Electron-, Mu-, and Tau-Number Conservation in a Supernova Core

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We study if the neutrino mixing parameters suggested by the atmospheric neutrino anomaly imply chemical equilibrium between μ- and τ-flavored leptons in a supernova (SN) core. The initial flavor-conversion rate would indeed be fast if the ν_μ-ν_τ-mixing angle were not suppressed by *second-order* refractive effects. The neutrino diffusion coefficients are different for ν_μ, $\bar{\nu}_\mu$, ν_τ and $\bar{\nu}_\tau$ so that neutrino transport will create a net μ and τ lepton number density. This will typically lead to a situation where the usual first-order refractive effects dominate, further suppressing the rate of flavor conversion. Altogether, neutrino refraction has the nontrivial consequence of guaranteeing the separate conservation of e, μ, and τ lepton number in a SN core on the infall and cooling time scales, even when neutrino mixing angles are large.

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I. INTRODUCTION

The evidence for neutrino flavor oscillations from atmospheric [1] and solar neutrinos [2] as well as the LSND experiment [3] is now so compelling that the debate in neutrino physics has fundamentally changed. The current experimental results point to a few very specific regions in the parameter space of mass differences and mixing angles. Therefore, the experimental task at hand is to verify the proposed solutions by independent means such as long-baseline oscillation experiments. From the astrophysical perspective, it is no longer an academic exercise to investigate systematically if neutrino mixings with these specific parameters lead to observable effects other than explaining the solar neutrino deficit.

One consequence of the apparent violation of e, μ, and τ flavor lepton number is that the leptons of a perfectly thermalized system are no longer characterized by six independent chemical potentials, but only by one for the charged leptons and one for the neutrinos, i.e. $\mu_e = \mu_\mu = \mu_\tau$ and $\mu_{\nu_e} = \mu_{\nu_\mu} = \mu_{\nu_\tau}$. In practice one knows of only two types of environment where neutrinos achieve thermal equilibrium and where this effect could be important, the hot early universe and the inner cores of collapsed stars. In the standard picture of the early universe, the neutrino chemical potentials are very small so that flavor equilibration would not make much of a difference except for the case of oscillations between active and sterile neutrinos [4], a possibility that we will mostly ignore.

In a supernova (SN) core, on the other hand, the electron lepton number of the progenitor star's iron core is trapped, leading to highly degenerate e and ν_e distributions. If the trapped electron lepton number were shared among all flavors, the impact on the effective equation of state would be large. Equilibration on the infall time scale of a few hundred milliseconds could well be disastrous for the SN explosion mechanism. However, in a previous paper [5] two of us found that even for maximal neutrino mixing, equilibrium between ν_e and the other flavors is not achieved on the time scale of a few seconds unless δm^2 exceeds about 10^5 eV^2. The rate of flavor conversion is very slow because the mixing angle is strongly suppressed by neutrino refractive effects. Since the current indications for neutrino oscillations point to much smaller mass differences, the electron lepton number in a SN core is almost perfectly conserved except for diffusive or convective transfer to the stellar surface.

We presently study the corresponding conversion between ν_μ and ν_τ, which has not yet been investigated. It is usually assumed that the chemical potentials for these flavors vanish in a SN core so that their distribution functions are equal, i.e. that flavor equilibrium exists from the start. However, the muon mass of 106 MeV is not

very large, considering that typical temperatures can exceed $T = 30$ MeV and that the average thermal energy of a relativistic particle is around $3T$. Beta equilibrium by reactions of the form $\nu_\mu + n \leftrightarrow p + \mu^-$ implies the condition $\Delta\mu \equiv \mu_n - \mu_p = \mu_\mu - \mu_{\nu_\mu}$, which is familiar from the electron flavor. The exact value of $\Delta\mu$ depends sensitively on the equation of state; a typical range is 50–100 MeV. Since initially the trapped μ lepton number is zero, and taking $\Delta\mu = 50$ MeV and $T = 30$ MeV as an example, we find $\mu_\mu \approx 32$ MeV and $\mu_{\nu_\mu} \approx -18$ MeV so that there is a significant excess of μ^- over μ^+, compensated by an equal excess of $\bar\nu_\mu$ over ν_μ. On the other hand, the large value $m_\tau = 1777$ MeV completely suppresses the presence of τ leptons. Together with the absence of trapped τ lepton number this implies $\mu_{\nu_\tau} = 0$. The initial thermal distributions are quite different between the μ and τ flavors!

