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Nuclear Matter Bulk Parameter Correlations from a Nonrelativistic Solvable Approach and Beyond.

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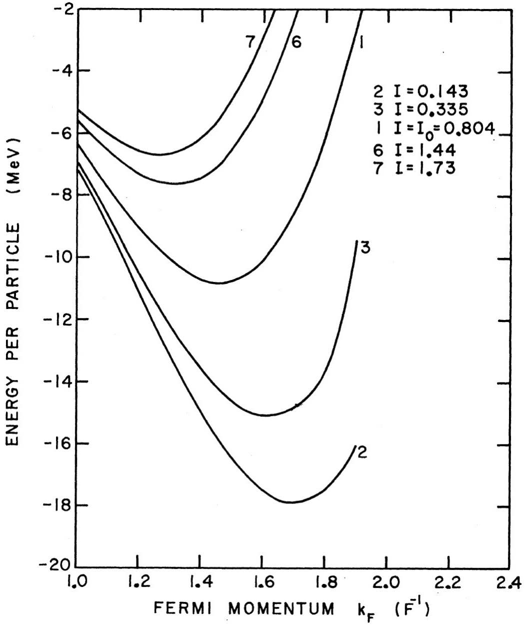
Summary

1. Nuclear matter energy per particle and saturation density correlation
2. Few-Body Scales
3. Non-relativistic limit for NLPC models
4. Correlations
5. Conclusions

Nuclear matter energy per particle and saturation density correlation

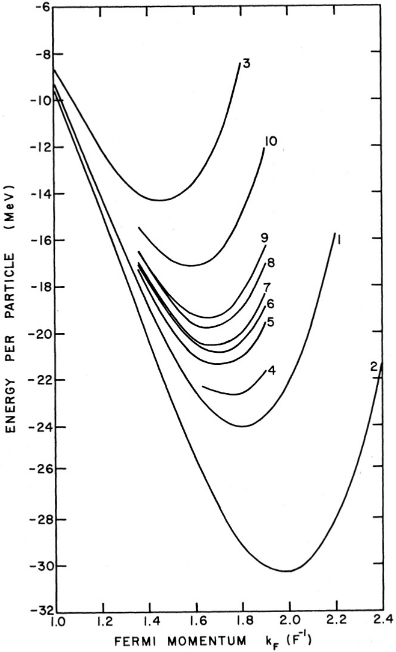


Nuclear force: SHCP

Remark: The binding energy per particle is calculated in the Bruckner approximation with self consistent single particle energies below the Fermi level.

Ref. [F. Coester, S. Cohen, B. Day, and C.M. Vincent, Phys. Rev. C **1**, 3 (l970)].

Where $\rho = \frac{2}{3\pi^2} k_F^3$.



Nuclear force: Yukava-Core Potential

Remark: The binding energy per particle is calculated in the Bruckner approximation with self consistent single particle energies below the Fermi level.

Ref. [F. Coester, S. Cohen, B. Day, and C.M. Vincent, Phys. Rev. C **1**, 3 (l970)].

Where $\rho = \frac{2}{3\pi^2} k_F^3$.

Few-Body Scales

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The few scales of nuclei and nuclear matter

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(AV18 + 3BF):

Ref. [R.B. Wiringa and S.C. Pieper, PRL **89**, 182501 (2002)]

Remark:

There is a systematic improvement of the Binding Energy results for He, Li, Be, and B isotopes simultaneously with the B_t when models are tuned to fit B_t.

Few-Body Scales

In the limit of a zero-range interaction, we write the binding energy of a nucleus with mass number A and isospin projection I_z, considering isospin breaking effects, as

$$B_{(A,I_z)} = AB_t\mathcal{B}(\beta_v, \beta_d, \beta_\alpha, A, I_z), \quad (1)$$

where $\beta_\alpha = B_a/B_t$ whit $a = v$, d and α.

According to the Tjon line, β_α remains approximately constant for a variety of two-nucleon potentials and the parametrization of the numerical results, given in MeV, for several two-nucleon potentials is

$$B_\alpha = 4.72(B_t - 2.48) \quad (2)$$

which for $B_t^{exp} = 8.48$ MeV gives $B_\alpha^{exp} = 28.32$ MeV.

Using (2) in (1), we obtain

$$R(A, I_z) = B_{(A,I_z)}/A = B_t \mathcal{R}(B_t, A, I_z), \tag{3}$$

where in the scaling function $\mathcal{R}(A, I_z)$ the values of B_d and B_v are fixed to the experimental values.

