$$R_t \approx \left(\frac{\sigma_{\rm T}}{4\pi} \frac{\mu\epsilon \, M_\odot}{m_p}\right)^{1/2} \equiv R_0\sqrt{\mu\epsilon} \approx 8 \times 10^{15}\sqrt{\mu\epsilon} \text{ cm} . \quad (23)$$

In particular a shell near the surface of the protoneutron star ($r \approx R \approx 1.5 \times 10^6$ cm) is optically thin only if

$$\mu\epsilon < 7 \times 10^{-20} . \quad (24)$$

In reality, the shell which is ejected by the prompt shock may be much more massive. During its acceleration it may remain optically thick or it may become optically thin before the end of the burst. Below we shall consider these various possibilities.

5.1. *Nonrelativistic Optically Thick Shell*

If the shell is massive enough, the fireball cannot accelerate it to relativistic energies. The total energy, Q, of the $e^+e^-\gamma$ plasma is expended as work done on the expanding shell. The maximum kinetic energy per nucleon follows from energy conservation:

$$\gamma - 1 \approx \frac{1}{\epsilon} \frac{Q}{M_\odot c^2} \ll 1 \quad \text{if} \quad \epsilon \gg \frac{Q}{M_\odot c^2} . \quad (25)$$

We note that the hydrodynamical calculations of Mayle & Wilson (1988) indicate that also the prompt shock imparts about 10^{51} ergs kinetic energy to the ejected shell. Consequently, for $\epsilon > 10^{-3}$, one may expect that the ejected material will be accelerated by the shock and/or the fireball to an energy per nucleon of about $(\gamma - 1)m_p c^2 \approx \text{MeV } \epsilon^{-1}$. For a typical ejected mass, $\Delta M \sim 0.04\ M_\odot$, found by Mayle & Wilson (1988), we expect an average kinetic energy per nucleon of about 25 MeV ($v \approx 70{,}000$ km s^{-1}.)

5.2. *Highly Relativistic Optically Thick Shell*

If the accelerated shell reaches high relativistic energies ($\gamma \gg 1$) then the average momentum transfer in a single Compton scattering becomes $\eta E_\gamma / c\gamma^2$. However, the deflected photons have an average scattering angle $\cos\theta_L \approx 1 - \eta/\gamma^2$ relative to the initial radial direction. This corresponds to a radial velocity $v \approx (1 - \eta/\gamma^2)c$ which is smaller than $v \approx (1 - 1/2\gamma^2)c$, the radial velocity of the shell. Consequently, after their first scattering, the photons moving with a radial velocity smaller than that of the shell will regain energy by collisions in the shell while those with a larger radial velocity lose energy. Therefore, the net result is that the photons retain their average energy but their angular distribution becomes isotropic in the shell rest frame while they are "locked" in the shell until it becomes optically thin. From momentum conservation it follows that

$$\frac{Q}{c} \langle \cos\theta_L \rangle + \epsilon M_\odot c\gamma_- = \frac{Q}{c} . \quad (26)$$

For an isotropic distribution in the shell rest frame, $\langle \cos\theta \rangle = 0$ and equation (16) yields $\cos\theta_L \approx 1 - 1/2\gamma^2$. Consequently,

$$\gamma^3 \approx \frac{1}{2\epsilon} \frac{Q}{M_\odot c^2} \quad \text{if} \quad \epsilon \gg 2 \times 10^{-6} . \quad (27)$$

The condition on ϵ was obtained as follows: If the luminosity of the fireball is constant during the burst, $dQ/dt = L = \text{const}$, equation (27) can be rewritten as

$$\gamma^3 \approx \frac{1}{2\epsilon} \frac{Lt}{M_\odot c^2} , \quad (28)$$

and together with equation (19) it can be used to solve for $r(t)$, the position of the shell when the photons that have been emitted at time t reach it:

$$t = \left(\frac{5r}{6c}\right)^{3/5} \left(\frac{2\epsilon M_\odot c^2}{L}\right)^{2/5} . \quad (29)$$

The minimum value of ϵ is obtained by requiring that the shell becomes transparent only when the last photons from the burst reach it, i.e., that $r = R_t = R_0(\epsilon)^{1/2}$ with R_0 given by equation (23), when $t = \Delta t$. For $L \approx 10^{50}$ ergs s^{-1} and $\Delta t \approx 10$ s this condition yields the value $\epsilon \approx 2 \times 10^{-6}$.

