Tensor force and the nuclear symmetry energy

Isaac Vidaña1, Artur Polls2 and Constança Providência1

1 Centro de Física Computacional. Department of Physics, University of Coimbra, PT-3004-516 Coimbra, Portugal
2 Departament d’Estructura i Constituents de la Matèria and Institut de Ciències del Cosmos, Universitat de Barcelona, ES-08028 Barcelona, Spain

E-mail: ividana@fis.uc.pt
E-mail: artur@ecm.ub.es
E-mail: cp@teor.fis.uc.pt

Abstract. We analyze the contribution of the different terms of the nucleon-nucleon interaction to the nuclear symmetry energy \( S_0 \) and its slope parameter \( L \) by using the Hellmann-Feynman theorem. The analysis is done using the Argonne V18 potential plus the Urbana IX three-body force within the framework of the microscopic Brueckner–Hartree–Fock approach. Our results show that the potential part of the nuclear Hamiltonian, and in particular the tensor component of the nuclear force, gives the main contribution to both \( S_0 \) and \( L \). The kinetic contribution to \( S_0 \) is very small and negative, in agreement with recent results.

1. Introduction

Isospin asymmetric nuclear matter is present in nuclei, especially in those far from the stability line, and in astrophysical systems, such as neutron stars [1]. Therefore, a well-grounded understanding of the properties of isospin-rich matter is a necessary ingredient for the advancement of both nuclear physics and astrophysics. However, some of these properties are not well established yet. In particular, the density dependence of the symmetry energy \( E_{\text{sym}}(\rho) \) is still an important source of uncertainties. Its value \( S_0 \) at saturation density \( \rho_0 \) is more or less well established (\( \sim 30 \) MeV), and its behavior below \( \rho_0 \) is now much better known [2]. However, for densities above \( \rho_0 \), \( E_{\text{sym}}(\rho) \) is not well constrained yet, and the predictions from different approaches strongly diverge. Why \( E_{\text{sym}}(\rho) \) is so uncertain is still an open question whose answer is related to our limited knowledge of the nuclear force, and in particular of its spin and isospin dependence [3]. Experimental information on the density dependence of the symmetry energy can be obtained from: (i) the analysis of data of giant [4] and pigmy resonances [5], (ii) isobaric analog states [6], (iii) isospin diffusion measurements [7], (iv) isoscaling [8], (v) meson production in heavy ion collisions [9], (vi) measurements of neutron skin thickness in heavy nuclei [10], (vii) the characterization of the crust-core transition in neutron stars [11], (viii) the analysis of power-law correlations, such as the relation between the radius of a neutron star and the equation of state [12], or the novel constraints recently reported by Steiner and Gandolfi [13] on the basis of neutron star mass and radius measurements, driven partially by the strong correlation between the symmetry energy and its derivative obtained in quantum Monte Carlo calculations of neutron matter.
Theoretically $E_{\text{sym}}(\rho)$ has been determined using both microscopic and phenomenological many-body approaches. Microscopic approaches start from realistic nucleon-nucleon (NN) interactions that reproduce the scattering and bound state properties of the free two-nucleon system and include naturally the isospin dependence [14]. The in-medium correlations are then built using many-body techniques that microscopically account for isospin asymmetric effects such as, for instance, the difference in the Pauli blocking factors of neutrons and protons in asymmetric matter. Among this type of approaches the most popular ones are the Brueckner–Bethe–Goldstone (BBG) [15] and the Dirac–Brueckner–Hartree–Fock (DBHF) [16] theories, the variational method [17], the correlated basis function (CBF) formalism [18], the self-consistent Green’s function technique (SCGF) [19] or, recently, the $V_{\text{lowk}}$ approach [20]. Phenomenological approaches, either relativistic or non-relativistic, are based on effective interactions that are frequently built to reproduce the properties of nuclei [21]. Since many of such phenomenological interactions are built to describe systems close to the symmetric case, their predictions at high asymmetries should be taken with care. Skyrme–Hartree–Fock [22] and relativistic mean field [23] calculations are the most popular ones among them. Nevertheless, in spite of the experimental and theoretical [24] efforts carried out to study the properties of isospin-asymmetric nuclear systems, $E_{\text{sym}}(\rho)$ is still uncertain.

