NEUTRINO OSCILLATION FROM MAGNETIZED STRANGE STARS

Nissim Fraija1 and Enrique Moreno Méndez2

ABSTRACT

Strange-quark matter (SQM) is a likely candidate of the ground state of nuclear matter. Along with many other equations of state (EoSs), SQM seemed to be severely constrained by the recent discoveries of the 1.97 $M_\odot$ PSR J1614-2230 and the 2.01 $M_\odot$ PSR J0348+0432. However with new, $O(\alpha_C^2)$, perturbative calculations, SQM seems to be able to accommodate masses as large as $\sim 2.75 M_\odot$. The literature of SQM stars or strange stars includes estimates of internal magnetic fields as large as $10^{20}$ G, which are unlikely to be formed as they would require $\sim 10^{57}$ erg to be produced. Nonetheless, if strange stars may hold magnetar-strength fields ($10^{15}$ G), their internal fields are likely to reach magnetic fields as large as $10^{17}$ G. We consider neutrinos with energies of some MeV and oscillation parameters from solar, atmospheric and accelerator experiments. We study the possibility of resonant oscillation of neutrinos in strange stars.

Key Words: neutrinos — stars: evolution — stars: magnetic field — stars: massive — stars: neutron

1. INTRODUCTION

Stars that produce Fe cores end up their lives as gravitational-core-collapse (GCC) supernovae (SNe), most likely driven by neutrinos (Bethe et al. 1979; Bethe 1990), if the compact object ends below a mass of some $\sim 1.97 M_\odot$. The threshold mass is determined by the nuclear equation of state (EoS) but it could be partially influenced by other factors like centrifugal support on millisecond (ms) pulsars (PSRs) and/or magnetic pressure in, e.g., magnetars (see, e.g., [Lattimer & Prakash 2010] for a review).

Recently, the nuclear EoS has been further constrained by the discovery of two massive pulsars, the 1.97-$M_\odot$ PSR J1614-2230 (Demorest et al. 2010), and the 2.01-$M_\odot$ PSR J0348+0432 (Antoniadis et al. 2013). Nonetheless, the nuclear EoS is still not known and many possibilities exist. Among them, is the possibility that as the pressure and density grow in the inside of a neutron star (NS), the quarks that originally conform a baryon mix with the quarks from neighbouring baryons until all quarks become free and produce quark matter (QM; see, e.g., Itoh 1970, Witten 1984), suggested that the Fermi energy of QM could be lowered by allowing an extra degree of freedom in the quarks, namely, another flavor. Thus, roughly $\sim 1/3$ (minus the chemical potential of the new quark flavor), of the quarks are converted to the next lower mass quark, which is the strange quark. Thus this matter is named strange quark matter (SQM). According to $O(\alpha_C^2)$ estimates by (Kurkela et al. 2010), with $\alpha_C$ the strong coupling constant, the maximum mass of SQM stars can be $M_{\mathrm{SQM}} \geq 2.75 M_\odot$. Following Farhi & Jaffe (1984), Haensel et al. (1986), Alcock et al. (1986) and Glendenning (1997), use the MIT-bag model to study the structure and characteristics of strange stars.

Using the MIT-bag model the values for the bag constant, the strong-coupling constant and the mass of the strange quark can be varied. For values above 0.5 for the strong coupling constant strange quark matter can become negatively charged as the density increases (Moreno Méndez et al. 2013 and Moreno Méndez & Page 2013, in preparation). Figure 1 shows how electrons are replaced by positrons once the baryonic density reaches values $0.8$ fm$^{-3}$. In such case, a region where $e^-$-$e^+$ pairs annihilate into neutrinos may be formed. In this
scenario neutrinos can be produced in this region even after the strange matter has cooled below 1 MeV.

SQM stars are bound by QCD and, thus, in principle, they could sustain a rather strong, internal, magnetic field. An internal field, perhaps as large as an equipartition one (during GCC), i.e., \( \sim 10^{17} \, \text{G} \) could in principle exist. Here we will test fields as large as \( 3 \times 10^{16} \, \text{G} \).

Now, the properties of neutrinos get modified when they propagate in a strongly magnetized medium. Depending on the flavor, a neutrino would feel a different effective potential because electron neutrinos (\( \nu_e \)) interact with electrons via both neutral and charged currents (CC), whereas muon (\( \nu_\mu \)) and tau (\( \nu_\tau \)) neutrinos interact 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 (Wolfenstein 1978). In this work, we study the propagation and resonant oscillation of thermal neutrinos in strange stars. We take into account the two-neutrino mixing (solar, atmospheric and accelerator parameters). Finally, we discuss our results in the strange stars framework.