Note that most of the fireball energy remains in gamma rays and only a small fraction is used to accelerate the ejected shell. An observer at a very large distance will see the gamma rays that were last scattered in his direction (scattering angle $\cos\theta_L \sim 1 - 1/2\gamma^2$). Consequently, the duration of the gamma-ray burst will be

$$\Delta t_\gamma \approx \frac{R_t}{2c\gamma^2} \approx 1.24 \times 10^7 (\epsilon)^{7/6} \text{ s} . \quad (30)$$

For instance, if $2 \times 10^{-6} \lesssim \epsilon \lesssim 10^{-4}$, then the expected duration of the burst is $3 \lesssim \Delta t_\gamma \lesssim 230$ s.

5.3. *Shell Becomes Transparent during Burst*

For $\epsilon < 2 \times 10^{-6}$ the shell becomes transparent before the last photons from the burst reach it. Its energy keeps increasing only until it becomes transparent. The final value of γ which is reached can be obtained by substituting equation (29) with $r = R_t$ into equation (28):

$$\gamma \approx \left(\frac{5}{3} \frac{R_0 L}{M_\odot c^3}\right)^{1/5} (\epsilon)^{-1/10} ; \quad 2 \times 10^{-6} \gg \epsilon \gg 4 \times 10^{-20} . \quad (31)$$

In the limit $r_i = R = R_t$ (where R_t is given by equation [23]), when the shell is already transparent near the surface of the neutron star, equation (31) reduces to equation (20) if the average lab angle, $\langle \cos\theta_L \rangle \approx 1 - 1/2\gamma^2$, of the photons locked in the optically thick shell is replaced by the average scattering angle of the photons from a free electron, $\langle \cos\theta_L \rangle \approx 1 - \eta/\gamma^2$.

Note that in this case essentially all the fireball energy escapes as a gamma-ray burst with a duration of a few seconds and only a tiny fraction of the fireball energy is converted into kinetic energy of the ejected shell.

6. GAMMA-RAY BURSTS FROM ACCRETION-INDUCED COLLAPSE

In §§ 3 and 5, we have shown that both in the case of a naked neutron star and the case of ejected matter whose mass is less than $10^{-3}\ M_\odot$, most of the fireball energy, $\sim 10^{51}$ ergs, is emitted as a gamma-ray burst and only a small fraction is used in the second case to accelerate the ejected shell to relativistic energy. The duration of the gamma-ray burst, depends on the total ejected mass and ranges between a few seconds for total ejected mass below $10^{-5}\ M_\odot$ to a few thousand seconds for masses around $10^{-3}\ M_\odot$ with corresponding luminosities of $\sim 10^{50}$ ergs s^{-1} and $\sim 10^{48}$ ergs s^{-1}, respectively. For a detection threshold of about 10^{-7} ergs cm^{-2} s^{-1} (see, for instance, Higdon & Lingenfelter 1990), such a gamma-ray burst could be seen from cosmological distances. The detection rate of such cosmological gamma-ray bursts would be $(4\pi D^3/3)\Sigma_g f_g n_g$, where f_g is the birthrate of naked neutron stars per galaxy of type g and n_g is the number density of such galaxies in the universe. We assume that for all types of galaxies, including our own, f_g is proportional to L_g, the luminosity of the galaxy. The mean luminosity density of the universe is (Schechter 1976; Yahil, Sandage, & Tammann 1980; Felten 1985) $\approx 2.4 \times 10^8 h\ L_\odot\ \mathrm{Mpc}^{-3}$, where $\frac{1}{2} < h < 1$ is the Hubble constant in units of 100 km/Mpc · s. The luminosity of our Galaxy is $\approx 3 \times 10^{10}\ L_\odot$. Thus, we obtain that the observed rate of gamma ray bursts (see, for instance, Higdon & Lingenfelter 1990) of $\sim 100\ \mathrm{yr}^{-1}$, limits the birthrate of naked neutron stars to less than 1 in about 10^6 yr and of neutron stars with a $\Delta M \leq 10^{-3}\ M_\odot$ shell to less than 1 in about 10^3 yr. Redshift effects in a Friedmann-Lemaître universe that is constrained by observations can reduce the bound on the birthrate of naked neutron stars to about one in 10^5 yr.

Could any of the above gamma-ray bursts fit into the experimentally observed bursts?

The cosmological nature of the present source ensures the observed isotropy. Also, the suggested sources are nonrepeating, as are most of the gamma-ray bursts. The typical temperature, $T_\gamma \sim 140$ keV, of the gamma-ray bursts is also consistent with the observed average gamma-ray energies. The pulse duration Δt can be as short as the neutrino emission time $\Delta t \leq 10$ s for $\epsilon \leq 10^{-5}$ and can extend up to 10^3 s for $\epsilon < 10^{-4}$, in rough agreement with the observations. However, a substantial fraction, $\sim 30\%$, of the gamma-ray bursts have γ line features, most likely corresponding to cyclotron resonance absorption lines in a strong, $\sim 10^{12}$ G, magnetic field. Such fields are present only near the surface of neutron stars but not at distances where the fireball becomes transparent. Moreover, as we pointed out before, the fireball blows off the magnetic field of the neutron star.