We suppose that going to the infinite isospin symmetrical nuclear matter, $A \to \infty$ and $I_z = 0$, the limit

$$\frac{B_A}{A} \doteq \frac{B_t}{A} \lim_{A\to\infty} \mathcal{B}(\beta_v, \beta_d, \beta_\alpha, A, I_z = 0) = B_t\, \mathcal{G}(\beta_v, \beta_d, \beta_\alpha), \tag{4}$$

is well defined and expresses the correlation between the binding energy of the nucleon in nuclear matter with the few-nucleon scales. The Fermi energy

$$E_F = B_t\, \mathcal{E}_F(\beta_v, \beta_d, \beta_\alpha), \tag{5}$$

will be correlated as well to the few-nucleon binding energies.

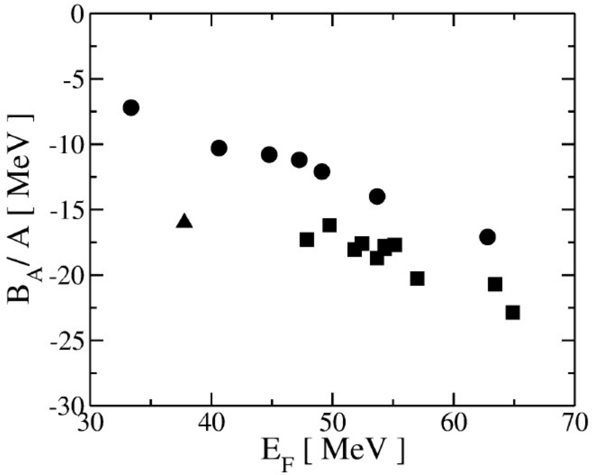


Figure: Infinite nuclear matter binding energy as a function of E_F extracted from Ref. [R. Machleidt, Adv. Nucl. Phys. **19** (1989) 189] (solid circles and squares). The squares includes the single particle contribution in the continuum. The full triangle is given by the empirical values.

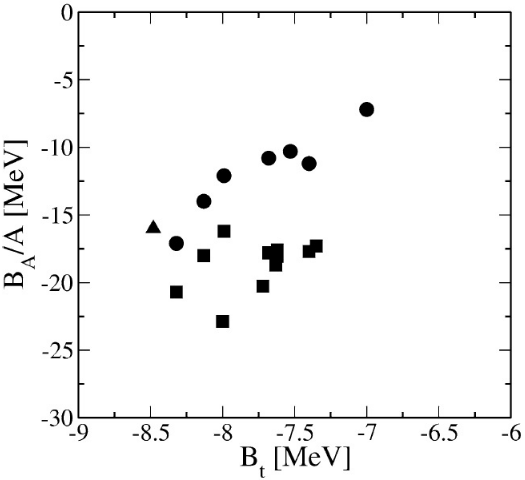


Figure: B_A/A as a function of Bt extracted from Ref. [R. Machleidt, Adv. Nucl. Phys. **19** (1989) 189] (solid circles and squares). The squares includes the single particle contribution in the continuum. The full triangle is given by the empirical values.

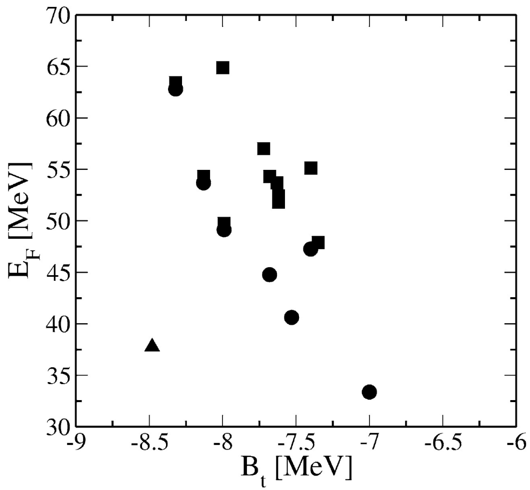


Figure: E_F as a function of Bt extracted from Ref. [R. Machleidt, Adv. Nucl. Phys. **19** (1989) 189] (solid circles and squares). The squares includes the single particle contribution in the continuum. The full triangle is given by the empirical values.