2. NEUTRINO PRODUCTION FOR LARGE COUPLING

Using the MIT-bag model the values for the bag constant, the strong coupling constant and the mass of the strange quark can be varied. For values above 0.5 for the strong coupling constant strange quark matter can become negatively charged as the density increases. The figure on the right shows how electrons are replaced by positrons once the baryonic density reaches values 0.8 fm^{-3}. In such case, a region where e^−e^+ pairs annihilate into neutrinos forms. In this scenario neutrinos can be produced in this region even after the strange matter has cooled below 1 MeV.

3. NEUTRINO EFFECTIVE POTENTIAL

We use the finite temperature field theory formalism to study the effect of heat bath on the propagation of elementary particles. The effect of magnetic field is taken into account through Schwinger's proper-time method (Schwinger 1951; Sahu et al. 2009a,b). 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 (Fraija 2013, submitted)

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

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 the Dirac algebra, the dispersion relation \( V_{eff} = k_0 - |k| \) as a function of Lorentz scalars can be written as

\[
V_{eff} = b - c \cos \phi - a |k| \sin^2 \phi,
\]

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

**Fig. 1.** Plot of the baryon density vs. the log of the frequency (in MeV). The upper five curves are plasma frequencies, whereas the lower five curves are the gyrofrequencies. The upper continuous line is the total plasma frequency. The Dash-dot-dot-dot line is for the up quark. The slimmer continuous line is for the down quark, and the dotted line is the strange quark. The dashed curve is for the electrons (lower density) and the positrons (higher density). The next top-to-down continuous curve is for the total gyrofrequency. The dashed one is for the electrons and positrons. The dash-dotted line is for the up quarks. The last two curves are for down and strange quarks. This plot is for a \( 10^{12} \, \text{G} \) magnetic field, the gyrofrequency curves get displaced upwards linearly for higher magnetic fields.
3.1. One-loop neutrino self-energy

The total one-loop neutrino self-energy in a magnetized medium is given by [Erdas et al. 1998; Fraija 2014b]

\[ R \Sigma_W(k) = R[a_{W_\perp} k_\perp + b_W k + c_W b] L \]  

where the Lorentz scalars are given by the following equations

\[ a_{W_\perp} = \frac{\sqrt{2} G_F}{M_W^2} \left[ \left( N_e - \bar{N}_e \right) + k_3 \left( N_e^0 - \bar{N}_e^0 \right) \right] + \frac{eB}{2\pi^2} \int_0^\infty dp_3 \sum_{n=0}^{\infty} \left( 2 - \delta_{n,0} \right) \left( \frac{m_e^2}{E_n} - \frac{H}{E_n} \right) (f_{e,n} + \bar{f}_{e,n}), \]  

\[ b_W = \frac{\sqrt{2} G_F}{\pi^2} \left[ \frac{3 m_e^2}{M_W^2} + \frac{E_e^2}{M_W^2} \right] (N_e - \bar{N}_e) + \frac{eB}{M_W^2} \int_0^\infty dp_3 \sum_{n=0}^{\infty} \left( 2 - \delta_{n,0} \right) \left( k_3 \left( N_e^0 - \bar{N}_e^0 \right) \right) - \frac{eB}{2\pi^2 M_W^2} \int_0^\infty dp_3 \sum_{n=0}^{\infty} \left( 2 - \delta_{n,0} \right) \left( 2k_3 E_n \delta_{n,0} + 2E_{nu} \left( E_n - \frac{m_e^2}{2E_n} \right) \right) \left( f_{e,n} + \bar{f}_{e,n} \right), \]  

\[ c_W = \frac{\sqrt{2} G_F}{\pi^2} \left[ \frac{1}{2} \left( \frac{m_e^2}{M_W^2} - \frac{k_3^2}{M_W^2} \right) \right] (N_e^0 - \bar{N}_e^0) + \frac{eB}{M_W^2} \int_0^\infty dp_3 \sum_{n=0}^{\infty} \left( 2 - \delta_{n,0} \right) \left( E_{nu} \delta_{n,0} + k_3 \right) \left( E_n - \frac{3 m_e^2}{2E_n} - \frac{H}{E_n} \right) \left( f_{e,n} + \bar{f}_{e,n} \right). \]  

The effective potential is

\[ V_{eff} = \frac{\sqrt{2} G_F m_e^3}{\pi^2 B_e} \sum_{l=0}^{\infty} \left( -1 \right)^l \sinh \alpha K_l(\sigma) \left[ 1 + C_{V_\alpha} + \frac{m_e^2}{m_W^2} \left( 1 - \frac{2E_{nu}}{m_e^2} + \frac{B}{B_e} \right) \right] \]  

where \( f(E_{e,0}) = \frac{1}{e^{\beta(E_{e,0} - \mu)} + 1} \) and the electron energy in the lowest Landau level is

\[ E_{e,0}^2 = (p_e^2 + m_e^2). \]  

