Chapter 7

Relativistic Description of Nuclear Matter in Hartree Fock using Generalized Hybrid Derivative Coupling

7.1 Introduction

Relativistic nuclear structure calculation achieved considerable success in recent years. The relativistic mean field (RMF) approach has yielded reasonable descriptions of various ground state properties including binding energy, deformation, and charge radius through out the periodic table. The correct spin orbit splitting is obtained naturally from the Lagrangian [1, 2]. Another class of relativistic models is the derivative coupling model introduced by Zimanyi and Moszkowski [4], where a derivative coupling between the scalar meson and the nucleon was employed. This provided, in a natural way, a correct value of the bulk modulus in nuclear matter. However, one important problem with this model is that it predicted a high value of the effective nucleon mass and, consequently, a lower spin-orbit splitting. A better description was obtained using both direct and derivative coupling between the scalar meson and the nucleon [5, 6]. This remedied the situation to a certain extent and provided a better description of the ground state of finite spherical nuclei [7, 8]. This model was successfully applied to deformed nuclei also, particularly in the light mass region as described in chapters 3 and 4.

A lot of features still remain to be studied in hybrid derivative coupling model. Hartree level calculations were performed for nuclear matter, neutron matter and spherical closed and open shell nuclei as well as for deformed nuclei [6, 7, 8] and chapters 3 and 4. A nucleus consists of identical fermions and ex-
change terms, the so called Fock terms, are expected to contribute significantly to the different properties. It has been suggested that the inclusion of Fock term results in a renormalization of the free parameters without improving the description \[2\]. However, the exchange effects included in the mean field level through the Fock terms are important and may substantially modify some results. For example, the liquid gas phase transition in neutron matter observed in Hartree level in $\sigma - \omega$-model is absent if one takes into account exchange terms \[1\]. Recent studies in finite nuclei, particularly in nuclei with extreme neutron proton ratio, indicate the importance of isospin dependent terms. Thus a study of the effect of the isovector mesons, the pion and the $\rho$-meson, may yield interesting results. Hartree Fock calculations allow nonzero contribution of the pion. The pion, being a pseudoscalar meson, cannot contribute in the Hartree level. The effect of two pion correlated exchange, an important contribution as suggested by the meson exchange theories for nucleon-nucleon scattering, is simulated by the exchange of the sigma meson. The isovector-vector rho meson also contributes to symmetric nuclear matter in the Fock level. Tensor interaction mediated by the $\rho$ meson may be included at this level. The Fock terms contribute significantly to symmetry energy. This in turn reduces the rho vector coupling from an unphysically high value required in the relativistic Hartree approximation to a level consistent with the one boson exchange potential (OBE) \[9\]. The effect of the exchange terms in the conventional RMF approach was studied in \[10, 11\]. The same effect was studied for the Zimanyi-Moszkowski Lagrangian and the Glendenning Lagrangian \[12\] also. We propose to study these effects in nuclear matter and neutron matter with generalized hybrid derivative coupling for the $\sigma - N$ interaction. One feature that is expected is a decrease in the effective nucleon mass. The class of derivative coupling models have a problem in describing the spin-orbit splitting which is predicted to be too low. The Fock terms tend to lower the effective mass and hence increase the spin-orbit splitting.

In section 7.2 we outline the theory. In section 7.3 our results for nuclear matter and neutron matter are presented. Also included are the results for the mass and radius of neutron stars obtained for the present Lagrangian.

7.2 Theory

Adding the contributions coming from the $\pi$ meson and the tensor and the vector-tensor terms from the $\rho$ meson, the starting point is the following scaled Lagrangian

$$\mathcal{L} = \bar{\psi} (i \gamma_\mu \partial^\mu - M^* - g_{\rho} \omega^\mu \gamma_\mu - g_{\rho} \gamma^\mu T\cdot \rho_\mu) \psi$$
The symbols have their usual meaning as in [6]. The effective nucleon mass $M^*$ given by

$$\frac{1}{2} \left( \partial_\mu \sigma \partial^\mu \sigma - m_\sigma^2 \sigma^2 \right) - \frac{1}{4} \omega_\mu \omega^{\mu \nu} \frac{1}{4} \rho_\mu \rho^{\mu \nu} + \frac{1}{2} m_\rho^2 \rho^\mu \rho^\mu + L'. \quad (7.1)$$

