GeV gamma-rays and TeV neutrinos from very massive compact binary systems: The case of WR 20a

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ABSTRACT
Massive Wolf-Rayet stars in a compact binary system are characterized by very strong winds which collide creating a shock wave. If the wind nuclei, from helium up to oxygen, accelerated at the shock can reach large enough energies, they suffer disintegration in collisions with soft thermal radiation from the massive stars injecting relativistic protons and neutrons. Protons collide with the matter of the wind and a fraction of neutrons collide with the massive stars producing γ-rays and neutrinos via decay of pions. We calculate γ-ray fluxes from inverse Compton e± pair cascades initiated by primary γ-rays and leptons, which are produced by protons, and the neutrino fluxes, produced by protons and neutrons, for the example compact massive binary WR 20a. From normalization of these cascade γ-ray spectra to the fluxes of the EGRET sources, 2EG J1021-5835 and 2EG J1049-5847 observed in the direction of WR 20a, we conclude that this massive binary can be detected by the IceCube type neutrino detector. The estimated muon neutrino event rate inside the 1 km² detector during 1 yr is between a few up to a few tens, depending on the spectrum of accelerated nuclei.

Key words: binaries: close - stars: WR 20a - radiation mechanisms: non-thermal - gamma-rays: neutrinos: theory - sources: 2EG J1021-5835; 2EG J1049-5847

The observation of neutrinos from discrete sources should give a unique information on the acceleration of hadrons in our Galaxy and on the origin of cosmic rays at energies above ~ 10^{12} eV. Possible production of neutrinos by relativistic hadrons accelerated inside the binary systems has been discussed in the case of accreting neutron stars (Gaisser & Stanev 1985, Kolb, Turner & Walker 1985, Berezinsky, Castagnoli & Galeotti 1985, Anchordoqui et al. 2003), energetic pulsars (Harding & Gaisser 1990), and accreting solar mass black holes called microquasars (Levinson & Waxman 2001, Distefano et al. 2002, Romero et al. 2003). For recent review of these models see e.g. Bednarek, Burgio & Montaruli (2005). Most of these models predict neutrino fluxes which should be easily detectable by the 1 km² neutrino detector of the IceCube type and in some cases by detectors of the present AMANDA II type. Recent searches of the discrete neutrino sources by the AMANDA II detector, which base on the analysis of the three-year experimental data (Ahrens et al. 2004, Ackermann et al. 2005), show no clear evidence of a significant neutrino flux excess above the background.

In this paper we show that also compact binary systems containing two massive stars can be interesting sources of TeV neutrinos likely to be observed by large scale neutrino telescopes. In fact, between 227 Wolf-Rayet (WR) stars (van der Hucht 2001), 39% form binary systems (including probable binaries), and about 15% of them form compact massive binaries with orbital periods < 10 days (see Table 15 in van der Hucht 2001). One extreme example of a binary system, WR 20a, contains two WR stars with masses ~ 70 M⊙, surface temperature $T = 10^7 T_5$, K ≈ 4 × 10^4 K, and radii ~ 20 R⊙ (where $R_\odot = 7 × 10^{10}$ cm is the radius of the Sun), on an orbit with the semimajor axis $a/2 \sim 26 R_\odot$ (Rauw et al. 2004, Bonanos et al. 2004). Let’s consider such a very massive binary composed of two massive WR type stars characterized by very strong winds, with mass loss rate $M = 10^{-5} M_{\odot}$ M⊙ yr⁻¹ ~ (0.8—8) x 10⁻⁵ M⊙ yr⁻¹, wind terminal velocities $v_\infty = 3 \times 10^3 v_\infty$ km s⁻¹ ~ (1—5) x 10⁶ km s⁻¹. The radiatively driven stellar winds accelerate from the surface of the star according to $v_\infty(t) = v_\infty + (v_\infty^0 - v_\infty^0)(1 - 1/r)^\beta \approx v_\infty^0 (1 - 1/r)$, where $v_\infty(t)$ is the wind velocity at the distance, $R = r R_{\text{WR}}$, from the star, $v_\infty^0 \approx 0.01 v_\infty^0$, $\beta \approx 1$ (e.g. Castor & Lamers 1979). The surface magnetic fields of WR stars can grow up to $B_\odot \sim 10^{15}$ Gs (Maheswaran & Cassinelli 1994). For the massive stars in WR 20a binary system we apply the typical values of $B = 10^9 B_3$ G, i.e. significantly below $10^9$ G. Such stellar winds collide creating a double shock structure separated by the contact disconti-
nuity which distance from the stars depends on their wind parameters.