Therefore, we do not expect gamma-ray bursts with cyclotron resonance absorption lines to be produced by neutron star formation in AIC.

While the present scenario is unlikely to explain all the bursts, it could account for a fraction of the gamma-ray bursts which are devoid of the γ-ray absorption lines. Note also that the $e^+e^-\gamma$ fireball sweeps away any matter orbiting near a neutron star formed in AIC. In particular, in scenarios where gamma-ray bursts are produced by comets impinging on neutron stars (Harwit & Salpeter 1973; Tremaine & Zytkov 1985), their mass must be about

$$M_m \approx 5 \times 10^{17}\ \mathrm{g} \tag{32}$$

in order to release $\sim 10^{38}$ ergs gravitational energy in a gamma-ray burst when they fall onto a neutron star. However, if such comets are present at a distance of the order of the distance to the Oort meteorite cloud in the solar system ($D \sim 10^3$ AU), they will be either evaporated by the gamma-ray burst or swept away by the shell blown off from the neutron star. This excludes scenarios where gamma-ray bursts are produced by comets impinging on neutron stars that were formed in AIC (Harding 1990).

7. COSMIC RAYS FROM ACCRETION-INDUCED COLLAPSE

The idea that collapses yielding neutron stars generate cosmic rays has been considered before (Colgate & White 1966). The specific suggestion that the shocks in Type II supernovae can achieve this has not been born out by recent observations of SN 1987A (see, for instance, Arnett et al. 1989 and references therein). However, the $e^+e^-\gamma$ fireball (and the shock itself) do accelerate the small amount of matter near the neutron star formed in AIC, injecting cosmic-ray particles into the Galaxy. In this section we would like to discuss the bounds which follow from the demand that we will not generate via this mechanism excessive amount of cosmic rays, and conversely, the possibility that this mechanism does indeed provide a main source of Galactic cosmic rays.

We note that for ϵ values in the range $0.001 < \epsilon < 0.1$ the bounds on the rate of neutron star formation in accretion-induced collapses from gamma-ray bursts deteriorate because of the spreading of the burst duration with increasing mass. In that case, independent bounds from cosmic-ray observations are therefore useful.

For $10^{-3} < \epsilon < 10^{-1}$ the $\epsilon 10^{57}$ ejected nucleons can be accelerated only to nonrelativistic energies. Thus we expect that a substantial portion of the fireball energy ($\sim 10^{51}$ ergs) is converted to the particle kinetic energy. The average energy per nucleon is then $\bar{E}/A \approx 10^{-3}\epsilon^{-1}$ GeV, i.e., between 1 GeV and 10 MeV in the ϵ range of interest.

The observed cosmic-ray intensity near Earth implies that the total power input into the Galaxy in particles of energies exceeding E is $3 \times 10^{40} E^{-0.2}$ ergs s^{-1}, where E is in GeV (Ginzburg & Syrovatskii 1964). Comparing this input with the above energy content in accelerated particles per collapse, we find a lower bound on the time interval between collapses in the Galaxy with overlaying masses less than $\epsilon\ M_\odot$ ($10^{-3} < \epsilon < 10^{-1}$),