The new terms are included in $L'$ and are given by

$$L' = \frac{1}{2} \left( \partial_\mu \pi \cdot \partial^\mu \pi - m_\pi^2 \pi^2 \right) + \frac{f_\pi}{2M} \bar{\psi} \sigma_{\mu \nu} \partial^\nu \rho^\mu \cdot \tau \psi + L_{\pi N} \quad (7.3)$$

The $\pi$-$N$ coupling may be pseudoscalar or pseudovector. Both of them give identical results on shell if the coupling constants satisfy certain relations. However, it is well known that chiral symmetry consideration and pair suppression mechanism in NN potential suggest that pseudovector coupling gives a better result in the one pion exchange approximation. Hence, following [10], we have assumed pseudovector coupling

$$L_{\pi N} = -\frac{f_\pi}{m_\pi} \bar{\psi} \gamma_\mu \gamma_5 \partial^\mu \pi \cdot \tau \psi. \quad (7.4)$$

In the equivalent Lagrangian the term proportional to $M^*$ can be written in the following form

$$\bar{\psi} M^* \psi = \bar{\psi} (M + \hat{m} g_\sigma \sigma) \psi = \bar{\psi} M \psi + L_{\sigma N}, \quad (7.5)$$

where

$$\hat{m} = \frac{1}{1 + a g_3 \sigma / M}. \quad (7.6)$$

From Eq. 7.6 it is clear that the quantity $\hat{m} g_3 \sigma$ will contribute to the scalar self energy.

The sigma field equation is

$$(\partial_\mu \partial^\mu + m_\sigma^2) \sigma = g_\sigma \hat{m}^2 \bar{\psi} \psi. \quad (7.7)$$

Because of the sigma-dependence of the $\sigma$-$N$ coupling, the $\sigma$ operator is very complicated. Following the methods followed in [11, 12], the scalar field equation is linearized with respect to the sigma operator, approximating the $\sigma$ operator in $\hat{m}$ by its ground state expectation value $\sigma_0$, we have

$$(\partial_\mu \partial^\mu + m_\sigma^2) \sigma = g_\sigma \hat{m}^2 \bar{\psi} \psi, \quad (7.8)$$
where
\[ \tilde{m} = \frac{1}{1 + \frac{a g_\sigma \sigma_0}{M}}. \] (7.9)

The linearization of the sigma field equation allows us to write the sigma operator as
\[ \sigma(x) = g_\sigma \tilde{m}^2 \int S_c(x,y) \bar{\psi}(y)\psi(y) dy, \] (7.10)
where the propagator \( S_c(x,y) \) satisfies the equation
\[ \left( \partial_\mu \partial^\mu + m_\sigma^2 \right) S_c(x,y) = \delta(x-y). \] (7.11)

The equations for the other mesons remain unchanged. The \( \sigma-N \) interaction term in the Lagrangian, \( L_{\sigma N} \), is also replaced by its linearized form obtained through replacing \( \tilde{m} \) by \( m \).

The Hamiltonian operator is obtained using the general Legendre transformation and is
\[ \mathcal{H} = \int \left[ \bar{\psi}(x)(-i\gamma \cdot \vec{V} + M)\psi(x) + c.c. \right] d^3x + \frac{i}{2} \sum_\Gamma \int \psi(x)\bar{\psi}(y) \Gamma_\Gamma(1,2) S_c(x,y) \bar{\psi}(y)\psi(x) dy d^3x. \] (7.12)

The operators \( \Gamma_\Gamma(1,2) \) are formally equivalent to that of [12] and are given in Table 7.2.

<table>
<thead>
<tr>
<th>( i )</th>
<th>( \Gamma_\Gamma )</th>
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<tbody>
<tr>
<td>( \sigma )</td>
<td>( g_\sigma^2 \tilde{m}^3(2 + \tilde{m}) )</td>
</tr>
<tr>
<td>( \omega )</td>
<td>( g_\omega^2 \gamma_\mu(1)\gamma^\nu(2) )</td>
</tr>
<tr>
<td>( \pi )</td>
<td>( -\frac{(f_\pi/m_\pi)^2}{2}(q\gamma_\mu)(q\gamma_\nu) )</td>
</tr>
<tr>
<td>( \rho )</td>
<td>( g_\rho^2 \gamma_\mu(1)\gamma^\nu(2) )</td>
</tr>
<tr>
<td>( \rho_T )</td>
<td>( i(f_\rho/2M)^2 q_\mu q_\nu \sigma^{\mu\nu}(1)\sigma_{\rho\lambda}(2) )</td>
</tr>
<tr>
<td>( \rho_{VT} )</td>
<td>( i(f_\rho/2M)[\gamma_\mu(2)\sigma^{\mu\nu}(1)q_\nu - \sigma^{\mu\nu}(2)q_\nu\gamma_\mu(1)] )</td>
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</table>