Apart from these initial differences, it was recently recognized that the opacities for ν_μ, $\bar\nu_\mu$, ν_τ and $\bar\nu_\tau$ are all different from each other if nucleon recoil effects and muonic beta reactions are taken into account, leading to a transient build-up of μ and τ lepton number as neutrinos diffuse out of the star [6]. Again, it is important to understand if the μ and τ lepton numbers are locally conserved on the diffusion time scale or if the condition $\mu_{\nu_\mu} = \mu_{\nu_\tau}$ is enforced by flavor-conversion processes.

The SuperKamiokande measurement of the atmospheric neutrino anomaly [1] suggests maximal ν_μ-ν_τ-mixing and δm^2 roughly between 10^{-3} and 10^{-2} eV^2. We now study if these mixing parameters lead to chemical equilibrium between ν_μ and ν_τ on the diffusion time scale of a few seconds.

II. FLAVOR CONVERSION

Chemical equilibrium between charged leptons ℓ and neutrinos ν_ℓ of a given flavor is established by beta processes of the form $\ell^- + p \leftrightarrow n + \nu_\ell$. They are "infinitely fast" relative to all other time scales of interest. Chemical equilibrium between μ- and τ-flavored leptons, on the other hand, is achieved by neutrino oscillations and by collisions which break the coherence of mixed neutrino states. The evolution of the ν_μ-ν_τ-system is best described in terms of the 2×2 density matrices $\rho_{\mathbf{p}}$ for each momentum $\mathbf{p}$. In the flavor basis, the diagonal entries of $\rho_{\mathbf{p}}$ are the usual ν_μ and ν_τ occupation numbers, respectively, while the off-diagonal entries represent phases produced by oscillations. The evolution of $\rho_{\mathbf{p}}$, and of $\bar\rho_{\mathbf{p}}$ for anti-neutrinos, is governed by a generalized Boltzmann collision equation which simultaneously includes the effects of oscillations and collisions [7,8].

For our simple estimates, however, it will suffice to use a "pedestrian" method where the neutrino ensemble is represented by a single mode with momentum $\mathbf{p}$ or energy $E \approx |\mathbf{p}|$. Without collisions, the flavor content of this mode would oscillate forever. Interactions with the medium, however, destroy the coherence between the flavor components, leading eventually to an equal, incoherent mixture. Maximally mixed neutrinos approach this state of flavor equilibrium with a rate [9]

$$\Gamma_{\rm flavor} = \omega_{\rm vac}^2 \, \frac{D}{D^2 + (\delta V)^2}, \tag{1}$$

where

$$\begin{aligned} \omega_{\rm vac} \equiv \frac{\delta m^2}{2E} &= 1.7 \times 10^{-11}\ {\rm eV}\ \Delta_3\, E_{30}^{-1} \\ &= 2.5 \times 10^{4}\ {\rm s}^{-1}\ \Delta_3\, E_{30}^{-1} \end{aligned} \tag{2}$$

is the oscillation frequency in vacuum. Here, $\Delta_3 \equiv \delta m^2/10^{-3}$ eV^2 represents the lower end of the mass range implied by SuperKamiokande and $E_{30} \equiv E/30$ MeV where 30 MeV is a typical temperature in a SN core. Further, δV is the energy difference between ν_μ and ν_τ of equal momenta caused by the medium, i.e. δV is the difference of the medium's weak potential for our two neutrino flavors.

Finally, D is the damping or decoherence rate, i.e. the rate by which interactions with the medium "measure" the flavor content of a mixed neutrino state. Typically D is of the order of the neutrino collision rate, but the exact relationship between the two quantities is not trivial. For example, in a situation where one of the neutrino flavors scatters with a rate $\Gamma_{\rm coll}$ while the other is sterile, one finds $D = \Gamma_{\rm coll}/2$ [9]. Equation (1) applies in the "strong damping limit" defined by $D \gg \omega_{\rm vac}$, a condition which is satisfied in our scenario.

