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Pulsars Covered by the Dense Envelopes as High-Energy Neutrino Sources

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A newly-born pulsar is covered by the expanding dense envelope, and high-energy charged particles accelerated near the pulsar are effectively converted into high-energy neutrinos with energies of $10^{13}\sim10^{14}$ eV through meson production and its decay. Such a opaque phase continues for a couple of months. This model was first proposed by Berezinsky, but we will show that his flux estimation is too high by a factor of $\sim10^5$. This reduction is needed because the particle flux is limited by the accompanying phenomena such as the source power of high-energy cosmic rays and the light element production rate by spallation nuclear reaction.

§ 1. Introduction

The rotational energy of a neutron star is believed to be liberated by a high-energy particle flux[1] and/or a coherent strong wave due to a huge magnetic field.[2),3] A newly-born neutron star is covered by a dense envelope which was thrown away by supernova explosion, and the liberated energy flux interacts with this envelope matter, changing its energy into other modes of energy. The effects of this interaction in the early phase of supernova nebula have been discussed in relation to an explosion itself,[4] light curve[5),6] and element composition of cosmic rays.[7] Recently, in addition to them, Berezinsky proposed a possible source model of high-energy neutrinos.[8] In this model, neutrinos are produced through nuclear collisions of high-energy particle flux with the envelope matter; multiple meson production at the nuclear collision and decays of mesons into neutrinos.

The cosmic-ray neutrinos with energies as high as $10^{13\sim14}$ eV are interesting not only for astrophysics but also for particle physics: This flux is expected to exceed the atmospheric neutrino flux in such high-energy region and only the cosmic-ray neutrinos can provide such high energy as $\sim10^{14.5}$ eV where the unitarity limit of the cross section is reached.[9] Further, the experimental project named DUMAND (deep under water muon and neutrino detection) is now going on and their detection may be hopeful in the near future.[9]

The purpose of this paper is to elaborate Berezinsky's model and to discuss related astrophysical phenomena. In §§ 2 and 3, we discuss an absorption of particle flux in the expanding envelope and it will shown that the high-energy particle flux is effectively converted into neutrinos in early stages during a couple of months after the neutron star formation. In § 4, we review a particle acceleration mechanism in a pulsar and apply them to estimate an attainable energy of the neutrinos. In § 5, we discuss related astrophysical phenomena such as Galactic cosmic-ray origin, production of light elements by spallation of heavier elements and light curve of supernova, and, finally, estimation of the metagalactic background neutrino flux due to this type of sources is given.

§ 2. Interaction of particle flux with expanding envelope matter

Since particle acceleration is supposed to occur in the central compact region of the envelope, a high-energy particle will propagate to the radial direction. Assuming a homologous expansion of the envelope, we write a density distribution as

$$\rho(t, q) = \frac{3M_e g(q)}{4\pi R(t)^3}, \qquad (2\cdot 1)$$

where q is a comoving coordinate defined by $q=r/R(t)$, a distribution function $g(q)$ is normalized as $\int_0^1 g(q) q^2 dq = 1/3$, M_e and $R(t)$ are a mass and a radius of the envelope respectively.

Through the nuclear interaction of the cross section σ_n, the particle flux decreases in the envelope by the factor

$$I_p(t) = \exp\{-(R_c\xi/R(t))^2\}, \qquad (2\cdot 2)$$

where

$$R_c = \left(\frac{3M_e\sigma_n}{4\pi m_p}\right)^{1/2} = 10^{15.5} M_e^{1/2}\ \mathrm{cm} \qquad (2\cdot 3)$$

and

$$\xi \simeq \left(\int_0^1 g(q)\, dq\right)^{1/2} + O(\dot{R}/c), \qquad (2\cdot 4)$$

M_e being in the solar mass unit and σ_n being taken as 43 mb. Since a time delay effect of the order of $(\dot{R}/c)$ is neglected in our discussion, ξ is time-independent and solely determined by the density distribution. If we take such a distribution as $g(q) \propto e^{-aq}$, ξ becomes like $\xi \simeq 1$ for $a \ll 1$, i.e., an uniform distribution, and $\xi = a/\sqrt{3}$ for $a \gg 1$. In general, ξ takes a large value for centrally condensed type distributions.

