On finite–temperature and –density radiative corrections to the neutrino effective potential in the early Universe

R.Horvat

“Rugjer Boskovic” Institute, P.O.Box 1016, 10001 Zagreb, Croatia.

Abstract

Finite–temperature and –density radiative corrections to the neutrino effective potential in the otherwise $CP$–symmetric early Universe are considered in the real–time approach of Thermal Field Theory. A consistent perturbation theory endowed with the hard thermal loop resummation techniques developed by Braaten and Pisarski is applied. Special attention is focused on the question whether such corrections can generate any nonzero contribution to the $CP$–symmetric part of the neutrino potential, if the contact approximation for the $W$–propagator is used.

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For neutrinos propagating in matter, the vacuum energy–momentum relation is no longer respected. The modification of the neutrino dispersion relation can be represented in terms of an index of refraction [1] or an effective potential [2]. However, in the framework of Thermal Field Theory (TFT) it is based on the real part of the matter–induced neutrino self–energy [3].

The subject of neutrino propagation in matter became popular firstly when Wolfenstein calculated neutrino refractive index in matter [4]. Later on, Mikheyev and Smirnov recognized resonant nature of flavor oscillations triggered by matter effects [5]. The hypothesis based on these effects has even become the most popular explanation for the solar–neutrino deficit [6].

Neutrino oscillations could also be cosmologically important. The oscillations between a standard left–handed neutrino (ν_e, ν_μ or ν_τ) and a SU(2)_L singlet ν_S has attracted considerable attention recently [7]. Oscillations into new neutrino states would distort successful Big Bang Nucleosynthesis (BBN) and hence a bound on oscillations parameters can be derived. The constraints on mixing between active and sterile neutrinos are such that excluded regions include the large angle ν_e − ν_S MSW solution [8] as well as the ν_μ − ν_S mixing solution [9] to the atmospheric neutrino problem, provided that there was no significant CP asymmetry in the early Universe. In addition, for Dirac neutrinos endowed with anomalous magnetic dipole moments, the magnetically induced neutrino oscillations between left–handed (ν_L) and right–handed (ν_R) states may occur [10]. In such oscillations the BBN arguments constrain the product of neutrino MDM’s and a present–day intergalactic field strength.

In all above situations in the early Universe (and also in other environments) the neutrino refractive effects are of crucial importance. In the zero–temperature and –density (0TD) limit, it suffices to calculate an index of refraction, obtained directly from the neutrino forward–scattering amplitude. For lowest order calculations at finite–temperature and –density (FTD), one usually uses the old version of TFT formulated in the real time, as developed first by Dolan and Jackiw [11] (in the real time formalism it is more convenient to separate the vacuum effects from the effects of the medium). However, this old version is
not applicable to higher-order calculations at FTD, since it is plagued with ill-defined pinch singularities which do not cancel. Instead, one should use a consistent TFT perturbative approach initiated by Niemi and Semenoff \cite{12}, and amplified later on by Braaten and Pisarski \cite{13} through the resummation program for the soft regime.

In terms of standard electroweak interactions the almost $CP$ symmetry of the medium implies that the lowest order refractive effects of order $G_FT^3$, which dominate in stars, nearly cancel. However, the $CP$–symmetric contribution, which arises from an expansion of gauge–boson propagators, turns out to be suppressed by a very small factor of order $T^2/M_W^2$, where the temperature $T$ in the epoch of interest $T \sim 0.1 - 10$ MeV. Since the only nonzero contribution is of order $G_FT^5$, the smallness of such tree–level result calls for an investigation of its radiative corrections. It was found \cite{14} that the 0TD radiative correction to the neutrino index of refraction in the early Universe is about 20% for $\nu_e$ and 50% for $\nu_\mu$ and $\nu_\tau$. The FTD radiative corrections have not been considered yet, except the general proof for cancellation of infrared and mass (or collinear) singularities at order $\alpha$, as given recently in \cite{15}. In the present note, we shall investigate FTD radiative corrections to the lowest order result (the zeroth order in the expansion of the $W$–propagator, i.e., the contact approximation), to see if they could produce any term not proportional to the particle–antiparticle asymmetry, namely of order $\alpha G_FT^3$. The main difference between the 0TD and FTD radiative corrections is the appearance of thermal ($e^\pm, \gamma$) loops in the latter, characterized by circulation of real particles from the heat bath. Thus, for example, the FTD radiative correction to the process $\nu_e e^\pm \rightarrow \nu_e e^\mp$ brings in some extra processes like $\nu_e e^\mp \gamma \rightarrow \nu_e e^\pm \gamma$, $\nu_e \gamma \rightarrow \nu_e \gamma$, where $\gamma'$s are from the heat bath.