In this work we analyze the contribution of the different terms of the NN interaction to $S_0$ and the slope parameter $L = 3\rho_0 \left( \frac{\partial E_{\text{sym}}(\rho)}{\partial \rho} \right)_{\rho_0}$. The analysis is carried out with the help of the Hellmann–Feynman theorem [25] within the framework of the microscopic Brueckner–Hartree–Fock (BHF) approach [15]. We employ the Argonne V18 (Av18) potential [26] supplemented with the Urbana IX three-body force [27] which for the use in the BHF approach is reduced to an effective two-body density-dependent force by averaging over the coordinates of the third nucleon [28]. We find that the tensor term of the nuclear force gives the largest contribution to both $S_0$ and $L$ [29].

The paper is organized in the following way: In Sec. 2 we briefly review the BHF approach of asymmetric nuclear matter. Our results are presented in Sec. 3. Finally, a summary and our main conclusions are given in Sec. 4.

2. BHF approach of asymmetric nuclear matter

Assuming charge symmetry of nuclear forces, the energy per particle of asymmetric nuclear matter can be well approximated in terms of the isospin asymmetry parameter, $\beta = (N-Z)/(N+Z) = (\rho_n - \rho_p)/\rho$, as

$$\frac{E}{A}(\rho, \beta) \sim E_{\text{SNM}}(\rho) + E_{\text{sym}}(\rho)\beta^2 + O(4),$$

(1)

where $E_{\text{SNM}}(\rho)$ is the energy per particle of symmetric nuclear matter, and $E_{\text{sym}}(\rho) = E/A(\rho, \beta = 1) - E_{\text{SNM}}(\rho)$ is a good approximation of the symmetry energy.

It is common to characterize the density dependence of the energy per particle of symmetric matter around the saturation density $\rho_0$ in terms of a few bulk parameters by expanding it in a Taylor series around $\rho_0$,

$$E_{\text{SNM}}(\rho) = E_0 + \frac{K_0}{2} \left( \frac{\rho - \rho_0}{3\rho_0} \right)^2 + \frac{Q_0}{6} \left( \frac{\rho - \rho_0}{3\rho_0} \right)^3 + O(4).$$

(2)

The coefficients denote, respectively, the energy per particle, the incompressibility coefficient, and the third derivative of symmetric matter at saturation,

$$E_0 = E_{\text{SNM}}(\rho_0), \quad K_0 = 9\rho_0^2 \left. \frac{\partial^2 E_{\text{SNM}}(\rho)}{\partial \rho^2} \right|_{\rho=\rho_0}, \quad Q_0 = 27\rho_0^3 \left. \frac{\partial^3 E_{\text{SNM}}(\rho)}{\partial \rho^3} \right|_{\rho=\rho_0}.$$

(3)
Similarly, the behaviour of the symmetry energy around saturation can be also characterized in terms of a few bulk parameters,

\[ E_{\text{sym}}(\rho) = S_0 + L \left( \frac{\rho - \rho_0}{3\rho_0} \right) + \frac{K_{\text{sym}}}{2} \left( \frac{\rho - \rho_0}{3\rho_0} \right)^2 + \frac{Q_{\text{sym}}}{6} \left( \frac{\rho - \rho_0}{3\rho_0} \right)^3 + \mathcal{O}(4), \]

where \( S_0 \) is the value of the symmetry energy at saturation, and the quantities \( L, K_{\text{sym}} \) and \( Q_{\text{sym}} \) are related to its slope, curvature and third derivative, respectively, at such density,

\[ L = 3\rho_0 \left. \frac{\partial E_{\text{sym}}(\rho)}{\partial \rho} \right|_{\rho=\rho_0}, \quad K_{\text{sym}} = 9\rho_0^2 \left. \frac{\partial^2 E_{\text{sym}}(\rho)}{\partial \rho^2} \right|_{\rho=\rho_0}, \quad Q_{\text{sym}} = 27\rho_0^3 \left. \frac{\partial^3 E_{\text{sym}}(\rho)}{\partial \rho^3} \right|_{\rho=\rho_0}. \]

