NEUTRINO OSCILLATION FROM HIDDEN GRB JETS

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RESUMEN

Los Colápsares son probablemente la fuente de los destellos de rayos gama de larga duración (lGRBs por sus siglas en inglés). Se han observado lGRBs con duración de miles y hasta decenas de miles de segundos, lo que hace poco probable que sus jets sean producidos por motores centrales accionados por neutrinos. En éste contexto, el mecanismo de Blandford-Znajek es un buen candidato para explicar la producción de jets a lo largo del eje de rotación sin la necesidad de grandes tasas de acréción. Estos motores centrales requieren campos magnéticos de $10^{12} \, G < B < 10^{15} \, G$ atravesando la parte interna del disco de acréción y el agujero negro de Kerr. Obtenemos la auto-energía y el potencial efectivo a orden $O(1/M^4)$ para una bola de fuego hecha de electrones, protones, neutrones y sus antiparticle respectivas y con campo magnético débil y fuerte. Consideramos que la energía de los neutrinos producidos durante el colapso gravitacional, así como en fusión de objetos compactos o en la bola de fuego en sí, por procesos de aniquilación de pares, decaimiento beta y por procesos bremsstrahlung nucleónicos es de entre 1 y 100 MeV. Muchos de dichos neutrinos se propagan a través de la bola de fuego y pueden oscilar resonantemente. Usando la mezcla de dos neutrinos estudiamos la oscilación de neutrinos.

ABSTRACT

Collapsars are the likely progenitors of Long GRBs (IGRBs). IGRBs have been observed to last for thousands to tens of thousands of seconds, thus making unlikely the neutrino-driven engine as the main mechanism for driving the jets. In this context, the Blandford-Znajek mechanism seems likely to explain the production of rotational-axis directed jets without the need for large accretion rates. These engines, require magnetic fields between $10^{12} \, G < B < 10^{15} \, G$ threading the inner disk, Kerr-BH region to exist. We derive the neutrino self-energy and the effective potential up to $O(1/M^4)$ in a weakly and highly magnetized GRB fireball flow which is made up of electrons, protons, neutrons and their anti-particles. We consider neutrino energies of 1-100 MeV which are produced during stellar collapse, merger events or in the fireball itself by electron-positron annihilation, inverse beta decay and nucleonic bremsstrahlung processes. Many of these neutrinos propagate through the fireball and may oscillate resonantly. Using two-neutrino mixing we study the possibility of these oscillations.

Key Words: binaries: general — black hole physics — gamma-ray bursts: general — stars: evolution — stars: magnetic field — stars: massive

Collapsars (Woosley 1993; MacFadyen & Woosley 1999) are generally accepted as the progenitors of Long Gamma-Ray Bursts (lGRBs). This is supported by the prediction that IGRBs are accompanied by an energetic supernova (SN). There is strong evidence that six IGRBs are spectroscopically associated with SNe (GRB980425/SN 1998bw, GRB 030329/SN 2003dh, GRB 031203/SN 2003bw, GRB 060218/SN 2006aj and GRB 100316D/SN 2010bh, see Hjorth & Bloom 2011 and GRB120422/SN 2012bz, see Malesani et al. 2012). There is also evidence of such a relation in another two dozen events.

One of the major issues with the Collapsar model, is the issue of producing a massive stellar core with enough angular momentum (see, e.g., Heger et al. 2005). To solve this issue, one of the proposed mechanisms is to use a binary to transfer angular momentum back to the star late in its evolution (see, e.g. Paczynski 1998). Lee et al. (2002) studied several Galactic black-hole binaries (BHBs) and reproduced the natal spin of the BHs. Using these results, Brown et al. (2007), Brown et al. (2008) and Moreno Méndez et al. (2011) estimated the rotational energies of the BHs at the time they were born and estimated whether a Blandford-Znajek (BZ; Blandford & Znajek 1977) central engine may have triggered a IGRB in any of these BHBs. In Brown et al. (2008) it was found that, most likely, LMC X−3 is a likely IGRB relic, unlike most Galactic BHBs where too much rotational energy may have destroyed the BZ engines too soon. In Moreno Méndez et al. (2008) and Moreno Méndez (2011) it was shown that the large spins observed in some Galactic high-mass X-ray binaries (HMXBs) cannot be natal but they rather have to be ac-
quired post-BH formation. Thus, these HMXBs are not good candidates for relics of IGRBs. Woosley & Heger (2012) have suggested that this binary channel may not produce IGRBs as the internal magnetic fields and/or dynamos may extract the angular momentum off the stellar core. Moreno Mendoza (2014) shows that this difficulty may be overcome if the internal field is low enough and a magnetar-like field can be produced during core collapse.