Equation (1) is easily interpreted in two limiting cases. For $\delta V = 0$, the flavor content of a given state oscillates as $\frac{1}{2}[1 + \cos(\omega_{\rm vac} t)]$. The oscillations are interrupted by those collisions which "measure" the difference between ν_μ and ν_τ [7–9]. If this "measurement rate" or "decoherence rate" $D = \tau_D^{-1}$ is much larger than $\omega_{\rm vac}$, the flavor oscillation is interrupted when $\omega_{\rm vac} t \ll 1$ so that the flavor content evolves as $(\omega_{\rm vac} t/2)^2$ until it is interrupted. This happens with a rate τ_D^{-1} so that the rate of flavor conversion must scale as $(\omega_{\rm vac}\tau_D)^2 \tau_D^{-1}$ or

$$\Gamma_{\rm flavor} = \frac{\omega_{\rm vac}^2}{D}. \tag{3}$$

In this case the flavor conversion rate decreases with increasing D, and vanishes for infinite D. This situation is known as the "Quantum Zeno Paradox" or "Watched Pot Effect" [9]: The neutrino remains "frozen" in its flavor state because it is frequently "watched" or measured to be in this state by the interactions with the medium.

A more familiar limiting case obtains when $|\delta V| \gg D$ so that D^2 in the denominator of Eq. (1) can be neglected. Since $D \gg \omega_{\rm vac}$ by assumption, we also have $|\delta V| \gg \omega_{\rm vac}$, implying that the oscillation frequency is δV instead of $\omega_{\rm vac}$ because the energy difference between ν_μ and ν_τ of equal momenta is now dominated by δV, not by $\omega_{\rm vac}$. Since $|\delta V| \gg D$, the collisions are rare relative to the oscillation period.[1] Averaging over an oscillation period, and if we begin with one flavor, the average probability for the appearance of the other is $\frac{1}{2}\sin^2(2\Theta)$ where Θ is the in-medium mixing angle. If the oscillations are interrupted with a rate $\Gamma_{\rm coll}$, we have

$$\Gamma_{\rm flavor} = \sin^2(2\Theta)\, D \tag{4}$$

if D is interpreted as $\Gamma_{\rm coll}/2$. For maximally mixed neutrinos, the in-medium mixing angle is given by

$$\tan(2\Theta) = \frac{\omega_{\rm vac}}{\delta V}. \tag{5}$$

Since $|\omega_{\rm vac}/\delta V| \ll 1$ we have $\tan(2\Theta) \approx \sin(2\Theta)$ so that we recover Eq. (1). A large refractive energy difference between the flavors suppresses the conversion rate, an effect which evidently can be interpreted as a suppression of the mixing angle in the medium.

III. RATE OF DECOHERENCE

In order to estimate the rate of flavor conversion we are thus left with the task of estimating the damping rate, D, and the refractive effect δV for the conditions of a SN core. The largest possible conversion rate obtains for $\delta V = 0$ so that we first consider this case. If $\Gamma_{\rm flavor}$ were slow in this limit, a further discussion of refractive effects would be unnecessary.

We thus begin by comparing $\Gamma_{\rm flavor} = \omega_{\rm vac}^2 D^{-1}$ with the diffusion rate $\tau_{\rm diffusion}^{-1} \approx 1\ {\rm s}^{-1}$,

$$\Gamma_{\rm flavor}\tau_{\rm diffusion} = \frac{6.4\times10^8\ {\rm s}^{-1}}{D}\,\frac{\Delta_3^2}{E_{30}^2}, \tag{6}$$

where we have used Eq. (2). Evidently it is the low-energy neutrino modes with their fast oscillation frequency which are most effective for flavor conversion.

There are two conceptually different contributions to D. Oscillations are interrupted by scattering processes which are sensitive to the neutrino flavor such as $\nu_\mu + e^- \to \nu_e + \mu^-$, a process which has no analogue for ν_τ because of the large τ mass. In addition, a mixed neutrino can scatter to higher energies, for example by $\nu + e^- \to e^- + \nu$, without interrupting the oscillation process. However, after this "up-scattering," the probability for processes which do distinguish between the flavors increases significantly since the cross sections for all neutrino processes increase with energy. Moreover, at large energies the fast beta process $\nu_\mu + n \to p + \mu^-$ becomes kinematically allowed. Effectively, the up-scattering of a low-energy neutrino causes the oscillations to be interrupted, even if the up-scattering process itself is flavor-diagonal. Therefore, up to a numerical factor, the decoherence rate D of low-energy neutrinos is identical with their scattering rate to higher energies.