To begin, let us remind how the cancellation in the lowest order contribution results. We will be always working in the limit of perfect $CP$ symmetry, which, as a consequence, results in the equal amounts of particles and antiparticles in the early Universe (stated differently, the chemical potential $\mu = 0$ is assumed for all the species). Then, the standard bubble charged–current graph for $\nu_e$ gives the following contribution to the neutrino self–energy in the contact approximation,
\[ \Sigma_{st.}^{(W)} = 2\sqrt{2} G_F \int \frac{d^4 k}{(2\pi)^4} \gamma^\mu Li S_{11}(k) \gamma_\mu L , \]  

where the (11) component of the real–time electron propagator is given by

\[ i S_{11}(k) = (\not k + m_e) \left[ \frac{i}{k^2 - m_e^2 + i\varepsilon} - 2\pi \delta(k^2 - m_e^2) \sin^2 \phi_k \right] , \]

and

\[ \sin^2 \phi_k = \frac{1}{e^{\beta |k_0|} + 1} . \]

is the Fermi–Dirac distribution function. Insertion of the FTD part of (2) into (1) gives

\[ \Sigma_{st.}^{(W)} = 4\sqrt{2} G_F \left\{ \int \frac{d^3 k}{(2\pi)^3} \int_{-\infty}^{+\infty} dk_0 k_0 \delta(k_0^2 - k^2 - m_e^2) \frac{1}{e^{\beta |k_0|} + 1} \right\} \gamma_0 L , \]

where \( k \equiv |\vec{k}| \). In Eqs.(3) and (4), the four–velocity of the center of mass of the medium \( u^\mu \) is given by its value in the rest frame of the medium; \( u^\mu = (1, \vec{0}) \). The effective potential is obtained from a neutrino self–energy by omitting the Dirac part \( \not \! k L \). In a mathematical sense, it is easy to understand why Eq.(4) vanishes identically: We simply integrate the odd function of \( k_0 \) over the symmetric interval. Now, we are going to check up whether a mathematical feature like that survives FTD radiative corrections.

The correction to the first order (in \( G_F \)) FTD contribution to the self–mass of \( \nu_e \) due to photon radiation amounts to calculate the \( O(\alpha) \) correction to the bare electron propagator. This is performed using the Schwinger–Dyson equation for the full fermion propagator. In our analysis, for simplicity, the electron is always considered to be massless, i.e., \( m_e = 0 \), a standard assumption in the early Universe calculations around the BBN epoch. Summing over different internal vertices and using the “momentum–derivator formula”, the final result in the real–time framework can be written in the following form

\[ i \delta S_{11}^M(k) = \lim_{k^2 \to k^2} \left\{ \pi \frac{\partial}{\partial k^2} \delta(k_0^2 - \vec{k}^2) \right\} \! k \text{Re} \! \Sigma_M(k_0, \vec{k}) \! k - PP \frac{1}{(k_0^2 - \vec{k}^2)^2} \! k \text{Im} \! \Sigma_M(k_0, \vec{k}) \! k \right\} \]

\[ - 2 \lim_{k^2 \to k^2} \left\{ \pi \frac{\partial}{\partial k^2} \delta(k_0^2 - \vec{k}^2) \right\} \! k \left[ \text{Re} \! \Sigma_0(k_0, \vec{k}) + \text{Re} \! \Sigma_M(k_0, \vec{k}) \right] \! k \]

\[ - PP \frac{1}{(k_0^2 - \vec{k}^2)^2} \! k \left[ \text{Im} \! \Sigma_0(k_0, \vec{k}) + \text{Im} \! \Sigma_M(k_0, \vec{k}) \right] \! k \right\} \sin^2 \phi_k \]