The BHF approach is the lowest order of the BBG many-body theory [15]. It starts with the construction of all the G-matrices describing the effective interaction between two nucleons in the presence of a surrounding medium. They are obtained by solving the well known Bethe–Goldstone equation

\[ G_{\tau_1 \tau_2 \rightarrow \tau_3 \tau_4}(\omega) = V_{\tau_1 \tau_2 \rightarrow \tau_3 \tau_4} + \sum_{ij} V_{\tau_1 \tau_2 \rightarrow \tau_3 \tau_7} \frac{Q_{\tau_i \tau_j}}{\omega - \epsilon_i - \epsilon_j + i\eta} G_{\tau_7 \tau_4 \rightarrow \tau_3 \tau_4}(\omega), \]

where \( \tau = n, p \) indicates the isospin projection of the two nucleons in the initial, intermediate and final states, \( V \) denotes the bare nucleon-nucleon interaction, \( Q_{\tau_i \tau_j} \) the Pauli operator that allows only intermediate states compatible with the Pauli principle, and \( \omega \), the so-called starting energy, which corresponds to the sum of the non-relativistic single-particle energies of the interacting nucleons. The single-particle energy \( \epsilon_{\tau}(k) \) of a nucleon with momentum \( k \) is given by

\[ \epsilon_{\tau}(k) = \frac{\hbar^2 k^2}{2m_{\tau}} + \text{Re}[U_{\tau}(\bar{k})], \]

where the single-particle potential \( U_{\tau}(\bar{k}) \) represents the mean field “felt” by a nucleon due to its interaction with the other nucleons of the medium. In the BHF approximation \( U_{\tau}(\bar{k}) \) is calculated through the on-shell energy G-matrix, and is given by

\[ U_{\tau}(\bar{k}) = \sum_{\tau'} \sum_{|\bar{k}|<k_{F_\tau}} \langle \bar{k} | k_{\tau'} \rangle G_{\tau \tau' \rightarrow \tau \tau'}(\omega = \epsilon_{\tau}(k) + \epsilon_{\tau'}(\bar{k}')) | k\bar{k}' \rangle_A, \]

where the sum runs over all neutron and proton occupied states, and the matrix elements are properly antisymmetrized. Once a self-consistent solution of Eqs. (6) and (8) is obtained, the energy per particle of asymmetric nuclear matter (and consequently the symmetry energy) can be calculated in the BHF approach as

\[ E_A(\rho, \beta) = \frac{1}{A} \sum_{\tau} \sum_{|\bar{k}|<k_{F_\tau}} \left( \frac{\hbar^2 k^2}{2m_{\tau}} + \frac{1}{2} \text{Re}[U_{\tau}(\bar{k})] \right). \]

We note here that the so-called continuous prescription [30] for the in-medium potential has been adopted in our calculation. It has been shown by Song et al., [31] that the contribution to the energy per particle from \textit{three-hole-line} diagrams (which account for the effect of three-body correlations) is minimized with this prescription.
Table 1. Kinetic \( \langle T \rangle \) and potential \( \langle V \rangle \) energy contributions to \( E_{NM}, E_0, S_0 \) and \( L \). Units are given in MeV.

|       | \( E_{NM} \) | \( E_0 \) | \( S_0 \) | \( L \) |
|-------|-------------|----------|----------|--------|
| \( \langle T \rangle \) | 53.321      | 54.294   | -0.973   | 14.896 |
| \( \langle V \rangle \) | -34.251     | -69.524  | 35.273   | 51.604 |
| Total | 19.070      | -15.230  | 34.300   | 66.500 |