The Lagrangian in Eq. 7.1 can be shown, in the nonrelativistic limit, to give rise to repulsive contact interaction which should be suppressed in real nuclei due to short range correlations. To simulate the effect of short-range correlations, following [10, 11], we have subtracted the zero-rank tensor part of the NN potential coming from \( \pi \) and \( \rho \) exchanges given by
\[
\delta[\Gamma_{\pi,\sigma}(1,2)] = \frac{1}{4\pi q^2} \Gamma_{\pi,\sigma}(1,2).
\] (7.13)

Zero range interaction appears in \( \sigma \)-\( N \) and \( \omega \)-\( N \) interactions also. In the present chapter the coupling constants for \( \pi \) and \( \rho \) mesons are taken from meson exchange theories while that of \( \sigma \) and \( \omega \) are adjusted to reproduce nuclear saturation properties. Thus, the effects of short term correlation are expected to be absorbed in the fitted values of the coupling constants of the two mesons to the nucleon. This was observed in [12] where it was pointed out that fitting the nuclear matter properties to obtain the different parameters, one gets results almost identical with the case where the contact interaction has been explicitly subtracted in the case of \( \sigma \) and \( \omega \) interaction. We neglect the time dependence of the meson fields. This approximation does not affect the direct terms and, in the exchange picture, neglects the retardation effects. It is a very good approximation for the heavier mesons but comparatively poorer for the pion.

The energy is calculated from the Hamiltonian in the tree approximation. The details of the calculation is available in the literature. [10, 11] We give only the necessary expressions.

The kinetic energy is given by

\[
<T> = \frac{2}{\pi^2} \int_0^{p_p} p^2 dp \left( \hat{P}(p) + \hat{M}(p) \right).
\] (7.14)

The expressions \( \hat{P}(p) \) and \( \hat{M}(p) \) are defined as

\[
\hat{M}(p) = \frac{M^*}{E^*},
\] (7.15)

\[
\hat{P}(p) = \frac{P^*}{E^*},
\] (7.16)

where

\[
M^*(p) = M + \Sigma_s(p),
\] (7.17)

\[
P^*(p) = p + \Sigma_s(p),
\] (7.18)

\[
E^*(p) = E(p) - \Sigma_\sigma(p),
\] (7.19)

\[
E^*(p) = \sqrt{p^*^2 + M^*^2}.
\] (7.20)

Here, \( \Sigma \)'s are the contributions of meson exchange toward baryon self energy.

The potential energy density may be split up into a direct and an exchange parts.[10, 12]

\[
<V> = <V_D> + <V_E>.
\] (7.21)
The direct or Hartree term is

\[
< V_D > = \frac{1}{2} \left( \frac{g_x}{m_\sigma} \right)^2 \bar{m}^3 (2 + \bar{m}) \rho_\sigma^2 + \frac{1}{2} \left( \frac{g_x}{m_\omega} \right)^2 \rho_\omega^2 \\
+ \frac{1}{2} \left( \frac{g_x}{m_\rho} \right)^2 (\rho_{B(u)} - \rho_{B(d)})^2.
\]  
(7.22)

The exchange part is more complicated and is given in [10, 12]

\[
< V_E > = \frac{1}{2(2\pi)^4} \int_0^{\pi} dp dp' dp'' \left[ \sum_i A_i(p, p') + M(p) M(p') \sum_i B_i(p, p') + \hat{P}(p) \hat{P}(p') \sum_i C_i(p, p') + \hat{P}(p) M(p') D(p, p') \right].
\]  
(7.23)

The summations run over all the terms in Table 7.2. The quantities \( A_i, B_i, \) and \( C_i \) for the different interactions are tabulated in Table 7.2. The expressions include subtraction of contact interaction for \( \rho \) and \( \pi \) mesons. The only nonzero \( D \) occurs for the vector-tensor coupling and is given by

\[
D(p, p') = \frac{3g_x^2 g_\rho}{2M} (\rho_\theta - 2\rho'\phi). 
\]  
(7.24)