The most important energy-changing process for neutrinos in a SN core is the scattering on the nuclear medium. Ignoring the subdominant vector-current interaction, the differential cross-section for neutral-current scattering of a neutrino of energy E_1 to energy E_2 is

$$\frac{d\sigma}{dE_2} = \frac{3C_A^2\,G_F^2}{\pi}\,E_2^2\,\frac{S(E_1-E_2)}{2\pi}. \tag{7}$$

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[1] Sometimes this situation is referred to as "weak damping." However, we follow the convention of Ref. [9] where "weak damping" means $D \ll \omega_{\rm vac}$.

Here, $S(\omega)$ is the dynamical structure function for the axial-vector current interaction in the long-wavelength limit [10]. In a dilute medium $S(\omega) = 2\pi\delta(\omega)$, leading to the usual neutral-current elastic scattering cross section. In a nuclear medium, nucleon-nucleon interactions cause $S(\omega)$ to be a broadly smeared-out function which obeys the detailed-balancing condition $S(-\omega) = S(\omega)e^{-\omega/T}$. Unless spin-spin correlations and degeneracy effects are important, the structure function has the norm $\int S(\omega)d\omega/2\pi = 1$.

In contrast with elastic neutrino scattering processes, the cross section Eq. (7) does not vanish for small neutrino energies. With $E_1 = 0$, the neutrino scattering rate on nucleons is

$$\Gamma_{\rm nuc}(0) = \frac{3C_A^2 G_F^2}{\pi}\, T^2\, n_B \int_0^\infty dx\, x^2 \frac{TS(-Tx)}{2\pi} \tag{8}$$

where n_B is the baryon (nucleon) number density. With $C_A = 1.26/2$ for neutral-current processes, the coefficient before the integral is

$$0.98 \times 10^8 \text{ s}^{-1}\, \frac{\rho}{3 \times 10^{14} \text{ g cm}^{-3}}\, T_{30}^2, \tag{9}$$

where $T_{30} \equiv T/30$ MeV.

The dimensionless integral vanishes if $S(\omega)$ is very narrow, corresponding to the usual case of a vanishing elastic neutrino scattering cross section at $E_1 = 0$. The integral also vanishes when it is very broad because $S(\omega)$ decreases at least exponentially for large negative ω. If $S(\omega)$ is normalized, and if it is a smoothly varying broad function, the maximum possible value for the integral expression is about 0.25 which obtains when the width of $S(\omega)$ is of order the temperature, probably corresponding to realistic SN conditions. The rate gets reduced by nucleon degeneracy effects and spin-spin anticorrelations. Therefore, for low-energy neutrinos the upscattering rate is probably not larger than a few times 10^7 s^{-1}. For low-energy neutrinos, this rate is larger than any other up-scattering process that we could identify such as neutrino-electron or neutrino-neutrino scattering.

Below the muon production threshold in beta processes, the most important reactions which distinguish between ν_μ and ν_τ appear to be $\nu_\mu + e^- \to \mu^- + \nu_e$ and $\bar\nu_\mu + \mu^- \to e^- + \bar\nu_e$. We find that for typical conditions in a SN core and for neutrino energies around T the rate is at most a few times 10^7 s^{-1} and thus comparable to neutrino nucleon scattering.

Comparing these numbers with Eq. (2) reveals that $D \gg \omega_{\rm vac}$. Therefore, the strong damping limit indeed applies as had been assumed earlier.