3. Results
Since the BHF approach does not provide the correlated many-body wave function \( |\Psi\rangle \), this approach cannot give access to the separate contributions of the kinetic and potential energies in the correlated many-body state. However, it has been shown [32] that the Hellmann-Feynman theorem [25] can be used to calculate the ground-state expectation values of both contributions from the derivative of the total energy with respect to a properly introduced parameter. Writing the nuclear matter Hamiltonian as \( H = T + V \), and defining a \( \lambda \)-dependent Hamiltonian \( H(\lambda) = T + \lambda V \), the expectation value of the potential energy is given as

\[
\langle \Psi \rangle = \frac{\langle \Psi | V | \Psi \rangle}{\langle \Psi | \Psi \rangle} = \left( \frac{dE}{d\lambda} \right)_{\lambda=1} .
\]

Then, the kinetic energy contribution \( \langle T \rangle \) can be obtained simply by subtracting \( \langle V \rangle \) from the total energy \( E \). The kinetic and potential energy contributions to the energy of neutron matter \( E_{NM} \), symmetric matter \( E_0 \), \( S_0 \) and \( L \) at saturation \( (\rho_0 = 0.187 \text{ fm}^{-3} \) in our calculation) are shown in Table 1. Note that the kinetic contribution to \( S_0 \) is very small and negative in contrast with the result for a free Fermi gas (FFG), whose contribution at \( \rho_0 \) is \( \sim 14.4 \) MeV. A similar result has been recently found by Xu and Li [33] (see also Carbone et al., [34]). According to Xu and Li, this is due to the strong isospin dependence of the short-range nucleon-nucleon correlations (SRC) induced by the tensor force. In fact, they have shown that the increase of the kinetic energy of symmetric matter due to SRC is much larger than that of neutron matter, the kinetic part of the symmetry energy becoming then negative. Note also that the kinetic contribution to \( L \) is smaller than the corresponding one of the FFG \( (L_{FFG} \sim 29.2 \text{ MeV}) \). The major contribution to both \( S_0 \) and \( L \) is due to the potential part. Note that, in fact, this contribution is practically equal to the total value of \( S_0 \), and it represents \( \sim 78\% \) of \( L \).

In Tables 2 and 3 we show the partial wave, and the spin \((S)\) and isospin \((T)\) channel decompositions of the potential part of \( E_{NM}, E_0, S_0 \), and \( L \) at \( \rho_0 \). We have considered contributions up to \( J = 8 \). Note that the largest contribution to both \( S_0 \) and \( L \) is that of the spin-triplet \((S=1)\) and isospin-singlet \((T=0)\) channel, and in particular that of the \( ^3S_1 \) partial wave. This is due, as we show next, to the effect of the tensor component of the nuclear force that dominates the potential contribution to the symmetry energy and \( L \), mainly through the \( ^3S_1 - ^3D_1 \) channel. Note that this channel, which gives the major contribution to the symmetry energy, does not contribute to neutron matter. Note also that isospin-triplet \((T=1)\) channels give similar contributions to both \( E_{NM} \) and \( E_0 \) which almost cancel out in \( S_0 \).

Finally, we analyze the role played by the different terms of the nuclear force, particularly the one of the tensor, in the determination of \( S_0 \) and \( L \). To such end we apply the Hellmann–Feynman theorem to the separate components of the Av18 potential and the Urbana IX three-body force.
Table 2. Partial wave decomposition of the potential part of $E_{NM}$, $E_0$, $S_0$ and $L$. Units are given in MeV.