To the neutrinos resulting from the mesons produced at the nuclear interaction the envelope matter is completely transparent and their flux is proportional to the function

$$I_\nu(t) = 1 - I_p(t). \qquad (2\cdot 5)$$

In the very dense envelope, however, the produced mesons are absorbed before they decay into neutrinos. If we include this effect, Eq. (2·5) is modified as

$$I_\nu(t) \simeq (1 - I_p(t)) \frac{R(t)^3}{R(t)^3 + R_\pi^3}, \qquad (2\cdot 6)$$

assuming $R_c \gg R_\pi$, where R_π is defined as

$$R_\pi = \left(\frac{3\sigma_\pi M_e}{4\pi m_p}\tau_\pi \gamma c\right)^{1/3} = 10^{11.2}\gamma^{1/3}M_e^{1/3}\ \mathrm{cm}\,, \tag{2·7}$$

where τ_π, γ, σ_π are the decay life, the Lorentz factor, the absorption cross section of $\pi^\pm$, respectively. The γ-ray flux resulting from the π^0 decay is proportional to the function

$$I_\gamma(t) = \frac{\exp\{-(\sigma_\gamma/\sigma_n)(R_c\xi/R(t))^2\} - \exp\{-(R_c\xi/R(t))^2\}}{(1-\sigma_\gamma/\sigma_n)}\,, \tag{2·8}$$

where σ_γ is the absorption cross section of γ-rays and of order of $\sigma_\gamma/\sigma_n \sim 0.2$. $I_\gamma(t)$ has the maximum at $R(t) = R_c\xi\{(1-\sigma_\gamma/\sigma_n)/\ln(\sigma_n/\sigma_\gamma)\}^{1/2}$ and both I_ν and I_γ decrease like $(R_c/R(t))^2$ in the later stage like $R(t) > R_c$. The ratio of the integrated flux of I_ν and I_γ is given as

$$\int I_\nu dR \Big/ \int I_\gamma dR = 1 + \sqrt{\sigma_\gamma/\sigma_n}\ .$$

§ 3. Neutrinos, γ-rays and electrons

Assuming that an injection rate of high-energy particles is proportional to the liberation rate $L(t)$ of the rotational energy, the energy fluxes of each component emitted from the envelope are given as

$$J_i(t) = f_i I_i(t) \lambda L(t) \tag{3·1}$$

for $i = \nu$, γ, p (particle flux), where f_ν and f_γ represent some numerical coefficients determined by a partition of energy to ν and γ, $f_p = 1$ and $\lambda L(t)$ represents the injection rate of the particle flux. In the following discussion of an order of magnitude, the differences among f_i are not important and we use the same coefficient A such as $A = f_i\lambda$.

The energy liberation rate is supposed to take a form

$$L(t) = \frac{L_0}{(1+t/\tau_b)^\beta} \tag{3·2}$$

and $\beta = 2$ for a brake due to magnetic dipole radiation. In the following, we will assume a constant expansion velocity like $R = Vt$. Then, combining Eqs. (3·1) and (3·2), time variations of the energy fluxes can be drawn schematically as shown in Fig. 1, which is characterized by three time scales;

$$t_\pi \equiv \frac{R_\pi}{V} = 0.2\left(\frac{\gamma}{10^6}\right)^{1/3}\frac{M_e^{1/3}}{V_9}\ \mathrm{days}\,,\quad t_c \equiv \frac{R_c\xi}{V} = 1.2\,\frac{M_e^{1/2}\xi}{V_9}\ \mathrm{months}$$

and the initial braking time τ_b of rotation, here $V_9 = V/10^9$ cm/sec. The total emitted energies $W_i = \int J_i dt$ of each component satisfy the following relations;

$$W_\gamma \sim \frac{AW_{\mathrm{rot}}(t_c/\tau_b)}{(1+t_c/\tau_b)^\beta} \tag{3·3}$$