Now, LMXBs (low-mass XBs), which have too much rotational energy may initially trigger an internal jet which quickly dies after the BZ engine is destroyed as a SN is launched. Also, many successful IGRBs may not be pointed in our direction. Similarly, were some stars to quickly dies after the BZ engine is destroyed as a SN acquired post-BH formation. Thus, these HMXBs are not good candidates for relics of IGRBs.

The flavor of the neutrinos, they interact with different effective potentials because electron neutrinos ($\nu_e$) interact with electrons via both, neutral and charged currents, whereas muon and tau neutrinos ($\nu_\mu$ and $\nu_\tau$) interact with electrons only via the neutral current (NC). This induces a coherent effect in which maximal conversion of $\nu_e$ into $\nu_\mu$ ($\nu_\tau$) takes place even for a small, intrinsic, mixing angle. The resonant conversion of neutrino from one flavor to another due to the medium effect, is well known as the Mikheyev-Smirnov-Wolfenstein effect. In this work, we study the propagation and resonant oscillation of thermal neutrinos in both, a weakly- and highly-magnetized fireball. We take into account the two-neutrino mixing (solar, atmospheric and accelerator parameters). Finally, we discuss our results in the GRB framework.

1. NEUTRINO EFFECTIVE POTENTIAL

We use the finite-temperature, field-theory formalism to study the effect of a heat bath on the propagation of elementary particles (Sahu et al. 2007; Nieves 1990). The effect of the magnetic field is taken into account through Schwinger’s proper-time method (Schwinger 1951). The effective potential of a particle is calculated from the real part of its self-energy diagram. The neutrino field equation of motion in a magnetized medium is,

$$ [k - \Sigma(k)]\Psi_L = 0, $$

where the neutrino self-energy operator $\Sigma(k)$ is a Lorentz scalar which depends on the characterized parameters of the medium, as for instance, chemical potential, particle density, temperature, magnetic field, etc. Solving this equation and using Dirac’s algebra we can write the dispersion relation $V_{eff} = k_0 - |k|$ as a function of Lorentz scalars:

$$ V_{eff} = b - c \cos \phi - a_1 |k| \sin^2 \phi, $$

where $\phi$ is the angle between the neutrino momentum and the magnetic field vector. The Lorentz scalars $a$, $b$ and $c$ are functions of the neutrino energy, momentum and the magnetic field. They can be calculated from the neutrino self-energy due to charge-current and neutral-current interaction with the background particles.

1.1. One-loop neutrino self-energy

The total one-loop neutrino self-energy in a magnetized medium is given by

$$ \Sigma(k) = \Sigma_W(k) + \Sigma_Z(k) + \Sigma_i(k) $$

where $\Sigma_W(k)$ is the W-exchange, $\Sigma_Z(k)$ is the Z-exchange and $\Sigma_i(k)$ represents the tadpole. The W-exchange diagrams to the one-loop self-energy is

$$ -i\Sigma_W(k) = R \left[ \int \frac{d^4p}{(2\pi)^4} \left( -ig \gamma_\mu iS_\mu(p) \right) \right] L $$

where $g^2 = 4\sqrt{3}G_F M_W^2$ is the weak coupling constant, $W^{\mu\nu}$ depicts the W-boson propagator which, in the eB$\ll M_W$ limit and in unitary gauge, is given by

$$ W^{\mu\nu}(q) = \frac{g^{\mu\nu}}{M_W^2} \left( 1 + \frac{q^2}{M_W^2} \right) - \frac{g^{\mu s} g^{\nu t}}{M_W^2} + \frac{3ie}{2M_W^2} F^{\mu\nu}, $$

here $M_W$ is the W-boson mass, $g^{\mu\nu}$ is the metric tensor, $F^{\mu\nu}$ is the electromagnetic field tensor and $e$ being the magnitude of the electron charge. $S_\mu(p)$ stands for the charged lepton propagator which can be separated into two charged propagators; one in presence of a uniform background magnetic field ($S^0_\mu(p)$) and another in a magnetized medium ($S^3_\mu(p)$). It can be written as,