$$\tau(\epsilon) \gtrsim 10^3 \epsilon^{-0.2}\ \mathrm{yr}\ . \tag{33}$$

For $\epsilon = 10^{-3}$, corresponding to an average particle energy $E \sim 1$ GeV, equation (33) yields $\tau \gtrsim 10^3$ yr. Thus, very reasonable accretion-induced collapse rates in our Galaxy (one collapse per 10^3 yr with $10^{-3}\ M_\odot$ mass ejection) can supply the required number of cosmic rays with the correct average energy. It would seem unreasonable that we have a spectrum of ϵ values which conspires to give also the correct cosmic-ray energy spectrum (at source). In particular, we have seen that our mechanism can accelerate protons only up to ~ 300 GeV whereas the cosmic-ray spectrum extends far beyond that. Let us assume then that the above mechanism serves to inject particles into the interstellar medium with initial energies between 10 MeV and 1 GeV corresponding to $10^{-3} < \epsilon < 10^{-1}$. The observed power-law energy spectrum is fixed by subsequent "conventional" magnetic field and shock wave interstellar acceleration mechanisms. A total column density of 0.1–10 g cm^{-3} of material is required to stop 10–100 MeV protons. Thus, even such protons will travel 10^{23}–10^{25} cm before stopping. Thus, most of these particles are likely to be accelerated in the interstellar medium and form the generic cosmic-ray population. The collapses with $10^{-2} < \epsilon < 10^{-3}$, corresponding to such average initial proton energies, release 10–100 times more particles than the one with $\epsilon = 10^{-3}$. Consequently, if a lower rate of 1 accretion induced collapse per 10^4–10^5 yr in our Galaxy, with $10^{-2} < \epsilon < 10^{-1}$, respectively, would suffice for the required particle injection. In all the above discussion, ϵ, the mass fraction ejected in the accretion induced collapse was viewed as a free unknown parameter. Numerical simulation of Mayle & Wilson (1988) do suggest an appreciable ($\sim 4\%$) mass ejection making the $10^{-2} < \epsilon < 10^{-1}$ range relevant.

To conclude this section, we would like to comment on the cosmic-ray abundances expected from cosmic-ray injection by accretion-induced collapse or bare collapse of white dwarfs. The accelerated material may contain (*a*) accretion disk material which was accreted from the envelope of the companion star; (*b*) surface material from the white dwarf itself that was ejected by the prompt shock; or (*c*) surface material from the white dwarf that underwent explosive nucleosynthesis by the prompt shock that ejected it. The injected cosmic rays can be a proper combination of these three distinct compositions. The observed cosmic-ray abundances follow roughly the cosmological abundances, and this could also be the case in our present model if most of the accelerated matter comes from the surface of a companion star. However, there are a few well-established "anomalies" in the cosmic-ray abundances, such as the near equality between the oxygen and carbon abundances, that could be explained if the ejected material came from the surface of an evolved bare O-Mg-Ne core. Reliable estimates of the expected cosmic-ray abundances from the birth of nearly naked neutron stars, however, require detailed nuclear synthesis calculation incorporated into detailed stellar evolution and stellar collapse codes, which are beyond the scope of this research.

8. CONCLUSIONS

In this paper we have shown that the birth of naked or nearly naked neutron stars in accretion-induced collapse or in bare collapse of white dwarfs can produce cosmological gamma-ray bursts and can provide the required injection rate of cosmic rays into the interstellar space. In fact, only a minute fraction of all stellar collapses need to be naked or nearly naked in order to explain gamma-ray bursts of possible cosmological origin or to account for the cosmic ray injection rate. Conversely, one can use our results to set limits on the rate of stellar collapses resulting in neutron stars with overlaying masses less than $\sim 10^{-3}\ M_{\odot}$. The observed rate of gamma-ray bursts and cosmic-ray data limit this rate to be less than 1 per 10^3 yr.

We estimate that most of the e^+e^- pairs annihilate in flight on a short time scale in the vicinity of the neutron star. But even if the majority of them escape to the Galactic interstellar medium where they annihilate on much longer time scales, then the birthrate of naked or nearly naked neutron stars is bounded by observations of diffuse 0.511 MeV Galactic annihilation radiation. These observations require (Lingenfelter & Ramaty 1989) positron injection rate of $\sim 2 \times 10^{43}$ positrons per second, and the bulk of these positrons are thought to result from the decay of ^{56}Co produced in SN I explosions. Therefore, the injection of $\sim 10^{55}$ e^+e^- pairs per naked neutron star imply that the Galactic birthrate of such stars must be less than about 1 every 10^4 yr. Thus the potential escape of the pairs from the neutron star cannot invalidate the limit set by the gamma-ray bursts.

The current estimate of the Galactic pulsar birthrate is about one every 30–120 yr with a best estimate of about one birth every ~ 56 yr (Narayan 1987). In comparison, the Galactic SN II explosion rate is estimated to be one SN II every 25–100 yr (this follows from the observed SN II rate in external galaxies (Evans, van den Bergh, & McClure 1989) of about $(1.04 \pm 0.30)h^2$ SN II per century per $L_B = 10^{10}\ L_{\odot}$). The large uncertainties in both of these rates would allow a significant contribution to the pulsars birthrate from naked or nearly naked neutron star formation. But, as we have shown, the gamma-ray bursts, the 0.511 MeV Galactic annihilation radiation, and the cosmic rays exclude this possibility. Moreover, the upper bound on the Galactic birthrate of naked or nearly naked neutron stars of less than 1 in 10^3 yr makes it very unlikely that a neutrino burst unaccompanied by optical emission from the birth of a naked or nearly naked neutron star will be detected in the near future by the underground neutrino telescopes.

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