It is likely that the binary system WR 20a is responsible for one of the EGRET sources observed in this region, i.e. 2EG J1021-5835 and GeV J1025-5809 or 2EG J1049-5847 and GeV J1047-5840 (see Thompson et al. 1995, Lamb & MaComb 1997). They have flat spectra with the index close to 2 (Merck et al. 1996) and fluxes at the level of \( \sim 10^{-7} \text{ cm}^{-2} \text{ s}^{-1} \) above 1 GeV (Lamb & MaComb 1997).

Nuclei present in the stellar winds can be accelerated at such shock structure either due to the magnetic field reconnection (we call this possibility as model I) or the diffusive field reconnection mechanism (model II) (see Fig. 1). The acceleration occurs in the fast reconnection mode provided that following condition is fulfilled, 
\[
\beta \equiv \frac{2m_{\text{p}}/Am_{\text{n}}}{\rho_{\text{w}}v_{\text{w}}^{2}} \approx 10^{-3}/A < 8\pi\rho_{\text{w}}B_{\text{sh}}/B_{\text{sh}}^{2} < 3.3 \times 10^{-3}M_{-5}/B_{12}^{3}r^{3}v_{3}(r) \]  
(although the fast reconnection can also occur in specific situations when this condition is not met, e.g. Hanasz & Lesch 2003), assuming typical temperature of the plasma at the shock \( T_{\text{sh}} = 10^{6} \text{ K} \) (e.g. Stephens et al. 1992), density of the wind \( \rho = M/4\pi R_{\text{w}}^{2}v_{\text{w}} \approx 5.4 \times 10^{-10}M_{-5}/r^{2}v_{3}(r) \) cm\(^{-3}\). Then, the reconnection occurs with the speed close to the Alfvén velocity \( v_{A} = B_{\text{sh}}/\sqrt{4\pi\rho_{\text{w}}\beta} \approx 7 \times 10^{8}B_{12}v_{3}^{1/2}(r)/rM_{12}^{3/2} \text{ cm s}^{-1} \). Nuclei can reach maximum Lorentz factors of the order of 
\[
\gamma_{\text{max}} \approx \frac{Z\epsilon B_{\text{sh}}L_{\text{w}}v_{\text{w}}}{Am_{\text{p}}v_{A}} \approx 2.6 \times 10^{6}B_{12}L_{12}^{1/2}v_{3}(r)/r^{3}M_{-5}^{2}, \tag{1}
\]
where \( A/Z = 2, \ m_{\text{p}} \approx m_{\text{n}} \) is the nucleon mass, \( L_{\text{w}} \) is the length of the reconnection region, \( c \) is the velocity of light, and \( \epsilon \) is the proton charge. In order to estimate the magnetic field strength after the shock, we apply the model for external magnetic field structure in the case of strong outflowing gas (Usov & Melrose 1992). According to this model the magnetic field is a dipole type below the Alfvén radius, \( R_{A} \), becomes radial above \( R_{A} \), and at the largest distances, determined by the rotation velocity of the star, i.e. at \( R \approx 10R_{\text{WR}} \), the toroidal component dominates. We are interested mainly in the region in which the dipole and radial components dominate. Above \( \sim 10R_{\text{WR}} \) the product of the magnetic field strength at the shock and the length of the reconnection region is roughly constant and the maximum energy of accelerated nuclei does not depend on the distance from the star. The Alfvén radius for the example parameters considered in this paper, \( v_{3} \approx 1, \ B_{3} = 1, \) and \( M_{-5} = 3 \), is at \( R_{A} \approx 1.3R_{\text{WR}} \). Therefore, we estimate the magnetic field at the shock in the region of radial magnetic field from
\[
B_{\text{sh}}(r) \approx B_{\text{sh}}(R_{\text{WR}}/R_{A})^{3}(R_{A}/R_{\text{sh}})^{2} \approx 750B_{3}/r^{2} \text{ Gs}, \tag{2}
\]
where \( R_{\text{sh}} \) is the distance from the center of the star to the acceleration region at the shock, i.e. \( R_{\text{sh}} = rR_{\text{WR}} \).