The actual decoherence rate D is smaller than our estimates of the corresponding collision rates. We have already mentioned that D is *half* the collision rate in a situation where one flavor scatters while the other is sterile [9]. Altogether, we believe that

$$D \lesssim 3 \times 10^7 \text{ s}^{-1} \tag{10}$$

is a realistic upper limit for typical SN conditions of $\rho = 3 \times 10^{14}$ g cm^{-3} and $T = 30$ MeV, implying

$$\Gamma_{\rm flavor}\tau_{\rm diffusion} \gtrsim 20\,\Delta_3^2\,E_{30}^{-2}. \tag{11}$$

Averaging this expression over a thermal neutrino distribution, we note that $\langle E^{-2}\rangle = 0.38\,T^{-2}$ so that we finally conclude that for typical SN conditions

$$\langle\Gamma_{\rm flavor}\rangle\tau_{\rm diffusion} \gtrsim 10\,\Delta_3^2. \tag{12}$$

Near the upper end of the mass range suggested by SuperKamiokande we have $\delta m^2 = 10^{-2}$ eV^2 or $\Delta_3 = 10$, implying that the flavor conversion rate increases by two orders of magnitude relative to our estimate.

In summary, it appears that the flavor conversion between ν_μ and ν_τ in a SN core would be much faster than the diffusion time scale of about one second if refractive effects could be ignored.

IV. REFRACTIVE EFFECTS

Turning to neutrino refraction in the SN medium, we begin with the usual lowest-order weak potential,

$$\delta V^{(1)} = \sqrt{2}\,G_F\left(\Delta n_\mu + \Delta n_{\nu_\mu} - \Delta n_\tau - \Delta n_{\nu_\tau}\right), \tag{13}$$

where G_F is the Fermi constant while Δn_j stands for the number density of particle j minus the density of anti-particles. Since the total μ and τ lepton number trapped in a SN core is zero, we have initially $\delta V^{(1)} = 0$. Therefore, the flavor conversion rate would seem to be initially fast until a significant weak potential difference has built up.

However, second-order effects can not be neglected. At one-loop level, the neutral-current interactions of ν_μ and ν_τ with the nucleons are not identical because of the different charged-lepton masses in the loop. The second-order energy difference in a normal medium was found to be [11]

$$\begin{aligned} |\delta V^{(2)}| &= \frac{3G_F^2 m_\tau^2}{2\pi^2}\, n_B \left[\ln\left(\frac{m_W^2}{m_\tau^2}\right) - 1 + \frac{Y_n}{3}\right] \\ &= 6.3\times10^{-4}\ \text{eV}\,\frac{\rho}{3{\times}10^{14}\ \text{g/cm}^3}\,\frac{6.61+Y_n/3}{7} \end{aligned} \tag{14}$$

where n_B is the baryon density and Y_n the neutron number fraction. With 6.3×10^{-4} eV $= 9.6\times10^{11}$ s^{-1} we find $|\delta V^{(2)}| \gg D \gg \omega_{\rm vac}$ and that even initially the flavor conversion time scale (the inverse of $\Gamma_{\rm flavor}$) far exceeds the one for diffusion. This must be the only example where radiative corrections to the neutrino refractive index are of direct practical relevance!

We stress, however, that $\delta V^{(2)}$ is by no means exotically large. Its importance in the present context derives from the cancellation of $\delta V^{(1)}$ and from the smallness of $\omega_{\rm vac}$. For example, the in-medium mixing angle in the present case is

$$\tan(2\Theta) = 2.6\times10^{-8}\,\frac{3{\times}10^{14}\ \text{g/cm}^3}{\rho}\,\frac{7}{6.61+Y_n/3}, \tag{15}$$

i.e. second-order refractive effects suppress the mixing angle by about eight orders of magnitude! The real surprise is that $\omega_{\rm vac}$, despite its smallness, is large enough to cause flavor equilibrium if it were not for the refractive suppression of the mixing angle.

Another second-order contribution to δV arises from the low-energy tail of the $W^\pm$ and Z^0 resonance in neutrino forward scattering on other neutrinos and charged leptons [12]. This term is the dominant second-order correction in the early universe where the particle-antiparticle asymmetries are small, but in a SN core Eq. (14) is more important because of the huge baryon density.