| Partial wave | $E_{NM}$ | $E_0$ | $S_0$ | $L$     |
|--------------|----------|-------|-------|--------|
| $^1S_0$      | −23.070  | −19.660 | −3.410 | −3.459 |
| $^3S_1$      | 0        | −45.810 | 45.810 | 71.855 |
| $^1P_1$      | 0        | 4.904   | −4.904 | −18.601 |
| $^3P_0$      | −5.321   | −4.029  | −1.292 | −1.898 |
| $^3P_1$      | 16.110   | 10.720  | 5.390  | 21.949 |
| $^3P_2$      | −16.000  | −9.334  | −6.666 | −21.168 |
| $^1D_2$      | −5.956   | −3.201  | −2.755 | −11.033 |
| $^3D_1$      | 0        | 0.981   | −0.981 | −3.739 |
| $^3D_2$      | 0        | −3.982  | 3.982  | 16.010 |
| $^3D_3$      | 0        | −0.798  | 0.798  | 4.895  |
| $^1F_3$      | 0        | 0.694   | −0.694 | −3.348 |
| $^3F_2$      | −0.695   | −0.229  | −0.466 | −1.799 |
| $^3F_3$      | 2.000    | 0.821   | 1.179  | 4.883  |
| $^3F_4$      | −0.796   | −0.194  | −0.602 | −3.239 |
| $^1G_4$      | −0.812   | −0.247  | −0.565 | −3.036 |
| $^3G_3$      | 0        | −0.001  | 0.001  | 0.441  |
| $^3G_4$      | 0        | −0.213  | 0.213  | 0.449  |
| $^3G_5$      | 0        | −0.057  | 0.057  | 0.650  |
| $^1H_5$      | 0        | 0.029   | −0.029 | 0.107  |
| $^3H_4$      | 0.033    | 0.040   | −0.007 | 0.232  |
| $^3H_5$      | 0.225    | −0.033  | 0.258  | 0.968  |
| $^3H_6$      | 0.043    | 0.034   | 0.009  | 0.144  |
| $^1I_6$      | −0.082   | 0.023   | −0.105 | −0.591 |
| $^3I_5$      | 0        | −0.029  | 0.029  | 0.342  |
| $^3I_6$      | 0        | 0.067   | −0.067 | −0.819 |
| $^3I_7$      | 0        | −0.021  | 0.021  | 0.239  |
| $^1J_7$      | 0        | −0.027  | 0.027  | 0.385  |
| $^3J_6$      | 0.044    | 0.020   | 0.024  | 0.283  |
| $^3J_7$      | −0.062   | −0.060  | −0.002 | −0.313 |
| $^3J_8$      | 0.036    | 0.014   | 0.022  | 0.242  |
| $^1K_8$      | 0.031    | 0.021   | 0.010  | 0.169  |
| $^3K_7$      | 0        | −0.011  | 0.011  | 0.138  |
| $^3K_8$      | 0        | 0.038   | −0.038 | −0.491 |
| $^3L_8$      | 0.021    | 0.006   | 0.015  | 0.166  |

Table 3. Spin (S) and isospin (T) channel decomposition of the potential part of $E_{NM}$, $E_0$, $S_0$ and $L$. Units are given in MeV.

| $(S, T)$ | $E_{NM}$ | $E_0$ | $S_0$ | $L$     |
|----------|----------|-------|-------|--------|
| (0, 0)   | 0        | 5.600 | −5.600 | −21.457 |
| (0, 1)   | −29.889  | −23.064 | −6.825 | −17.950 |
| (1, 0)   | 0        | −49.836 | 49.836 | 90.561 |
| (1, 1)   | −4.362   | −2.224 | −2.138 | 0.450  |
Table 4. Separate contributions to $E_{NM}$, $E_0$, $S_0$ and $L$ from the various components of the Av18 potential (denoted as $\langle V_i \rangle$) and the reduced Urbana force (denoted as $\langle U_i \rangle$). Units are given in MeV.