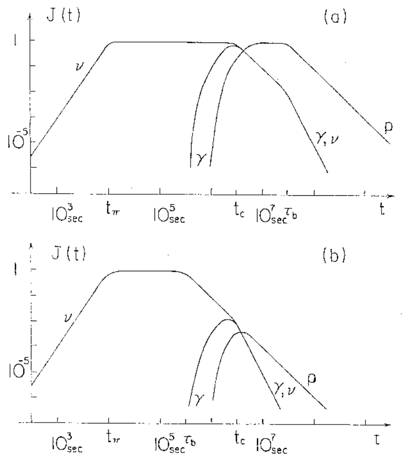

Fig. 1. Schematic feature of time variation of neutrino, γ-ray and high-energy particle fluxes $J_i(t)$ from a supernova envelope: Figure (a) is an example for $\tau_b \gg t_c$ ($\tau_b/t_c=10^2$) and Fig (b) is that for $\tau_b \ll t_c$ ($\tau_b/t_c=10^{-2}$). The ordinate is in an arbitrary unit.

both for $\tau_b<t_c$ and $t_c<\tau_b$, and

$$W_\nu \sim A W_{\rm rot}\,,\ W_p \sim W_\gamma \quad \text{if} \quad \tau_b<t_c\,, \tag{3·4}$$

and

$$W_\nu \sim A W_{\rm rot}\left(\frac{t_c}{\tau_b}\right),\ W_p \sim A W_{\rm rot} \quad \text{if} \quad \tau_b>t_c\,,$$

where $W_{\rm rot} \equiv L_0 \tau_b$. If $\tau_b \ll t_c$ as the case (b) of Fig. 1, the neutrino flux becomes much larger than γ-ray or particle fluxes, that is just the case considered by Berezinsky. In other high-energy neutrino production mechanisms such as collision between cosmic-ray proton and cosmic relict radiation,[10] the neutrino production is always accompanied by the γ-ray production of the same order. Therefore, in such models, the neutrino flux is rigorously limitted by the observation of γ-ray flux. In Berezinsky's model,[8] however, the production of neutrinos can exceed that of γ-rays by the factor of t_c/τ_b, supposing $t_c \gg \tau_b$.

The number of neutrinos caught by a detector is proportional to $W_\nu \sigma_0$ in the neutrino energy region where the cross section is given as $\sigma_\nu=\sigma_0 \varepsilon_\nu$ and $\sigma_0=0.8\times 10^{-38}\,{\rm cm}^2/{\rm GeV}$. It is worthwhile to mention that the detection rate is propotional only to the total energy of accelerated particles and independent of the particle energy and the details of meson production process.

The γ-rays from π^0 decay will give a start to electromagnetic cascade shower. The development of the shower is described by the age parameter s, which is approximately related to the matter thickness x in the radiation length as $s \simeq 3x/(x+2y)$, where $y \simeq \ln \varepsilon/\varepsilon_c$, ε and ε_c being the initial energy and the critical energy of the shower, respectively. For hydrogen gas, the radiation length is $X_0=62.8$ g/cm^2 and $\varepsilon_c=350$ MeV.[11] Substituting the thickness of envelope for x, we can know untill what stage the cascade shower has developed; the shower has already damped on the envelope surface if $s>1$ and it has not developed sufficiently even on the surface if $s<1$. The age parameter of the envelope shower at t is given as

$$s \simeq \frac{3}{1+2(R(t)/R_{s\,\max})^2} \tag{3·5}$$

and

$$R_{s\,\max}=\left(\frac{9}{7}\frac{\sigma_\gamma}{\sigma_n}\frac{1}{y}\right)^{1/2} R_c \simeq 0.27\left(\frac{3.5}{y}\right)^{1/2} R_c\,. \tag{3·6}$$