$$ S_\mu(p) = S^0_\mu(p) + S^3_\mu(p). $$

Assuming that the z-axis points in the direction of the magnetic field B, we can express the charged lepton propagator in presence of a uniform background magnetic field as,

$$ iS^0_\ell(p) = \int_0^\infty e^{i(p.s)} G(p, s) \, ds, $$
where the functions $\phi(p, s)$ and $G(p, s)$ are given by,
\[
G(p, s) = \sec^2 z \left[ A + iB \gamma_5 \right] + m_\ell (\cos^2 z - i\Sigma^3 \sin z \cos z) \right]
\]
\[
\phi(p, s) = is\left(p_0^2 - m_\ell^2\right) - is\left[p_0^2 + \frac{\tan z}{z} - p_1^2\right], \quad (8)
\]
here $m_\ell$ is the mass of the charged lepton, $p_0^2 = p_0^2 + p_3^2$ and $p_1^2 = p_1^2 + p_2^2$ are the projections of the momentum on the magnetic field direction and $z = eBs$. Additionally, the covariant vectors are given as follows,
\[
A_\mu = p_\mu - \sin^2 z\left(p \cdot u\right) \mu \cdot b, \quad B_\mu = \sin z \cos z\left(p \cdot u\right) \mu \cdot b, \quad (9)
\]
and
\[
\Sigma^3 = \gamma_3 b \mu.
\]
The charged lepton propagator in a magnetized medium is given by,
\[
S_\ell^\beta(p) = i\eta_\ell(p \cdot u) \int_{-\infty}^{\infty} e^{\phi(p, s)} G(p, s) ds, \quad (10)
\]
where $\eta_\ell(p \cdot u)$ contains the distribution functions of the particles in the medium which are given by:
\[
\eta_\ell(p \cdot u) = \frac{\theta(p \cdot u)}{e^{\beta(p \cdot u - \mu_\ell)} + 1} + \frac{\theta(-p \cdot u)}{e^{\beta(p \cdot u - \mu_\ell)} + 1},
\]
where $\beta$ and $\mu_\ell$ are the inverse of the medium temperature and the chemical potential of the charged lepton.

The Z-exchange diagram to the one-loop self-energy is
\[
-\Sigma_Z(k) = R \left[ \int \frac{d^4p}{(2\pi)^4} \left( -i\frac{g}{\sqrt{2} \cos \theta_W} \right) \gamma_\mu iS_\nu(p) \right] \left( -i\frac{g}{\sqrt{2} \cos \theta_W} \right) \gamma_\nu iZ^{\mu\nu}(q) L, \quad (11)
\]
$\theta_W$ is the Weinberg angle, $Z^{\mu\nu}(q)$ is the Z-boson propagator in vacuum, $S_\nu$ is the neutrino propagator in a thermal bath of neutrinos. The Tadpole diagram to the one-loop self-energy is
\[
i\Sigma_\ell(k) = R \left[ \left( -i\frac{g}{2 \cos \theta_W} \right)^2 \gamma_\mu iZ^{\mu\nu}(0) \int \frac{d^4p}{(2\pi)^4} \right] \text{Tr} \left[ \gamma_\nu \left( C_V + C_A \gamma_5 \right) iS_\ell(p) \right] L, \quad (12)
\]
where the quantities $C_V$ and $C_A$ are the vector and axial-vector coupling constants which come in the neutral-current interaction of electrons, protons ($p$), neutrons ($n$) and neutrinos with the $Z$ boson. Their forms are as follows,
\[
C_V = \begin{cases} 
-\frac{1}{2} + 2\sin^2 \theta_W & e \\
\frac{1}{2} & \nu_
u \\
-\frac{1}{2} - 2\sin^2 \theta_W & p \\
\frac{1}{2} & \nu_e \\
-\frac{1}{2} & \nu_n \\
\frac{1}{2} & \nu_\tau
\end{cases}, \quad (13)
\]
and
\[
C_A = \begin{cases} 
-\frac{1}{2} & \nu, p \\
\frac{1}{2} & e, n
\end{cases}.
\]