The functions \( \theta \) and \( \phi \) are given by

\[
\theta_i(m_i, p, p') = \ln \frac{m_i^2 + (p + p')^2}{m_i^2 + (p - p')^2};
\]  
(7.25)

\[
\phi(m_i, p, p) = \frac{p^2 + p'^2 + m_i^2}{4pp'} \theta_i(m_i, p, p') - 1.
\]  
(7.26)

<table>
<thead>
<tr>
<th>( i )</th>
<th>( A_i )</th>
<th>( B_i )</th>
<th>( C_i )</th>
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<tbody>
<tr>
<td>( \sigma )</td>
<td>( -g_x^2 \bar{m}^3 (2 + \bar{m}) \theta_\sigma )</td>
<td>( A_\sigma )</td>
<td>( 2g_x^2 \bar{m}^3 (2 + \bar{m}) \phi_\sigma )</td>
</tr>
<tr>
<td>( \omega )</td>
<td>( 2g_x^2 \theta_\omega )</td>
<td>( -2A_\omega )</td>
<td>( -4g_x^2 \phi_\omega )</td>
</tr>
<tr>
<td>( \pi )</td>
<td>( -3(f_x/m_\pi)^2 m_\pi^2 \theta_\pi )</td>
<td>( A_\pi )</td>
<td>( 6(f_x/m_\pi)^2 [(p^2 + p'^2) \phi_\pi - pp' \theta_\pi] )</td>
</tr>
<tr>
<td>( \rho_\omega )</td>
<td>( 6g_\rho^2 \theta_\rho )</td>
<td>( -2A_\rho )</td>
<td>( -12g_\rho^2 \phi_\rho )</td>
</tr>
<tr>
<td>( \rho_T )</td>
<td>( -3(f_x/2M)^2 m_\rho^2 \theta_\rho )</td>
<td>( 3A_\rho )</td>
<td>( 12(f_x/2M)^2 [(p^2 + p'^2 - m_\rho^2/2) \phi_\rho - pp' \theta_\rho] )</td>
</tr>
</tbody>
</table>

The different components of baryon self energy are

\[
\Sigma_s(p) = \left( \frac{g_x}{m_\sigma} \right)^2 \bar{m}^3 \rho_s \\
+ \frac{1}{4\pi^2} \int_0^{\pi} p' dp' \left[ \hat{M}(p') \sum_i B_i(p, p') \right] \\
+ \frac{1}{2} \hat{P}(p') D(p, p'),
\]  
(7.27)
\[ \Sigma_0(p) = \left( \frac{g_\omega}{m_\omega} \right)^2 \rho_B + \frac{g_\rho}{m_\rho} \left( \rho_{B(\omega)} - \rho_{B(\rho)} \right)^2 + \frac{1}{(4\pi)^2} \frac{1}{p} \int_0^{PP} p' dp' \sum_i A_i(p, p'), \tag{7.28} \]

\[ \Sigma_\pi(p) = \frac{1}{(4\pi)^2} \frac{1}{p} \int_0^{PP} p' dp' \left( \hat{P}(p') \sum_i C_i(p, p') \right) + \frac{1}{2} M(p') D_\rho(p, p'). \tag{7.29} \]

The Fermi momentum is related to the baryon density

\[ \rho_B = \frac{2}{3\pi^2} p_F^3. \tag{7.30} \]

The above expressions also contain scalar density, given by

\[ \rho_s = \frac{2}{\pi^2} \int_0^{PP} p^2 M(p) dp. \tag{7.31} \]

All these quantities are evaluated self-consistently for a particular parameter set given a fixed \( p_F \), or equivalently, \( \rho_B \) for nuclear matter.

For asymmetric nuclear matter and neutron matter the expressions are slightly different as the isospin degeneracy must be considered. For example, in neutron matter the baryon density for a particular Fermi momentum is half the value for nuclear matter as only one species of nucleon is present. The expressions for isovector mesons \( \rho \) and \( \pi \) in Table 7.2 are multiplied by a factor \( 1/3 \) to take into account the isospin degeneracy of neutron matter.