\[ + \lim_{k^2 \to k^2} \left\{ i \pi \frac{\partial}{\partial k^2} \delta(k_0^2 - \vec{k}^2) \right\} \! k \text{Im} \! \Sigma_M(k_0, \vec{k}) \! k - PP \frac{i}{(k_0^2 - \vec{k}^2)^2} \! k \text{Re} \! \Sigma_M(k_0, \vec{k}) \! k \right\} . \]
Since it is assumed that $m_e = 0$, the well known “mass–derivative formula” \[16\] has to be replaced with the “momentum–derivative formula”. Since we want to discuss the effects of the medium, only FTD parts are kept in (5). Also, note that the terms in the third curly bracket in (5) do not contribute to the real part of the neutrino self–energy. Moreover, we have split in (5) the electron self–energy function $\Sigma$ into the vacuum and the matter part. The parts associated to $\text{Re}\Sigma_0$, $\text{Re}\Sigma_M$ contain virtual photon corrections (at 0TD and FTD, respectively) and correspond to a cut through the internal line. On the other hand, the parts involving $\text{Im}\Sigma_0$, $\text{Im}\Sigma_M$ correspond to a cut through $\Sigma$ and are associated to a real process. Finally, Eq.(5) is understood as integrating over the $\delta$–function before taking the derivative.

If $m_e = 0$, the electron mass operator can be decomposed as

$$\Sigma(k_0, \vec{k}) = a(k_0, k)\gamma^0 + b(k_0, k)\vec{k} \cdot \vec{\gamma}. \quad (6)$$

Then, with the aid of generalized Cutosky rules at FTD \[17\], we obtain for the imaginary parts,

$$\text{Im}[a_0(k_0, k) + a_M(k_0, k)] = (4\pi\alpha\varepsilon(k_0) / \sin 2\phi_k) \times \begin{cases} \int_{\frac{1}{2}(k_0+k)}^{\frac{1}{2}(k_0-k)} \frac{xdx\Theta(x)}{32\pi k} \frac{\varepsilon(k_0 - x)}{\sinh \left[ \frac{1}{2}\beta(k_0 - x) \right] \cosh \left( \frac{1}{2}\beta x \right)} \\ + \int_{\frac{1}{2}(k_0-k)}^{-\frac{1}{2}(k_0+k)} \frac{xdx\Theta(x)}{32\pi k} \frac{\varepsilon(k_0 + x)}{\sinh \left[ \frac{1}{2}\beta(k_0 + x) \right] \cosh \left( \frac{1}{2}\beta x \right)} \end{cases} \quad , \quad (7a)$$

$$\text{Im}[b_0(k_0, k) + b_M(k_0, k)] = (4\pi\alpha\varepsilon(k_0) / \sin 2\phi_k) \times \begin{cases} -\frac{k_0}{k^2} \left[ \int_{\frac{1}{2}(k_0+k)}^{\frac{1}{2}(k_0-k)} \frac{xdx\Theta(x)}{32\pi k} \frac{\varepsilon(k_0 - x)}{\sinh \left[ \frac{1}{2}\beta(k_0 - x) \right] \cosh \left( \frac{1}{2}\beta x \right)} \right] \\ + \int_{\frac{1}{2}(k_0-k)}^{-\frac{1}{2}(k_0+k)} \frac{xdx\Theta(x)}{32\pi k} \frac{\varepsilon(k_0 + x)}{\sinh \left[ \frac{1}{2}\beta(k_0 + x) \right] \cosh \left( \frac{1}{2}\beta x \right)} \right] \\ + \left( k_0^2/2k^2 - 1/2 \right) \left[ \int_{\frac{1}{2}(k_0+k)}^{\frac{1}{2}(k_0-k)} \frac{dx\Theta(x)}{32\pi k} \frac{\varepsilon(k_0 - x)}{\sinh \left[ \frac{1}{2}\beta(k_0 - x) \right] \cosh \left( \frac{1}{2}\beta x \right)} \right] \\ - \int_{\frac{1}{2}(k_0-k)}^{-\frac{1}{2}(k_0+k)} \frac{xdx\Theta(x)}{32\pi k} \frac{\varepsilon(k_0 + x)}{\sinh \left[ \frac{1}{2}\beta(k_0 + x) \right] \cosh \left( \frac{1}{2}\beta x \right)} \right] \end{cases} \quad . \quad (7b)$$
Also, trivially, we have

\[ \text{Im} a_0(k_0, k) = -\frac{\alpha}{4} k_0 \Theta(k_0^2 - k^2), \]  
(8a)

\[ \text{Im} b_0(k_0, k) = \frac{\alpha}{4} \Theta(k_0^2 - k^2). \]  
(8b)