Nuclei can be also accelerated diffusively by the I order Fermi shock acceleration mechanism. In this case they obtain a power law spectrum. Nuclei are accelerated to \( \gamma_{A} \) during the time estimated by \( \tau_{\text{acc}} \approx (R_{A}/c)(c/v_{A})^{2} \), where \( R_{L} = Am_{\gamma}/(c\Delta ZB_{\text{sh}}) \approx 8 \times 10^{-6}v_{3}^{2}/B_{3} \) cm is the Larmor radius required to complete acceleration to \( \gamma_{A} = 10^{6}\gamma_{0} \). The maximum energy of nuclei is not limited by any radiation loss process but by their escape from the acceleration region. We identify this escape mechanism with the convection of relativistic nuclei with the wind flow along the surface of the shock. The escape time is estimated as
\[
\tau_{\text{esc}} \approx 3\tau_{\text{sh}}/v_{\text{w}} = 10^{4}R_{2}v_{3}(r) s, \tag{3}
\]
where the radius of the WR star, \( R_{\text{WR}} = 10^{12}R_{12} \) cm. The maximum energies of nuclei determined by the condition \( \tau_{\gamma} = \tau_{\text{acc}} \). Therefore, we expect the cut-off in the power law spectrum of nuclei at energies corresponding to
\[
\gamma_{\text{max}} = 3\epsilon_{\text{B}}B_{\text{sh}}R_{\text{sh}}v_{\text{w}} \approx 5.3 \times 10^{3}v_{3}(r)B_{\text{sh}}(r). \tag{4}
\]

Most efficient disintegrations occur when photon energies correspond to the maximum in the photo-disintegration cross-section of nuclei, which is at \( \sim 20 – 30 \text{ MeV} \) with the half width of \( \sim 10 \text{ MeV} \) for nuclei with the mass numbers between helium and oxygen (main composition of the surface of the WR stars) . Let us estimate the efficiency of the photo-disintegration process of nuclei in the considered here scenario. The average density of thermal photons from both stars (with similar temperatures), at the location of the shock, can be approximated by
\[
n_{\text{ph}} \approx \frac{4\sigma_{\text{SB}}T_{3}^{3}}{3c\kappa_{B}} \left( \frac{R_{\text{WR}}}{R_{\text{sh}}} \right)^{2} \approx 2 \times 10^{16}T_{3}^{3}r_{2}^{-3} \text{ ph. cm}^{-3}. \tag{5}
\]

![Figure 1](image-url)
mean free path for dissociation of a single nucleon from a nucleus is $A_{\Lambda} = (\sigma_{\gamma N}/\sigma_{\gamma A})^{-1} \approx 3.5 \times 10^{18} \text{cm}^2$, where for the photo-disintegration cross section the peak in the giant resonance is applied, $\sigma_{\gamma N} = 1.45 \times 10^{-22} \text{A}^2$ cm$^{-2}$ (Karakula & Tkaczyk 1993). For the considered here binary system WR 20a, for which $R_{gb} = a/2$ and so $r = 1.3$, the characteristic photo-disintegration time scale of nuclei with the mass number $A$, $\tau_{\Lambda} = A_{\Lambda}/c \approx 30/A \text{ s}$, is much smaller than the acceleration time and the convection escape time of nuclei from the acceleration region at the shock (estimated above). Nuclei with the initial mass numbers between 4 (helium) and 16 (oxygen), accelerated to energies given by Eqs. 1 and 3, should suffer complete fragmentation. Note that nuclei, and protons from their disintegration, can not escape far away from the shock region. For typical values of the binary system WR 20a, their Larmor radii, $R_L$, are significantly shorter than the characteristic distance scale of the considered picture defined by the semimajor axis of the binary system. Significant fraction of neutrons from disintegration of nuclei move toward the surface of the massive stars since the probability of dissociation of a single nucleon is the highest for the head on collisions of nuclei with thermal photons. These neutrons propagate along the straight lines and interact with the matter of stellar atmospheres. On the other hand, protons from disintegration of nuclei are convected outside the binary system along the surface of the shock structure in a relatively dense stellar winds.