Then, the shower development is in the maximum stage on the surface for the shower at $t_{s\,\max}=R_{s\,\max}/V$; the maximum shower is ejected from the envelope a little earlier than the epoch of the γ-ray maximum. The total number of the electron-positron pairs injected into interstellar space from these showers is estimated as

$$N_e \simeq 0.03\,\frac{AW_{\rm rot}}{\varepsilon_c}\,\frac{t_{s\,\max}/\tau_b}{(1+t_{s\,\max}/\tau_b)^{\beta}}\,. \tag{3·7}$$

§ 4. Particle acceleration and neutrino energy

As the particle acceleration in a pulsar, the following two mechanisms have been proposed; (i) an acceleration by a static electric potential drop induced near the neutron-star's surface and/or in the corotating magnetosphere, and (ii) an acceleration by a coherent strong wave in the wave zone.

(i) *potential drop*

In the original model of Goldreich and Julian,[1] it is supposed that there is a steady current flow from the polar cap region along the open field lines and the particles of the current suffer an acceleration somewhere near the light cylinder by an amount of

$$\varepsilon_{\rm pol} \sim e\Delta V^{\rm pol} \sim \frac{e\Omega^2 R^3 B_s}{2c^2} = 10^{11.2}\Omega^2 B_{12}\ {\rm eV}\,, \tag{4·1}$$

where $\Delta V^{\rm pol}$ is a potential difference between the pole and the polar cap edge, Ω is in units of rad sec^{-1}, $R=10^6$ cm and $B_{12}=B_s/10^{12}$ gauss. Here, Ω, R and B_s are the angular frequency of rotation, the radius and the surface magnetic field of a neutron star, respectively.

Later on, more elaborate models which explain the emission mechanism of radio wave have been proposed. The potential drop is classified into two types; (a) a potential drop induced just above the surface of the polar cap, accompaning the space charge limited current flow, which was proposed by Sturrock,[12] and (b) a potential drop induced in the vacuum gap by which a steady current system is interupted. An existence of such gap was suggested by Holloway for the outer gap along the zero charge cone,[13] and by Ruderman and Sutherland for the polar gap just above the surface.[14] Holloway criticized the Sturrock model claiming that this contains an inconsistency.[15] However, we will adopt this model to estimate energy for comparison.

In case (a), the potential drop is given essentially by $\Delta V^{\rm pol}$ in Eq. (4·1). When an accelerated particle runs along the magnetic lines, we have taken into account a reaction of curvature radiation.[12),14),16] The radiation limited energy ε_{RL} is obtained generally from the relation

$$eE = \frac{2}{3}\,\frac{e^2}{c^4}\left(\frac{\varepsilon_{RL}}{mc^2}\right)^4\left(\frac{c^2}{\rho}\right)^2,$$

where E is electric field and ρ is a curvature of the magnetic lines. If this energy is smaller than $\varepsilon_{\rm pol}$, the maximum energy attainable is given by ε_{RL} rather than $\varepsilon_{\rm pol}$. Such cases occur for proton when

$$\Omega > 10^{2.8}\rho_6^{-4/13}B_{12}^{-6/13}\ \text{sec}\,, \tag{4·2}$$

where $\rho_6 = \rho/10^6$ cm, and

$$\varepsilon_{RL} = 10^{15.7}B_{12}^{1/4}\Omega^{3/8}\rho_6^{1/2}\ \text{eV}\,. \tag{4·3}$$

Next, in case (b), a growth and a breakdown of the strong electric field are occurring recurrently with a time-scale of micro-sec.[14] In the case of the polar gap, the maximum potential drop reached is given as

$$\Delta V^{\rm gap} \sim \frac{\Omega B_s}{c} h^2 \sim 10^{12.1}B_{12}^{-1/7}\Omega^{1/7}\rho_6^{4/7}\ \text{Volt}\,, \tag{4·4}$$

where h is a height of the gap and it is determined by the condition of sparkling due to e^+e^- pair creation.[14] $\Delta V^{\rm gap}$ is smaller than $\Delta V^{\rm pol}$. In the case of the outer gap, the potential drop becomes larger than $\Delta V^{\rm gap}$ in Eq. (4·4) as discussed in Ref. 17).