### 7.3 Results

Our first task is to fix the values of the different parameters in the Lagrangian. The calculations have been performed with four mesons. The parameters in the theory are the four meson masses and the six coupling constants. The properties of nuclear matter we have to reproduce are the binding energy per nucleon (-16.0 MeV) and its saturation density (0.16 fm\(^{-3}\)). We fix the masses of the physical mesons to their physical values, \( m_\omega = 783 \text{ MeV} \), \( m_\rho = 770 \text{ MeV} \), and \( m_\pi = 138 \text{ MeV} \), while the bare nucleon mass is taken as \( M = 939 \text{ MeV} \). We fix the \( \pi-N \) and \( \rho-N \) coupling constants to their physical values from the meson exchange theories, namely \( f_\pi^2/4\pi = 0.08 \) and \( g_\rho^2/4\pi = 0.55 \). Following [10], we have fixed the value of \( f_\rho/g_\rho \) to 3.7. This is lower than that suggested from scattering theories. However, HF approximation tends to overestimate the short range contributions and a smaller value of \( f_\rho \) is expected to simulate a more realistic interaction. Hence we have now two coupling constants to be fitted, \( g_\rho \) and \( g_\omega \). The mass of the \( \sigma \)-meson, however, is not yet fixed since it corresponds to
representation of a two \( \pi \) exchange contribution. According to the construction of the N-N potential it should lie between 400 and 600 MeV. Calculations are performed for two possible values of \( m_\sigma \), 571 MeV and 440 MeV. The first value is from OBEP while the second one was obtained by fitting the charge radius of \( ^{16}\text{O} \) in [10]. As will be shown, the results in nuclear matter for the different values of \( m_\sigma \) are nearly identical.

For \( m_\sigma = 571 \) MeV, we obtain the values \( g_1^2/4\pi = 5.330 \) and \( g_2^2/4\pi = 7.301 \). We call it set I. For this set we get the bulk modulus \( K = 368.0 \) MeV and the reduced mass \( M^*/M = 0.670 \) at saturation density. For \( m_\sigma = 440 \) MeV, we get the values \( g_1^2/4\pi = 3.157 \) and \( g_2^2/4\pi = 7.505 \). We call these values set II. We obtain \( K = 374.0 \) MeV and \( M^*/M = 0.675 \) at saturation density for this set. For comparison, we have also performed a Hartree calculation for the scalar meson mass \( m_\sigma = 571 \) MeV. The different coupling constants used are \( g_1^2/4\pi = 6.800 \) and \( g_2^2/4\pi = 7.405 \). We have also adjusted the \( \rho \)-meson coupling constant to be \( g_2^\rho/4\pi = 1.35 \) so that the theoretical value of the symmetry energy matches with that obtained experimentally, viz. 33.2 MeV. The reduced mass is \( M^*/M = 0.73 \) and bulk modulus \( K = 301.0 \) MeV in our calculation. The reduced nucleon mass in Hartree-Fock calculation is seen to be much lower than the value obtained in Hartree calculation and is closer to the values in conventional RMF calculations. The bulk modulus is in turn higher.

### 7.3.1 Nuclear Matter

Fig. 7.1 shows the nucleon effective mass calculated at the Fermi surface for nuclear matter against Fermi momentum using both the parameter sets for Hartree-Fock calculation as well as for the Hartree calculation. In all the figures, we have used the following convention. Solid lines indicate results of Hartree-Fock calculation using set I (\( m_\sigma = 571 \) MeV) while dashed lines are for set II (\( m_\sigma = 440 \) MeV). Dashed-dotted lines are used to show the results of Hartree calculation. As already mentioned, the effective mass decreases considerably as a result of the inclusion of the exchange terms.

In Fig. 7.2 we plot the saturation curve for symmetric infinite nuclear matter obtained in the present chapter as a function of baryon density. The Hartree-Fock curve rises more steeply than the Hartree one as expected from the higher bulk modulus in the first case. Although the results of the present calculation agree with other ones involving derivative coupling around and below nuclear saturation density, at higher density the situation is different. Compared to the results obtained using the Zimanyi or Glendenning Lagrangian,[12] the curve rises sharply. At densities three times normal nuclear density, the energy per nucleon is twice in the present chapter compared to [12]. This is manifested in
Figure 7.1: The effective nucleon mass at Fermi surface for symmetric nuclear matter as a function of Fermi momentum. The solid and the dashed curves denote Hartree-Fock results for set I and II, respectively. The dashed-dotted curve indicates the Hartree results.
the higher bulk modulus value in the present calculation. However, our results are closer to conventional relativistic mean field results.[10]

The symmetry energy for nuclear matter is plotted as a function of saturation baryon density in Fig. 7.3. The value at normal nuclear saturation density is low compared to the accepted value, 33.2 MeV.