Once diffusion processes begin to build up a net μ or τ lepton number in the SN core, the first-order refractive effect becomes important. Its magnitude is understood if we calculate the refractive energy shift caused by neutrinos with a chemical potential $\mu_\nu \ll T$ which we express as $\eta_\nu = \mu_\nu/T$. To lowest order we have $\Delta n_\nu = T^3\eta_\nu/6 + \mathcal{O}(\eta_\nu^2)$ so that

$$\delta V_\nu^{(1)} = 0.074\ \text{eV}\ T_{30}^3\eta_\nu. \tag{16}$$

Likewise, for muons we find $\Delta n_\mu = (T^3/\pi^2)\,f(m_\mu/T)\,\eta_\mu$ where the function $f(m_\mu/T)$ is 1.077 for $m_\mu/T = 3$. Therefore, muons provide a similar energy shift. Altogether, if we use η as some characteristic chemical potential for the muons and neutrinos, the first-order energy shift is given by Eq. (16), leading to an in-medium mixing angle of

$$\tan(2\Theta) = 2.2\times10^{-10}\,\frac{\Delta_3}{\eta\,E_{30}T_{30}^3}. \tag{17}$$

Again, the mixing angle is hugely suppressed unless η is finely tuned to zero. Considering that $\Gamma_{\rm flavor} = \sin^2(2\Theta)D$ the suppression of the conversion rate is quadratic in the small in-medium mixing angle. For $D = \mathcal{O}(10^7\ \text{s}^{-1})$ even a tiny value for η is enough to suppress flavor conversion completely.

As μ and τ lepton number builds up it may happen that the first and second-order contributions to δV cancel for suitable conditions, allowing briefly for a fast flavor conversion rate. However, the slightest deviation from this condition again suppresses the mixing angle and thus quenches any further flavor conversion. Any attempt to reach flavor equilibrium by neutrino oscillations is robustly self-quenching.

We have mostly studied typical core conditions in the SN. In regions below the neutrino sphere the density is about three orders of magnitude less than our benchmark value of nuclear density, and the temperature may be a factor of 5 smaller than our standard figure of 30 MeV. For such conditions, the presence of muons is strongly suppressed so that the decoherence rate D is much smaller than what we have estimated for average core conditions. On the other hand, the in-medium mixing-angle Eq. (17) is still very small if we use $T_{30} = 0.2$ and $E_{30} = 0.2$ as characteristic values near the neutrino sphere. Therefore, our conclusion that flavor equilibrium cannot be achieved applies to conditions throughout the SN core.

V. CONCLUSIONS

The chemical equilibration between the μ and τ flavors in a SN core by neutrino oscillations would be fast on the neutrino diffusion time scale if the refractive energy shift δV between ν_μ and ν_τ were small. The usual first-order contribution indeed vanishes due to the lack of trapped μ- and τ-lepton number, but second-order contributions are large enough to suppress flavor conversion.

In the course of neutrino transport to the stellar surface, μ- and τ-lepton number will build up due to the differences between the ν_μ, $\bar{\nu}_\mu$, ν_τ, $\bar{\nu}_\tau$ diffusion constants [6]. It is conceivable that local conditions can be reached where the refractive energy shift δV vanishes to all orders. While this cancellation would momentarily lead to a fast rate of flavor conversion, the resulting re-distribution of μ and τ lepton number quickly produces a δV so large that the conversion rate is suppressed again. The equilibrium condition $\mu_{\nu_\mu} = \mu_{\nu_\tau}$ implies a δV so large that it can never be reached, i.e. the flavor conversion process is inevitably self-quenching.

For the small neutrino mass differences indicated by the experiments, the lack of conversion between ν_e and the other flavors is much easier to understand because the large amount of trapped electron-lepton number causes $\delta V^{(1)}$ to be so large that the in-medium mixing angle is easily seen to be vastly suppressed [5].

The phenomenon of neutrino mixing and neutrino oscillations may have a variety of astrophysical consequences which need to be explored. In the past, most of the attention was focussed on the possibility of detecting evidence for neutrino oscillations in the astrophysical context. However, as neutrino oscillations become more and more experimentally established, the problem of unravelling the neutrino mass matrix may become a less pressing astrophysical preoccupation than the reverse question: given the experimentally measured neutrino parameters, is the effect of flavor violation important or ignorable in a given environment?

We think it is intriguing that the core of a SN is protected from the consequences of flavor-lepton number violation by the phenomenon of neutrino refraction and that second-order effects have to be included. While the significance of the weak-interaction potential for the resonant enhancement of oscillations in the spirit of the MSW effect has been widely acknowledged, the suppression of oscillations in a SN core is another consequence of neutrino refraction with real and important astrophysical ramifications.

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