In the Ruderman and Sutherland model,[14] the ions are expected not to be stripped from the surface, because of the large binding energy per ion of the surface matter. However, in the early phase of pulsar, the surface temperature of neutron star is as high as 2×10^{7}°K,[18] and the thermaly emitted ions may provide an ion source, since the critical temperature for this is $T\sim1.2\times10^{7}$°K.[14]

(ii) *strong wave*

The acceleration by strong wave is characterized by a strength parameter defined by $\nu = e\sqrt{L/c}/(mc\Omega r)$. The maximum of ν is given by $\nu_c = e\sqrt{L/c}/mc^2 \simeq 10^{10.3}(L/10^{44}\ \text{erg/sec})^{1/2}$, since $r > c/\Omega$ in the wave zone. According to Gunn and Ostriker,[19] the phase locked acceleration is possible within the radius given by $r_{pL} \sim \nu_c^{2/3}(c/\Omega)$ and the Lorentz factor of a particle accelerated in this way is

$$\gamma_{pL} \sim \nu_c^{2/3}. \tag{4·5}$$

This energy is related to $\varepsilon_{\rm pol}$ as

$$\varepsilon_{pL} \sim (mc^2\varepsilon^2_{\rm pol})^{1/3}. \tag{4·6}$$

The effect of radiation reaction is negligible in these outer regions.

As the strong wave cannot propagate in the plasma, we have to assume a cavity of very thin density, which may be formed by pushing out the matter by the radiation pressure. Gaffet has assumed such cavity filled with the strong wave, in order to explain the light curve of supernova, in which the strong wave is assumed to be dissipated partially on the interface between the cavity and the dense envelope.[6]

Using the estimation of an acceleration energy given above, we can calculate

Table I. Examples of pulsar models and acceleration mechanisms. $\Omega(10^3\text{y})$ and $L(10^3\text{y})$ denote their values after 10^3 years. The energy ε is given in the unit of eV.

Model			A	B	C	D
Magnetic field B_s (gauss)			10^{12}	10^{12}	10^{13}	10^{13}
Initial brake time τ_b (years)			1	10^{-1}	10^{-2}	10^{-3}
Rotation energy $W_{\rm rot}$ (erg)			10^{52}	10^{53}	10^{52}	10^{53}
Rotation period $\log \Omega_0$(sec) ($\log \Omega(10^3y)$)			4.0 (2.7)	4.5 (2.5)	4.0 (1.5)	4.5 (1.5)
Liberation rate $\log L_0 \left(\frac{\rm erg}{\rm sec}\right)$ ($\log L(10^3y)$)			44.5 (38.5)	46.5 (38.5)	46.5 (36.5)	48.5 (36.5)
I	potential drop induced by current flow	$\log \varepsilon_\pi (\log \varepsilon_{RL})$	14.3 (17.2)	14.5 (17.4)	14.5 (17.4)	14.5 (17.6)
		$(\log \varepsilon_{\rm pol})$	(19.2)	(20.2)	(21.2)	(22.2)
	(case (a))	$\log A_{CR}^{\max}$	−6.4	−6.9	−5.4	−5.5
II	polar gap	$\log \varepsilon_\pi (\log \varepsilon_{\rm gap})$	11.0 (12.7)	11.1 (12.8)	11.2 (13.0)	11.3 (13.1)
	(case (b))	$\log A_{CR}^{\max}$	−3.7	−4.2	−2.9	−3.0
III	strong wave	$\log \varepsilon_\pi (\log \varepsilon_{pL})$	13.4 (16.0)	13.9 (16.7)	13.9 (16.7)	14.4 (17.4)
		$\log A_{CR}^{\max}$	−5.7	−6.5	−5.1	−5.5