The relativistic nuclei accelerated at the shock take a fraction, $\xi$, of the kinetic power of the two stellar winds,

$$P_\Lambda = \xi M v_\Lambda^2 \approx 6 \times 10^{37} \xi M_{-5} v_\Lambda^2(r) \text{ erg s}^{-1}. \quad (6)$$

The parameter $\xi$ is determined by the efficiency of acceleration of nuclei and by the solid angle $\Delta \Omega$ subtended by the active part of the shock (see Fig. 1). The active part of the shock is determined by the power of the stellar wind which falls onto the shock region in which nuclei can be accelerated above $\gamma_{\text{min}}$ (see Eq. 4). $\Delta \Omega$ can be estimated from the condition of acceleration of nuclei above $\gamma_{\text{min}}$ by applying Eqs. 1 (for model I) or 3 (for model II) and Eq. 2. By comparing Eq. 4 with Eq. 6 (for $\theta = 0^\circ$), we estimate the maximum distance of the active shock from the star on $R_{\text{max}} \approx 4(B_7^2 T_{9} L_{12} v_\Lambda^2/\gamma_2^2(r))^{1/3} R_{\text{WR}}$. Similar procedure, comparison of Eq. 6 with Eq. 4 allows us to estimate $R_{\text{max}}$ in model II, $R_{\text{max}} \approx 10^3 (R_7^2 B_7 T_{9} v_\Lambda^2(r))^{1/3} R_{\text{WR}}$. A part of the shock solid angle at which nuclei are accelerated above $\gamma_{\text{min}}$ is then $\Delta \Omega = 2\pi(1 - R_{\text{a}}/R_{\text{max}})/4\pi$, where the distance of the shock in the plane of the binary system is approximated by $R_{\text{a}} \approx a/2$. For the considered above parameters of WR 20a, $B_3 = 1$, $T_3 = 0.4$, $v_3^2 = 1$, $L_{12} = 0.3a/(10^{12} \text{cm})$, $a^2/2 = 1.3R_{\text{WR}}$, and $R_{\text{WR}} = 20R_{\text{0}}$. $\Delta \Omega \approx 0.2$ and 0.5 for the models I and II, respectively.

Provided that complete disintegration of nuclei occurs, as envisaged above in this scenario, the flux of nucleons dissolved from nuclei in the case of monoenergetic acceleration in the reconnection regions (model I) is

$$N_n = P_\Lambda/m_\gamma \gamma_{\Lambda} \approx 4 \times 10^{34} \xi M_{-5} v_\Lambda^2(r)/\gamma_{\Lambda} \text{ N s}^{-1}. \quad (7)$$

In the case of nuclei with the power law spectrum (model II) the differential spectrum of nucleons is,

$$dN_n/d\gamma_n dt = K \gamma_n^{-\delta}, \quad (8)$$

in the range $\gamma_2 = \gamma_{\text{min}}$ (Eq. 4) and applying $\gamma_1 = 10^3$, $\delta$ is the spectral index, and the normalization constant is equal to $K = A P_\Lambda/(A m_\gamma \ln(\gamma_2/\gamma_1))$ for $\delta = 2$ and $K = A P_\Lambda (\delta - 2)/(A m_\gamma (\gamma_1^{2-\delta} - \gamma_2^{-\delta}))$ for $\delta > 2$. The spectral index of particles accelerated at the shock depends on the Alfvénic Mach number of the plasma flowing through the shock. It is defined as $M = v_\Lambda/v_{\text{Alf}}$ (see e.g. Schlickeiser 2002), where $v_\Lambda$ is the Alfvén velocity defined above. For the parameters of WR 20a, we estimate the Alfvén Mach number as a function of distance from the star on $M \approx r$, where the density of the wind and the magnetic field strength at the shock location are calculated above. We conclude that the outer part of the shock, $r \approx 1$, fulfills the condition of the strong shock, $M \gg 1$, although in the inner part, the Alfvén shock is probably created. Therefore, we consider the power law spectrum of accelerated nuclei with indexes between 2 (characteristic for a strong shock) and 2.3 (for a weaker shock).