|       | $E_{NM}$ | $E_0$   | $S_0$ | $L$     |
|-------|----------|---------|-------|---------|
| $\langle V_1 \rangle$ | -31.212  | -32.710 | 1.498 | -5.580  |
| $\langle V_i \rangle$ | -4.957   | 3.997   | -8.954| -20.383 |
| $\langle V_{ij} \rangle$ | -0.319   | -0.382  | 0.063 | 2.392   |
| $\langle V_{ij}(\bar{\tau}_i, \bar{\tau}_j) \rangle$ | -5.724   | -11.388 | 5.664 | 2.521   |
| $\langle V_{ij} \rangle$ | -0.792   | 1.912   | -2.704| -4.998  |
| $\langle V_{ij}(\bar{S}, \bar{S}) \rangle$ | -4.989   | -37.592 | 32.603| 47.095  |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | -7.538   | -1.754  | -5.784| -12.251 |
| $\langle V_{ij}(\bar{S}, \bar{S}) \rangle$ | -2.671   | -6.539  | 3.868 | 3.969   |
| $\langle V_{ij} \rangle$ | 11.850   | 13.610  | -1.760| 1.521   |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | -2.788   | 0.270   | -3.058| -14.202 |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | 1.265    | 1.383   | -0.118| 1.405   |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | 0.051    | 0.008   | 0.043 | -0.341  |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | 4.194    | 5.682   | -1.488| -0.327  |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | 5.169    | -6.190  | 11.359| 31.368  |
| $\langle V_{ij} \rangle$ | 0.003    | 0.039   | -0.036| -0.022  |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | -0.017   | -0.106  | 0.089 | 0.042   |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | 0.004    | 0.079   | -0.075| -0.124  |
| $\langle V_{ij}(\bar{L}, \bar{S}) \rangle$ | -0.084   | -0.001  | -0.083| -0.331  |

The Av18 potential has 18 components of the form $v_p(r_{ij})O_{ij}^p$ with

$$O_{ij}^{p=1,18} = 1, \bar{\tau}_i \cdot \bar{\tau}_j, \bar{\sigma}_i \cdot \bar{\sigma}_j, (\bar{\sigma}_i \cdot \bar{\sigma}_j)(\bar{\tau}_i \cdot \bar{\tau}_j),$$

$$S_{ij}, S_{ij}(\bar{\tau}_i \cdot \bar{\tau}_j), \bar{L} \cdot \bar{S}, \bar{L} \cdot \bar{S}(\bar{\tau}_i \cdot \bar{\tau}_j), L^2, \bar{L}^2(\bar{\tau}_i \cdot \bar{\tau}_j), L^2(\bar{\sigma}_i \cdot \bar{\sigma}_j)(\bar{\tau}_i \cdot \bar{\tau}_j), (\bar{L} \cdot \bar{S})^2, \bar{L}^2(\bar{S}, \bar{S}) \cdot \bar{S}, \bar{L}^2(\bar{S}, \bar{S}) \cdot \bar{S}, T_{ij}, (\bar{\sigma}_i \cdot \bar{\sigma}_j)T_{ij}, S_{ij}T_{ij}, (\tau_{zi} + \tau_{zj})$$

being $S_{ij}$ the usual tensor operator, $\bar{L}$ the relative orbital angular momentum, $\bar{S}$ the total spin of the nucleon pair, and $T_{ij} = 3\tau_{zi}\tau_{zj} - \tau_i \cdot \tau_j$ the isotensor operator defined analogously to $S_{ij}$. Note that the last four operators break the charge independence of the nuclear interaction.

As we said above, the Urbana IX three-body force is reduced to an effective density-dependent two-body force when used in the BHF approach. For simplicity, in the following we refer to it as reduced Urbana force. This force is made of 3 components of the type $u_p(r_{ij}, \rho)O_{ij}^p$ where $O_{ij}^{p=1,3} = 1, (\bar{\sigma}_i \cdot \bar{\sigma}_j)(\bar{\tau}_i \cdot \bar{\tau}_j), S_{ij}(\bar{\tau}_i \cdot \bar{\tau}_j)$, introducing additional central, $\sigma \tau$ and tensor terms (see e.g., Baldo and Ferreira in Ref. [28] for details).