By evaluating each contribution of one-loop neutrino energy, the effective potential in the weak (Sahu et al., 2009b) ($B \ll m_e^2/e$) and strong ($B \gg m_e^2/e$) (Fraija & Moreno Mendoza, 2014b) -field limit is given respectively by
\[
V_{eff,w} = \frac{\sqrt{2} G_F m_\nu^3 B}{\pi^2} \sum_{l=0}^{\infty} (-1)^l \sinh \alpha L K_1(\sigma) \left( 1 + C_{Ve} \right) + \frac{m_\nu^2}{m_W^2} \left( \frac{3}{2} + \frac{2E_\nu}{m_\nu^2} \right) B \frac{B}{B_e} - \left( 1 - \frac{C_{Ve} + m_\nu^2}{m_W^2} \right) \times \left( 1 - \frac{2E_\nu}{m_\nu^2} + \frac{B}{B_e} \right) \cos \phi \right) - \frac{m_\nu^2}{m_W^2} E_\nu \sum_{l=0}^{\infty} (-1)^l \times \cosh \alpha \left( \frac{3}{4} L K_0(\sigma) + \frac{K_1(\sigma)}{\sigma} - \frac{K_1(\sigma)}{\sigma} \cos \phi \right), \quad (15)
\]
and
\[
V_{eff,s} = \sqrt{2} G_F m_\nu^3 \sum_{l=0}^{\infty} (-1)^l \sinh \alpha \left[ \left( 1 + \frac{3}{2} \frac{m_\nu^2}{M_W^2} - \frac{eB}{M_W} \right) \right] K_1(\sigma) \left( \frac{2}{\sigma} - \frac{B}{B_e} K_1(\sigma) \right) \frac{B}{B_e} \left( 1 + \frac{m_\nu^2}{2M_W^2} - \frac{eB}{M_W} \right) \times \frac{K_1(\sigma)}{M_W^2} \times K_0(\sigma) \left( 2 - \frac{4B}{B_e} + \frac{16}{\sigma^2} \right) K_1(\sigma), \quad (16)
\]
where,
\[
\alpha = \beta \mu(l + 1) \quad \text{and} \quad \sigma = \beta m_e(l + 1).
\]

2. TWO-NEUTRINO MIXING

Here we consider the neutrino oscillation process $\nu_e \leftrightarrow \nu_{\mu,\tau}$. The evolution equation for the propagation of neutrinos in the above medium is given by
\[
i \begin{pmatrix} \nu_e \\ \nu_{\mu} \end{pmatrix} = \left( V_{eff} - \Delta \cos 2\theta \begin{pmatrix} \frac{1}{2} & \sin 2\theta \\ \frac{1}{2} & 0 \end{pmatrix} \begin{pmatrix} \nu_e \\ \nu_{\mu} \end{pmatrix} \right), \quad (18)
\]
where $\Delta = \delta m^2/2E_\nu$, $V_{eff}$ is the potential difference between $V_{\nu_e}$ and $V_{\nu_{\mu,\tau}}$, $E_\nu$ is the neutrino energy and $\theta$...
is the neutrino mixing angle. The conversion probability for the above process at a given time \( t \) is given by

\[
P_{\nu_e \rightarrow \nu_\mu (\nu_\tau)}(t) = \frac{\Delta^2 \sin^2 2\theta}{\omega^2} \sin^2 \left( \frac{\omega t}{2} \right),
\]

with

\[
\omega = \sqrt{(V_{\text{eff}} - \Delta \cos 2\theta)^2 + \Delta^2 \sin^2 2\theta}.
\]

The potential for the above oscillation process is given by eq. (2). The oscillation length for the neutrino is given by

\[
L_{\text{osc}} = \frac{L_\nu}{\sqrt{\cos^2 2\theta (1 - \frac{V_{\text{eff}}}{\Delta \cos 2\theta})^2 + \sin^2 2\theta}},
\]

and the resonance condition can be written as

\[
V_{\text{eff}} - \frac{\delta m^2}{2E_\nu} \cos 2\theta = 0.
\]

Following Fraija (2014a) we apply the neutrino oscillation parameters in the resonance condition (eq. 22).