Fig. 7.4 shows the three components of self-energy in symmetric nuclear matter calculated at Fermi momentum as functions of the saturation baryon density. As is common with relativistic theories, the binding energy comes from the sensitive cancellation between the large scalar and vector potentials. An interesting effect is the saturation of the scalar self energy $\Sigma_s$ at high density compared to vector self energy $\Sigma_v$. This gives rise to the sharp increase in energy per nucleon at high density observed in the present chapter. Our values are also close to Dirac-Brueckner-Hartree-Fock results.[14]

7.3.2 Neutron Matter

Apart from nuclear matter, a test of the present approach can be its description of neutron matter. With this in mind, we have calculated different properties of neutron matter. The effective nucleon mass in neutron matter is shown as a function of saturation Fermi momentum in Fig. 7.5. Fig. 7.6 shows the energy per particle in neutron matter as a function of Fermi momentum. No liquid-gas phase transition, a feature observed in some other relativistic calculations, is obtained.

The neutron matter equation of state (EOS) is plotted in Fig. 7.7 along with the causal limit $p = E$, where $p$ is the pressure and $E$ is the energy density. Using this EOS, the Tolman-Oppenheimer-Volkoff equations for the general relativistic metric[15] is integrated assuming spherical symmetry to get the neutron star mass as a function of central density of the star. The mass of the neutron star is plotted against the central density in the lower panel of Fig. 7.8. On the right hand side, we plot the neutron star mass ($M_\odot$) in terms of solar mass, $M_\odot$. The maximum star mass values are comparable to other calculations involving only neutron matter. For sigma mass 571 MeV the maximum neutron star mass is $2.12M_\odot$ and radius at maximum mass is 9.54 km and for sigma mass 440 MeV these two values are 2.13 $M_\odot$ and 9.56 km respectively. The corresponding values for Hartree results 2.04$M_\odot$ and 8.95 km, respectively. The maximum limit predicted is more than observed maximum value of around 1.5-1.6 $M_\odot$. However, at the high density present inside the neutron star, one may have to deal with other interacting baryons as well as quarks. This is expected to soften the EOS and reduce the maximum limit. The upper panel of the figure shows the radius of the star as a function of the central density.
Figure 7.2: The energy per nucleon in symmetric nuclear matter as a function of baryon density. See caption of Fig. 7.1 for details.
Figure 7.3: The symmetry energy per nucleon in nuclear matter plotted against saturation baryon density. See caption of Fig. 7.1 for details.
Figure 7.4: The different components of self-energy in symmetric nuclear matter plotted against saturation baryon density. See caption of Fig. 7.1 for details.
Figure 7.5: The effective nucleon mass at Fermi surface for neutron matter as a function of Fermi momentum. See caption of Fig. 7.1 for details.
Figure 7.6: The energy per nucleon in neutron matter shown as a function of Fermi momentum. See caption of Fig. 7.1 for details.
7.4 Summary

The effect of the exchange terms in NN interaction for nuclear matter and neutron matter is studied in hybrid derivative coupling model. Pseudovector coupling between the pion and the nucleon and both vector and tensor couplings between the ρ-meson and the nucleon is included. Contact terms are subtracted for ρ - N and π - N interactions. One attractive feature of the inclusion of the exchange terms is the reduction in the value of the effective nucleon mass. The other important feature that comes out of the present calculation is that the ρ-meson coupling constant obtained from meson exchange theories can be used to get a reasonable value of symmetry energy. The maximum limit of the neutron star mass is high compared to observation. This points to the inadequacy of the present Lagrangian at such a high density as it is restricted only to nucleonic degrees of freedom. Other baryonic degrees may be important and are expected to reduce the mass limit. This Lagrangian works well at ordinary nuclear saturation density. The reduction in effective nucleon mass is expected to result in a larger spin-orbit splitting in finite nuclei.
Figure 7.7: The equation of state in neutron matter. Also shown is the causal limit $p = E$. See caption of Fig. 7.1 for details.
Figure 7.8: The radius (upper panel) and mass (lower panel) of neutron star in the different parameter sets as a function of central density. See caption of Fig. 7.1 for details.
Bibliography