Since the winds of the massive stars are very dense, protons from disintegration of nuclei have chance to interact with the matter of the winds. The characteristic time scale for collision of relativistic protons with the matter of the wind is $\tau_{\text{pp}} \approx (\sigma_{\gamma p} \rho)^{-1} \approx 10^4 r^3 v_\Lambda(r)/M_{-5} s$, where $\rho$ is the density of the stellar wind estimated above. For the applied above parameters of the wind, $v_\Lambda = 1$ and $M_{-5} = 3$, $\tau_{\text{pp}}$ is shorter than escape time scale of protons along the shock, $\tau_\gamma$, at distances less than $r \approx 5$, i.e. in the main part of the active shock.

We calculate the $\gamma$-ray spectra from decay of pions which are produced by protons in collisions with the matter of the wind by integrating the injection spectra of protons (monoenergetic or the power law) over the active part of the shock and applying the break model for hadronic interactions developed by Wdowczyk & Wolfendale (1987) which is suitable for the considered energies of relativistic protons. Only single interaction of proton with the matter has been included. These high energy $\gamma$-rays originate relatively close to the surface of the massive stars, $r \sim 1.3 - 5$, and therefore they can suffer absorption in the thermal radiation coming from the stellar surfaces. In Fig. 2a, we show the optical depths for $\gamma$-rays as a function of their energies. It is clear that most of the $\gamma$-rays with energies above a few tens GeV are absorbed and initiate the inverse Compton (IC) $e^\pm$ pair cascades in the radiation field of the massive stars. In order to determine the final $\gamma$-ray spectra which escape from the binary system, we apply the Monte Carlo code developed by Bednarek (2000) which follow the IC $e^\pm$ pair cascade in the anisotropic radiation field of the massive star. It is assumed that primary $\gamma$-rays are produced by protons isotropically at the active part of the shock. The phase averaged $\gamma$-ray spectra escaping from the binary system are shown in Fig. 2b for the monoenergetic and the power law spectra of injected protons, assuming the distance to the WR 20a equal to 5 kpc (the average value estimated from the reported range (2.5-8) kpc, Churchill et al. 2004). These spectra have been normalized to the $\gamma$-ray fluxes observed from the EGRET sources, 2EG J1021-5835 and GeV J1025-5809 and 2EG J1049-5847 and GeV J1047-5840 (see Thompson et al. 1995, Lamb & MaComb 1997), which are equal to $\sim 10^{-7} \text{ ph cm}^{-2} \text{s}^{-1}$ above 1 GeV. Based on these normalizations we derive the required acceleration efficiencies of nuclei at the shock equal to $\sim 13.5\%$ (for the monoenergetic injection) and $\sim 12.3\%$ and $\sim 8\%$ for the power law
injection with the spectral indexes 2 and 2.3, respectively. The \( \gamma \)-ray spectra predicted by this model (shown in Fig. 2b) should be detectable above 100 GeV by the Cherenkov telescopes on the Southern hemisphere such as the HESS and CANGAROO III or at lower energies by the MAGIC telescope (e.g. HESS II).

The \( \gamma \)-ray fluxes produced by protons at the shock region in collisions with the matter of the winds via decay of pions are also accompanied by the high energy neutrinos. Let us estimate the muon neutrino event rate from our example binary system WR 20a. The spectra of muon neutrinos, calculated exactly as the primary spectra of \( \gamma \)-rays, are shown in Fig. 3a. They are clearly above the atmospheric neutrino background (ANB) and above the sensitivity limit of the 1 km\(^2\) neutrino detector of the IceCube type. In the case of monoenergetic injection of nuclei they are also above the present sensitivity of the AMANDA II neutrino detector. However due to the localization of the source on the Southern hemisphere this model cannot be at present excluded.

The number of events produced by the muon neutrinos in the 1 km\(^2\) neutrino detector of the IceCube type can be estimated from

\[
N_\mu = \frac{S}{4\pi D^2} \int P_{\nu\rightarrow\mu}(E_\nu) \frac{dN_\nu}{dE_\nu d\Omega} dE_\nu,
\]

(9)

where \( S = 1 \text{ km}^2 \) is the surface of the detector, \( P_{\nu\rightarrow\mu}(E_\nu) \) is the energy dependent detection probability of muon neutrino (Gaisser & Grillo 1987).