3. RESULTS AND CONCLUSIONS

We have plotted the resonance condition for neutrino oscillation in weakly and strongly magnetized collapsars at the base of the jet when the temperature is in the range of 1-25 MeV. We have used effective potential \( V_{\text{eff}} \) up to order \( M_4^W \) in the strong- and weak-field limit, the best parameters for the two- solar (Aharmim & et al. 2011), atmospheric (Abe & et al. 2011) and accelerator (Athanassopoulos & et al. 1996, 1998) neutrinos and the neutrino energies \( E_\nu = 1, 5, 20, 50 \) and 100 MeV). The analysis of resonance condition shows that, the temperature as a function of chemical potential is degenerate for the weak limit and not in the strong limit. In both cases, the chemical potential is biggest when accelerator parameters are used (19 eV - 50 keV and 5 keV - 0.15 MeV, respectively) and smallest for solar parameters (0.08 eV - 5 eV and 0.25 - 50 eV, respectively).

The baryon load, resonance length and leptonic symmetry (two and three flavors) will be estimated and studied in Fraija & Moreno M´endez in progress.

NF gratefully acknowledges a Luc Binette-Fundación UNAM Postdoctoral Fellowship. EMM was supported by a CONACyT fellowship and projects CB-2007/83254 and CB-2008/101958. This research has made use of NASAs Astrophysics Data System as well as arXiv.
Fig. 2. Plot of temperature ($T/m_e$) as a function of chemical potential ($\mu = m_e 10^p$) for which the resonance condition is satisfied. We have used the best parameters of the two-flavor solar (top), atmospheric (middle) and accelerator (bottom) neutrino oscillation, $B=10^9$, and taken five different neutrino energies: $E_{\nu} = 1$ MeV (red dashed line), $E_{\nu} = 5$ MeV (blue dot-dashed line), $E_{\nu} = 20$ MeV (magenta dotted line), $E_{\nu} = 50$ MeV (black thin-solid line) and $E_{\nu} = 100$ MeV (brown thick-solid line).

REFERENCES

Abe, K. & et al. 2011, Physical Review Letters, 107, 241801
Aharmim, B. & et al. 2011, ArXiv e-prints
Athanassopoulos, C. & et al. 1996, Physical Review Letters, 77, 3082
—. 1998, Physical Review Letters, 81, 1774
Blandford, R. D. & Znajek, R. L. 1977, MNRAS, 179, 433
Brown, G. E., Lee, C., & Moreno Mendoza, E. 2007, ApJ, 671, L41
—. 2008, ApJ, 685, 1063
Fraija, N. 2014a, MNRAS, 437, 2187
—. 2014b, arXiv:1401.1581
Fraija, N. & Moreno Mendoza, E. 2014, arXiv:1401.1908
Fraija, N. & Moreno Mendoza, E. 2014, arXiv:1401.3787
García, A. B. & Sahu, S. 2007, Modern Physics Letters A, 22, 213
Heger, A., Woosley, S. E., & Spruit, H. C. 2005, ApJ, 626, 350
Hjorth, J. & Bloom, J. S. 2011, ArXiv e-prints
Lee, C., Brown, G. E., & Wijers, R. A. M. J. 2002, ApJ, 575, 996
MacFadyen, A. I. & Woosley, S. E. 1999, ApJ, 524, 262
Malesani, D., Schulze, S., Kruehler, T., Fynbo, J. P. U., Hjorth, J., Milvang-Jensen, B., Watson, D., de Ugarte Postigo, A., Tanvir, N. R., Tagliaferri, G., Sollerman, J., Xu, D., Stritzinger, M. D., & De Cia, A. 2012, Central Bureau Electronic Telegrams, 3100, 2
Moreno Mendoza, E. 2011, MNRAS, 413, 183
—. 2014, ApJ, 781, 3
Moreno Mendoza, E., Brown, G. E., Lee, C., & Park, I. H. 2008, ApJ, 689, L9
Moreno Mendoza, E., Brown, G. E., Lee, C., & Walter, F. M. 2011, ApJ, 727, 29
Nieves, J. F. 1990, Phys. Rev. D, 42, 4123
Paczynski, B. 1998, ApJ, 494, L45
Sahu, S., Fraija, N., & Keum, Y.-Y. 2009a, Phys. Rev. D, 80, 033009
—. 2009b, J. Cosmology Astropart. Phys., 11, 24
Schwinger, J. 1951, Physical Review, 82, 664
Woosley, S. E. 1993, ApJ, 405, 273
Woosley, S. E. & Heger, A. 2012, ApJ, 752, 32