the neutrino energy. In Table I, the characteristic physical quantities are shown, assuming typical values of surface magnetic field and an initial rotation frequency. We have chosen these values so that they give a large neutrino flux, and these values might be unreasonably large. $L(t)$ is calculated from

$$L \simeq \frac{B_s^2 R^6 \Omega^4}{c^3}, \qquad (4\cdot7)$$

which is correct both for the particle flux model and the magnetic radiation model. The meson energy is calculated from the acceleration energy of a charged particle using the relation of Fermi's fire-ball model for multiple meson production such as[11]

$$\varepsilon_\pi \sim 10^{8.2} (\varepsilon/m_p c^2)^{3/4}\ \text{eV}. \qquad (4\cdot8)$$

The neutrino energy is easily obtained from ε_π.

§ 5. Relations to other astrophysical phenomena

In this section, we discuss several astrophysical phenomena which are related to the high-energy neutrino source model in this paper. Through these observational facts, we will try to find the upper limit of the neutrino flux, which seems very much smaller compared with that of Berezinsky.[10]

(i) *high-energy cosmic rays*

We assume an injection spectrum to have a peak at ε_p and a width of the order ε_p. The high-energy particles escape into the intersteller space without interaction after t_c and their contribution to Galactic cosmic rays at the corresponding energy should be smaller than the observed flux;

$$\frac{W_p(\tau_{CR}/\tau_{sn})}{V_G}c<J_{CR}^{\rm obs}(>\varepsilon_p), \tag{5·1}$$

where τ_{CR}, τ_{sn}, V_G are the confinement time of cosmic rays, the supernova rate, the volume of Galaxy, respectively, and $J_{CR}^{\rm obs}(\varepsilon_p)$ is the integral energy flux of cosmic rays above ε_p. From Eqs. (5·1) and (3·4), we can get the maximum value for A as

$$A_{CR}^{\rm max}=10^{-3.6}\alpha_{CR}\left(\frac{10^{15}\ {\rm eV}}{\varepsilon_p}\right)^{0.6}\left(\frac{AW_{\rm rot}}{W_p}\right), \tag{5·2}$$

where

$$\alpha_{CR}=\frac{J_{CR}^{\rm obs}(>10^9\ {\rm eV})}{W_{\rm rot}(\tau_{CR}/\tau_{sn})c/V_G} \tag{5·3}$$

is a ratio of the total cosmic-ray energy flux $J_{CR}^{\rm obs}(>10^9\ {\rm eV})$ to a hypothetical flux which is obtained if all of $W_{\rm rot}$ is converted into cosmic rays. This α_{CR} is estimated as

$$\alpha_{CR}=10^{-1.5}\left(\frac{10^{52}\ {\rm erg}}{W_{\rm rot}}\right), \tag{5·4}$$

taking $\tau_{CR}=10^{15}$ sec, $\tau_{sn}=10^9$ sec, $V_G=10^{68}{\rm cm}^3$.

(ii) *high-energy γ-ray flux*

Since γ-rays are not confined in Galaxy, they contribute to the metagalactic background. A similar argument as above gives the upper limit of A as

$$A_\gamma^{\rm max}=10^{4.2}\left(\frac{J_\gamma(>\varepsilon_\gamma)}{J_{CR}(>\varepsilon_p)}\right)_{\rm obs}A_{CR}^{\rm max}, \tag{5·5}$$

where the numerical factor comes from the term $\tau_{CR}H_0/(4\pi\tau_{SN}V_G\nu_{\rm sneg})=10^{4.2}$, $\nu_{\rm sneg}=10^{-2}/(1Mpc)^3100{\rm y}$ being the rate of supernovae per volume in the metagalaxy and H_0 the Hubble constant. This relation tells us that $A_\gamma^{\rm max}$ is a stronger condition than $A_{CR}^{\rm max}$ if the γ-ray component of cosmic rays around $10^{13}\sim10^{15}$ eV is smaller than nucleon component by the factor more than $10^{4.2}$. However, the present knowledge about the extensive air shower tells us very few about this point.