The separate contributions to $E_{NM}$, $E_0$, $S_0$ and $L$ from the various components of the Av18 potential and the reduced Urbana force are given in Table 4. The contribution from the
tensor component to \( S_0 \) and \( L \) (contributions \( \langle V_{S,i} \rangle \) and \( \langle V_{S,i}(\vec{\tau} \cdot \vec{\tau}) \rangle \) from the Av18 potential, and \( \langle U_{S,i}(\vec{\tau} \cdot \vec{\tau}) \rangle \) from the reduced Urbana force) is 36.056 MeV and 69.968 MeV, respectively. These results clearly confirm that the tensor force gives the largest contribution to both \( S_0 \) and \( L \). The contributions from the other components are either negligible, as for instance the contribution from the charge symmetry breaking terms \( \langle \langle V_{T,i} \rangle \rangle, \langle V(\vec{\tau} \cdot \vec{\sigma})_i T_{ij} \rangle, \langle V_{S,i} T_{ij} \rangle \) and \( \langle V(\tau_{S,i} + \tau_{S,j}) \rangle \), or almost cancel out.

4. Summary and conclusions

In summary, using the Hellmann–Feynman theorem we have evaluated the separate contribution of the different terms of the nuclear force to the nuclear symmetry energy \( S_0 \) and the slope parameter \( L \). Our study has been done within the framework of the BHF approach using the Av18 potential plus an effective density-dependent two-body force deduced from the Urbana IX three-body one. Our results show that the potential part of the nuclear Hamiltonian gives the main contribution to both \( S_0 \) and \( L \). The kinetic contribution to \( S_0 \) is very small and negative in agreement with the recent results of Xu and Li [33]. We have performed a partial wave, and a spin-isospin channel decomposition of the potential part of \( S_0 \) and \( L \), showing that the major contribution to them is given by the spin-triplet \( (S = 1) \) and isospin-singlet \( (T = 0) \) channel. This is due, as we have explicitly shown, to the dominant effect of the tensor force which gives the largest contribution to both \( S_0 \) and \( L \). In conclusion, our results confirm the critical role of the tensor force in the determination of the symmetry energy and its density dependence.

Acknowledgments

This work is partly supported by COMPSTAR, an ESF (European Science Foundation) Research Networking Programme; by the initiative QREN financed by the UE/FEDER through the Programme COMPETE under the projects, PTDC/FIS/113292/2009, CERN/FP/109316/2009, CERN/FP/116366/2010 and CERN/FP/123608/2011; by the Consolider Ingenio 2010 Programme CPAN CSD2007-00042, Grant No. FIS2008-01661 from MEC and FEDER (Spain) and Grant No. 2009GR-1289 from Generalitat de Catalunya (Spain).

References

[1] Baran V, Colonna, Greco V and Di Toro M 2005 Phys. Rep. 410 355
[2] Tsang M B et al. 2011 Prog. Part. Nuc. Phys. 66 400
[3] Pandharipande V R and Garde V K 1972 Phys. Lett. B 39 608
[4] Klimkiewicz A et al. 2007 Phys. Rev. C 76 051603 (R)
[5] Klimkiewicz A et al. 2007 Phys. Rev. C 76 051603 (R)
[6] Danielewicz P and Lee J 2009 Nucl. Phys. A 818 36
[7] Chen L W, Ko M and Li B A 2005 Phys. Rev. Lett. 94 032701
[8] Shetty D V, Yennello S J and Souliotis G A 2009 Phys. Rev. C 76 024606
[9] Li B A, Yong G C and Zuo W 2005 Phys. Rev. C 71 014608
[10] Brown B A 2000 Phys. Rev. Lett. 85 5296

Typel S and Brown B A 2001 Phys. Rev. C 64 027302

doi:10.1088/1742-6596/420/1/012091