Assuming that the chemical potentials (\( \mu \)) of the electrons and positrons are much smaller than their energies (\( \mu \leq E_e \)), the fermion distribution function can be written as a sum given by

\[ f(E_{e,0}) = \frac{1}{e^{[m(E_{e,0} - \mu)](l + 1)}}. \]  

In the strong magnetic field approximation, the energy of charged particles is modified confining the particles to the lowest Landau level (\( n = 0 \)). The number density of electrons will become

\[ N_e^0 = \frac{eB}{2\pi^2} \int_0^\infty dp_3 f_{e,0}. \]
The effective potential for the above oscillation process is given by eq. (10). The oscillation length for the neutrino is given by

\[ L_{osc} = \frac{L_v}{\sqrt{\cos^2 2\theta \left(1 - \frac{V_{eff} - \delta m^2}{E_\nu \cos 2\theta}\right)^2 + \sin^2 2\theta}} \]

(14)

where \( L_v = \frac{2\pi}{\Delta} \) is the vacuum oscillation length. The resonance length can be written as

\[ L_{res} = \frac{L_v}{\sin 2\theta} \]

(15)

to obtain the previous equation, we applied the resonance condition given by

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

(16)

NF gratefully acknowledges a Luc Binette-Fundación UNAM Posdoctoral 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.

5. CONCLUSIONS

We have plotted the resonance condition for neutrino oscillation in strongly magnetized strange stars. We have taken into account the temperature is in the range of 1-10 MeV and magnetic field \( B = 10^3 B_c \). We have used effective potential \( V_{eff} \) up to order \( M^4 W \), 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 \( E_\nu = 1, 5, 10, 20 \) and 30 MeV). From the plots in fig. 2, we can see that the chemical potential for electrons from accelerator parameters is largest and the one from solar parameters the smallest. Another important characteristic from fig. 2 is that the temperature is degenerated for solar and atmospheric parameters unlike the case for accelerator parameters. Further strange star observables (for instance leptonic asymmetry, minimum baryon load, etc.) will be estimated in Fraija & Moreno Méndez (2013, in preparation).

REFERENCES

Abe, K. & et al. 2011, Physical Review Letters, 107, 241801
Aharmim, B. & et al. 2011, ArXiv e-prints
Alcock, C., Farhi, E., & Olinto, A. 1986, ApJ, 310, 261
Antoniadis, J., Freire, P. C. C., Wes, N., Tauris, T. M., Lynch, R. S., van Kerkwijk, M. H., Kramer, M., Bassa, C., Dhillon, V. S., Driebe, T., Hessels, J. W., Kaspi, V. M., Kondratiev, V. I., Langer, N., Marsh, T. R., McLaughlin, M. A., Pennucci, T. T., Ransom, S. M., Stairs, I. H., van Leeuwen, J., Verbiest, J. P. W., & Whelan, D. G. 2013, Science, 340, 448
Athanassopoulos, C. & et al. 1996, Physical Review Letters, 77, 3082
—. 1998, Physical Review Letters, 81, 1774
Bethe, H. A. 1990, Reviews of Modern Physics, 62, 801
Bethe, H. A., Brown, G. E., Applegate, J., & Lattimer, J. M. 1979, Nuclear Physics A, 324, 487
Demorest, P. B., Pennucci, T., Ransom, S. M., Roberts, M. S. E., & Hessels, J. W. T. 2010, Nature, 467, 1081
Erdas, A., Kim, C. W., & Lee, T. H. 1998, Phys. Rev. D, 58, 085016
Farhi, E. & Jaffe, R. L. 1984, Phys. Rev. D, 30, 2379
Fraija, N. 2014a, MNRAS, 437, 2187
Fraija, N. 2014b, ArXiv: 1401.1581
Glendenning, N. K., ed. 1997, Compact stars. Nuclear physics, particle physics, and general relativity
Haensel, P., Zdunik, J., & Schaeffer, R. 1986, Astron. Astrophys., 160, 121
Itoh, N. 1970, Progress of Theoretical Physics, 44, 291
Kurkela, A., Romatschke, P., & Vuorinen, A. 2010, Phys. Rev. D, 81, 105021
Lattimer, J. M. & Prakash, M. 2010, ArXiv e-prints
Moreno Méndez, E., Page, D., Patiño, L., & Ortega, P. 2013, ArXiv e-prints
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
Witten, E. 1984, Phys.Rev., D30, 272
Wolfenstein, L. 1978, Phys. Rev. D, 17, 2369