(iii) *source of e^+e^- cosmic rays*

The number of cosmic-ray electrons with energy above 100 MeV is less than that of cosmic-ray protons by the factor of 10^{-2}. Combining this fact and Eq. (3·7), we have a further condition for A as

$$A_{\rm el}^{\max}=10^{-2.3}\alpha_{CR}\left(\frac{t_c}{t_b}\right) \qquad (5\cdot6)$$

if $\beta=2$, and $t_{s\,\max}\gg\tau_b$. Therefore $A_{CR}^{\max}$ gives a stronger condition than $A_{\rm el}^{\max}$ if $\varepsilon_p>10^{13}$ eV.

(iv) *production of light element*

Since the envelope matter has been bombarded by the particle flux, the abundance of light elements such as Li, Be, B is increased as the heavier elements fragment by nuclear spallation reaction. The cosmic abundance of light element is $[L]\sim10^{-8}$. The total amount of $[L]$ increased by this process in Galaxy is given as $\Delta[L]=(M_e/M_G)\times$(number of supernovae in Galaxy)$\times[L]_e$, where M_G is mass of Galaxy and $[L]_e$ is the enriched abundance in the envelope. Therefore, the condition $\Delta[L]<[L]$ implies that $[L]_e<10^{-6}$.

The total number of spallation reaction $\mathfrak{N}_s$ in the envelope is written as $\mathfrak{N}_s\sim SAW_{\rm rot}/\varepsilon_p$, where S represents an average number of nuclear reactions initiated by a particle with energy ε_p. Then, $[L]_e$ is given as

$$[L]_e\sim\frac{\mathfrak{N}_s m_p}{M_e}[H]_e\,, \qquad (5\cdot7)$$

where $[H]_e$ is the heavy element abundance in the envelope. Using these relations, we finally get the limit for A as

$$A_{[L]}^{\max}=\frac{10^4}{S}\left(\frac{10^{52}\,{\rm erg}}{W_{\rm rot}}\right)\left(\frac{\varepsilon_p}{10^{15}\,{\rm eV}}\right). \qquad (5\cdot8)$$

The estimation of S is very difficult but it may be much less than $\varepsilon_p/100$ MeV. Then, $A_{CR}^{\max}$ gives a stronger condition than $A_{[L]}^{\max}$ unless $\tau_b<10^{-2}t_c(10^{15}\,{\rm eV}/\varepsilon_p)^{0.6}$.

(v) *light curve*

If we adopt the model where the supernova light curve is explained by a dissipation of $L(t)$ in the dense envelope, we get a relation between the maximum luminosity of supernova and $L(t)$ in the early phase which we are considering. A rough feature of this luminosity is given as follows; the dissipated thermal energy flows out by diffusion after t_d

$$t_d\simeq\left(\frac{\kappa M_e}{cV}\right)^{1/2}\sim10^{6.4}\left(\frac{M_e}{V_9}\right)^{1/2}\,{\rm sec}\,, \qquad (5\cdot9)$$

where the electron scattering opacity is taken for κ. Before t_d, the time scale of adiabatic cooling is faster than that of heat conduction. Writing the energy supply into heat as $BL(t)$, the maximum luminosity and the temperature are roughly estimated as