The easiest way to account for the terms involving \( \text{Re} \Sigma \) is by writing down a dispersion relation for the matter part of the self–energy at both poles,

\[ \text{Re} a_M(\pm k_0, k) = \pm \frac{PP}{\pi} \int_{-\infty}^{+\infty} dk'_0 \frac{\text{Im} a_M(k'_0, k)}{k'_0 \mp k_0}, \]  
(9)

and the same for \( \text{Re} b_M(\pm k_0, k). \) Eq.(9) is based on the fact that \( \Sigma_M \sim \exp(-\beta |p_0|) \) as \( |p_0| \to \infty. \) Gathering all together, one finds, after performing angular integrations, that the FTD radiative correction at \( O(\alpha) \) to the standard result (2) is obtained by the replacement:

\[ \int_{-\infty}^{+\infty} dk_0 k_0 \delta(k_0^2 - k^2) \left\{ \int_{-\infty}^{+\infty} dk'_0 \left[ \frac{\partial}{\partial k^2} \delta(k_0^2 - k)^2 \right] (k_0^2 + k^2) \text{Re} a_M(k_0, k) \
+ 2k_0 k^2 \text{Re} b_M(k_0, k) \right\} \left( -\frac{1}{2} + \frac{1}{e^{\beta|p_0|} + 1} \right) 
+ \frac{1}{2\pi} PP \frac{1}{(k_0^2 - k^2)^2} \left[ (k_0^2 + k^2) \text{Im} a_M(k_0, k) + 2k_0 k^2 \text{Im} b_M(k_0, k) \right] 
- \frac{1}{\pi} PP \frac{1}{(k_0^2 - k^2)^2} \left[ (k_0^2 + k^2) \text{Im}[a_0(k_0, k) + a_M(k_0, k)] \
+ 2k_0 k^2 \text{Im}[b_0(k_0, k) + b_M(k_0, k)] \right] \frac{1}{e^{\beta|p_0|} + 1} \right\} . \]  
(10)

By a direct inspection of the above expressions, one concludes that

\[ \text{Im}[a_0(-k_0, k) + a_M(-k_0, k)] = -\text{Im}[a_0(k_0, k) + a_M(k_0, k)], \]  
(11a)

\[ \text{Im}[b_0(-k_0, k) + b_M(-k_0, k)] = -\text{Im}[b_0(k_0, k) + b_M(k_0, k)], \]  
(11b)

\[ \text{Re} a_M(-k_0, k) = -\text{Re} a_M(k_0, k), \]  
(11c)

\(^1\)Within the above formalism, it can be easily seen that the term involving the renormalized part of \( \text{Re} \Sigma_0 \) vanishes, in agreement with the results obtained earlier \([17]\).
Hence, upon inclusion of the FTD radiative corrections at $\mathcal{O}(\alpha)$, one finds that one odd function of $k_0$ is replaced by another odd function of $k_0$, and the net result is again zero.

However, strictly speaking, the above analysis is adequate only for hard loop momenta of order $T(k_0, k \sim T)$, whereas for soft momenta $(k_0, k \sim eT)$ the resummation program developed by Pisarski [13], Braaten and Pisarski [13] and Frenkel and Taylor [20] should be applied. The starting point is that, in the sense of “hard thermal loop” [13,20], one has to make a distinction between hard loop momenta of order $T$ and soft momenta of order $eT$. The thermal mass of the electron, being of order $eT$, is generated by a loop integral where the momentum running inside the loop is hard. The hard momentum contribution to the thermal self–energy of the electron is called a hard thermal loop [13,20]. Finally, only soft lines need to be resummed (the HTL resummed propagators are used), whereas for hard lines the bare perturbation series can still be used. Usually, the arbitrary intermediate energy–momentum cut–off $k_c$ of order $\sqrt{eT}$ is put by hand to separate the two regimes (the final result should be independent of $k_c$ [21]).

The matter part of the resummed electron propagator is given by (only the real part of it is kept)

$$iS^{M,R}_{11}(k_0, \vec{k}) = (1 - 2 \sin^2 \phi_k) \text{Re} \left( \frac{i}{\vec{k} - \Sigma(k_0, \vec{k}) + i\varepsilon} \right).$$