Applying estimated above values for acceleration efficiency, the distance to WR 20a equal to 5 kpc, and typical wind parameters of the WR stars, \( M_\ldots = 3, v_\ldots = 1 \), we predict in the case of the model I \( N_\mu \approx 23.5 \) (23) muon neutrino events inside 1 km\(^2\) detector in 1 yr for the case of neutrinos arriving from the directions close to the horizon (no absorption inside the Earth) and from the nadir (partial absorption inside the Earth). The neutrino event rate predicted in the case of model II is \( N_\mu \approx 21.6 \) (20.5) and \( \sim 3.7 \) (3.5) muon neutrino events in 1 km\(^2\) in 1 yr provided that nuclei are accelerated with the power law spectrum and spectral indexes equal to 2 and 2.3, respectively. In these calculations we applied the Earth shadowing factors from Gandhi (2000). The high event rates predicted for the nuclei injected with the flat spectra should be easily tested by the future 0.1 km\(^2\) Antares telescope.

A fraction of neutrons, \( \eta \), dissolved from nuclei can fall onto the surfaces of the massive stars. \( \eta \) can be calculated as a function of the shock distance from the WR star, \( r \), and the Lorentz factor of neutrons, \( \gamma_n \), from the formula

\[
\eta(\gamma_n) = \frac{2F(\cos \theta_D, 1)}{F(\cos \theta_D, 1) + F(\cos \theta_D, -\cos \theta_D)}
\]

(10)

where \( F(x,y) = \int_x^y \sigma_A n_{WR}(\varepsilon, \cos \theta)(1 + \cos \theta) d\varepsilon d\cos \theta \), \( \theta_D \) is the angle intercepted by the massive star observed from the distance \( r \). We have performed such calculations and found that \( \eta \) weekly depends on the Lorentz factor of nuclei, except for the Lorentz factors for which efficient fragmentation starts to occur. It is equal to \( \eta \sim 0.52, 0.33, 0.22 \) for \( R_{sh} = 1.3R_{WR}, 2R_{WR}, 3R_{WR} \), respectively. The Lorentz factors of neutrons extracted from nuclei are equal to the Lorentz factors of the parent nuclei \( \gamma_A = \gamma_n \). Neutrons which fall onto the surface of the massive stars produce charged pions in collisions with the matter of the stellar atmospheres. Pions from the first interaction are produced at the characteristic densities \( \rho \sim 3 \times 10^{14} \text{ cm}^{-3} \) (estimated from the p-p cross section and applying the characteristic dimension of the stellar atmosphere \( \sim 0.1R_{WR} \)). In such conditions pions with the Lorentz factors \( \gamma_n = 10^6 \) decay before interacting with the matter since they pass only \( \sim 0.3 \text{ g cm}^{-2} \). This conclusion has been checked by performing the Monte Carlo simulations of the interaction process of neutrons with the stellar atmosphere which preliminary results has been published in Bartosik, Bednarek & Sierpowska (2003).

We have calculated the spectra of neutrinos produced by neutrons in the atmospheres of the massive stars applying the scale break model for hadronic interactions, and taking
above), i.e. within $\alpha$ period of the binary system as postulated by the emission produced by neutrons should be modulated with the orbital period of the system. However, in the case of WR 20a, which inclination angle is $\sim 74.5^\circ$ (Bonanos et al. 2004), the observer is located inside the eclipsing cone. The neutrino signal produced by neutrons should be modulated with the orbital period of the binary system as postulated by the emission angle $\alpha$. This may help extracting the neutrino signal produced by neutrons from the atmospheric background.

The $\gamma$-ray and neutrino signals from such close binaries as WR 20a can be strongly variable due to the very unstable position of the wind collision region (short cooling time scales of thermal plasma). Small change in one of the winds may push the wind collision region onto the surface of one of the stars changing somewhat conditions in the acceleration region.

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Figure 3. The differential spectra of muon neutrinos produced by protons interacting with the matter of the stellar winds (a) and by neutrons interacting with the matter of the stellar atmosphere (b). Specific neutrino spectra are calculated for the parameters as the corresponding $\gamma$-ray spectra shown in Fig. 2b (these same curve styles). The dot-dot-dot-dashed curves indicate the atmospheric neutrino background (horizontal - upper curve, and vertical - lower curve) within a $1^\circ$ of the source (Lipari 1993), the thin dotted line shows the 3 yr sensitivity of the IceCube detector (Hill 2001), and the thick dotted line shows the upper limit obtained by AMANDA II detector (Ackermann et al. 2004).