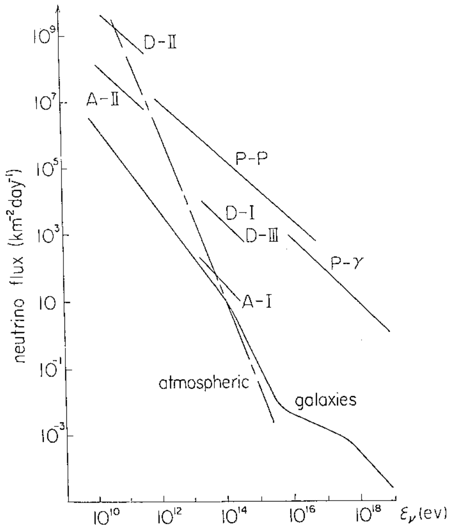

Fig. 2. Number flux of the metagalactic neutrinos. A-I, D-I,⋯ represent the models given in the Table; the pulsar model is specified by A~D and the acceleration model is done by I~III. Following Berezinsky and Smirnov,[10] other estimations of the flux are also shown; "atmospheric" neutrinos, the neutrino flux from the normal "galaxies" and the upper limits from the γ- and X-ray backgrounds, which are denoted by p-p and p-γ. p-p means "proton-proton" collision in the metagalactic space and p-γ does "proton-cosmic relict radiation" collision.

$$L_{SN}^{\max}\sim B\frac{c}{V}L(t_d)\sim 10^{14}B\left(\frac{L(t_d)}{10^{46}\ \mathrm{erg/sec}}\right)\frac{L_\odot}{V_9} \qquad (5\cdot 10)$$

and

$$T_{SN}\sim\left(\frac{3B}{4\pi}\frac{L(t_d)}{aV^3t_d^2}\right)^{1/4}\sim 10^{4.3}B^{1/4}\left(\frac{L(t_d)}{10^{46}\ \mathrm{erg/sec}}\right)^{1/4}\frac{M_e^{1/4}}{V_9^{1/2}}\ {}^\circ\mathrm{K}\,, \qquad (5\cdot 11)$$

The time scale of the luminosity maximum may be of order of $(V/c)t_d$. The conversion rate into heat, B, should satisfy $B>A$, because an appreciable fraction of the high-energy particle flux is dissipated into heat by the nuclear collision considered in this paper. Beside this dissipation mechanism via the high-energy particle flux, other mechanisms such as a dissipation via more low-energy particles may be also effective. The situation where both A and B are much less than unity implies that most of the liberated energy in the early phase is consumed only to accelerate the envelope expansion.

The rough estimate of $L_{SN}^{\max}$ in Eq. (5·10) seems too high to explain the observation, if $L(t)$ is taken as large as 10^{46} erg/sec and $B\sim 1$.

(vi) *metagalactic high-energy neutrinos due to this model*

The estimation of the background neutrino flux emitted in this way is given as

$$10^{2.8}\left(\frac{10^{14}\ \mathrm{eV}}{\varepsilon_\nu}\right)\left(\frac{A}{10^{-5}}\right)\left(\frac{W_{\mathrm{rot}}}{10^{52}\ \mathrm{erg}}\right)\mathrm{km}^{-2}\ \mathrm{day}^{-1}, \qquad (5\cdot 12)$$

which is shown in Fig. 2 for several choices of the parameters. Though this estimation contains some unfixed parameters, it is likely that this flux exceeds the atmospheric neutrino flux for $\varepsilon_\nu>10^{14}$ eV. By the typical DUMAND detector which contains 6×10^{38} atoms,[9] we may be able to catch them by the rate

$$10^{0.1}\left(\frac{A}{10^{-5}}\right)\left(\frac{W_{\rm rot}}{10^{52}\,{\rm erg}}\right){\rm year}^{-1}, \tag{5·13}$$

which is too small for practical observation. Our estimation of the detection rate is different from the Berezinsky's one by a factor of A, even if we take the same value for $W_{\rm rot}$. According to the arguments given in this section A is restricted less than $\sim10^{-5}$ for the neutrinos of $10^{13\sim14}$ eV.

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