With the explicit expressions for the fermionic HTL (first determined by Klimov [22] and Weldon [23]), it takes the form

$$iS^{M,R}_{11} = \left( -\frac{1}{2} + \sin^2 \phi_k \right) \left\{ (\gamma_0 + \vec{\gamma} \cdot \hat{k}) \text{Im} \frac{1}{A_0 + A_S} + (\gamma_0 - \vec{\gamma} \cdot \hat{k}) \text{Im} \frac{1}{A_0 - A_S} \right\},$$

where the two functions with opposite chirality/helicity ratio are given by

$$\frac{1}{A_0 \mp A_S} = \frac{1}{k_0 \mp k - \frac{m_i^2}{2k}} \left[ \frac{1}{1 \mp \frac{k_0}{k}} \ln \frac{k_0 \mp k}{k_0 \pm k} \pm 2 \right],$$

and $m_i^2 = e^2T^2/8$ is the thermal mass of the electron. Note, that there are four poles in $S^{M,R}_{11}$ [24]. Hence, for soft lines, the replacement (10) is actually given by
\[
\int_{-k_c}^{k_c} dk_0 k_0 \delta (k_0^2 - k^2) \frac{1}{e^{\beta |k_0|} + 1} \rightarrow \int_{-k_c}^{k_c} \frac{dk_0}{2\pi} \left( \text{Im} \frac{1}{A_0 + A_S} + \text{Im} \frac{1}{A_0 - A_S} \right) \left( \frac{1}{2} - \frac{1}{e^{\beta |k_0|} + 1} \right) .
\]

(15)

Above the light cone, the renormalized propagator is determined by quasiparticles which are collective excitations. We find for \( k_0 > k_0 > 0 \),

\[
\text{Im} \frac{1}{A_0 + A_S + i\epsilon} + \text{Im} \frac{1}{A_0 - A_S + i\epsilon} = (\pi/2m^2_e) \left[ \delta (k_0 - \omega_- (k)) (\omega_-^2 (k) - k^2) + \delta (k_0 - \omega_+ (k)) (\omega_+^2 (k) - k^2) \right] ,
\]

(16)

and for \(|k_0| > k_0 < 0\),

\[
\text{Im} \frac{1}{A_0 + A_S - i\epsilon} + \text{Im} \frac{1}{A_0 - A_S - i\epsilon} = (\pi/2m^2_e) \left[ \delta (k_0 + \omega_- (k)) (\omega_-^2 (k) - k^2) + \delta (k_0 + \omega_+ (k)) (\omega_+^2 (k) - k^2) \right] ,
\]

(17)

where \( \omega_{\pm} \) denote the two dispersion laws \[24\].

Instead, below the light cone, the renormalized propagator is determined by the imaginary part of the HTL, which is nonzero owing to the Landau damping mechanism \[13,19\]. Thus, we find for \(|k_0| < k\),

\[
\text{Im} \frac{1}{A_0 - A_S + i\epsilon} = -\frac{\pi (m^2_e/2k) \left( 1 - \frac{k_0}{k} \right) \varepsilon (k_0)}{\left( k_0 - k - m^2_e/2k \left[ \left( 1 - \frac{k_0}{k} \right) \ln \left| \frac{k_0 + k}{k_0 - k} \right| + 2 \right] \right)^2 + m^2_e/4k^2 \left( 1 - \frac{k_0}{k} \right)^2 \pi^2} ,
\]

(18a)

\[
\text{Im} \frac{1}{A_0 + A_S + i\epsilon} = -\frac{\pi (m^2_e/2k) \left( 1 + \frac{k_0}{k} \right) \varepsilon (k_0)}{\left( k_0 + k - m^2_e/2k \left[ \left( 1 + \frac{k_0}{k} \right) \ln \left| \frac{k_0 + k}{k_0 - k} \right| - 2 \right] \right)^2 + m^2_e/4k^2 \left( 1 + \frac{k_0}{k} \right)^2 \pi^2} .
\]

(18b)

By noting that

\[
\text{Im} \frac{1}{A_0 \mp A_S + i\epsilon} (-k_0, k) = -\text{Im} \frac{1}{A_0 \mp A_S + i\epsilon} (k_0, k) ,
\]

(19)

we can see that the new function is again odd with respect to \( k_0 \), and therefore its contribution vanishes.
In conclusion, we have considered previously ignored higher–order corrections to the neutrino effective potential in the early Universe, namely the FTD radiative corrections at $O(\alpha)$, in a theory where only the contact part of the $W$–boson propagator is kept. Searching for $\alpha G_F T^3$ corrections, we have applied a consistent TFT in the real time. The FTD radiative corrections comprise, beside the usual virtual corrections at 0TD, the virtual corrections at FTD as well as the real corrections. We have found that they share the same feature as the lowest order result: They vanish in a CP–symmetric plasma. Because of the resummed character of a perturbation theory, it is easy to show that the same feature persists to all–order perturbation contributions.

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