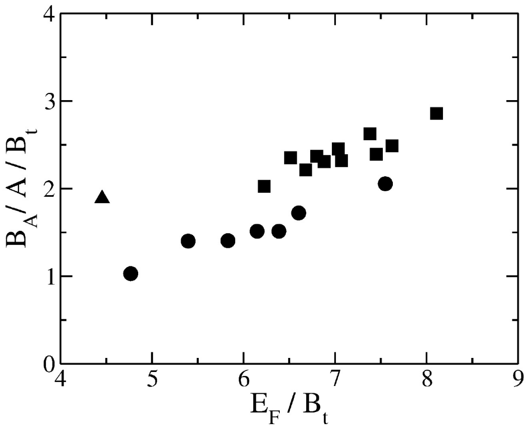


Figure: Infinite nuclear matter binding energy as a function of EF , both in units of the triton binding energy. The calculation results are extracted from Ref. [R. Machleidt, Adv. Nucl. Phys. **19** (1989) 189] (solid circles and squares). The squares includes the single particle contribution in the continuum. The full triangle represents the empirical values.

Non-relativistic limit for NLPC models

Non-relativistic limit for NLPC models

$$\mathcal{L}_{\text{NLPC}} = \bar{\psi}(i\gamma^{\mu}\partial_{\mu} - M)\psi - \frac{1}{2}G_{\text{V}}^{2}(\bar{\psi}\gamma^{\mu}\psi)^{2} + \frac{1}{2}G_{\text{S}}^{2}(\bar{\psi}\psi)^{2} + \frac{A}{3}(\bar{\psi}\psi)^{3} + \frac{B}{4}(\bar{\psi}\psi)^{4} - \frac{1}{2}G_{\text{TV}}^{2}(\bar{\psi}\gamma^{\mu}\vec{\tau}\psi)^{2}, \quad (6)$$

We perform the nonrelativistic limit of the NLPC models by rewriting the small component (χ) of the fermion field ψ in terms of the large one (ϕ) in the Dirac equation

$$(\boldsymbol{\sigma} \cdot \boldsymbol{k}\, B\, \boldsymbol{\sigma} \cdot \boldsymbol{k} \; + \; M + S + V)\phi \quad = \quad E\phi \qquad (7)$$

where B have been expanded from the parameter $x = (\epsilon - S - V)B_0 = (E - M - S - V)B_0$. Thus,

$$\begin{aligned} B &= B_0 \frac{1}{1 + (\epsilon - S - V)B_0} \simeq B_0 + B_0^2(S + V - \epsilon), \\ B_0 &= \frac{1}{2(M + S)} \end{aligned} \qquad (8)$$

Reduces to the Schrödinger equation:

$$\hat{H}^{class}\varphi^{class} = \epsilon\varphi^{class}, \qquad \text{where} \qquad \varphi^{class} = \hat{I}^{1/2}\phi, \tag{9}$$

$$\hat{H}^{class} = \hat{I}^{-1/2}\left[\boldsymbol{\sigma}\cdot\boldsymbol{k}\,B_0\,\boldsymbol{\sigma}\cdot\boldsymbol{k} + S + V + \boldsymbol{\sigma}\cdot\boldsymbol{k}\,B_0^2(S+V)\,\boldsymbol{\sigma}\cdot\boldsymbol{k}\right]\hat{I}^{-1/2}, \tag{10}$$

$$\hat{I} = 1 + \boldsymbol{\sigma}\cdot\boldsymbol{k}\,B_0^2\,\boldsymbol{\sigma}\cdot\boldsymbol{k} = 1 + \frac{z^2}{4} = 1 + x(x+1). \tag{11}$$

These expansions lead to the vector and scalar densities

$$\begin{aligned} \rho &= \phi^\dagger\phi + \chi^\dagger\chi = |\varphi^{class}|^2, \\ \rho_s &= \phi^\dagger\phi - \chi^\dagger\chi = \rho(1 - z^2/2), \end{aligned} \tag{12}$$

and to the single-particle energy $\hat{H}^{class}$, that now reads

$$H^{class} = \frac{k^2}{2(M+S)} + S + V \tag{13}$$

$$H^{class} = \frac{k^2}{2(M+S)} + (G_V^2 - G_s^2)\rho - A\rho^2 - B\rho^3 + 2B_0^2k^2\rho(G_s^2 + 2A\rho + 3B\rho^2). \tag{14}$$

With this procedure and using the continuous limit in H^{class}, we have

$$\mathcal{E}_{NR} = c_1\rho^2 + c_2\rho^3 + c_3\rho^4 + c_4(\rho)\frac{3}{40}\left(\frac{3\pi^2}{2}\right)^{2/3}\rho^{8/3} + \frac{3}{10M}\left(\frac{3\pi^2}{2}\right)^{2/3}\rho^{5/3}, \quad \text{where} \tag{15}$$

$$c_1 = G_V^2 - G_s^2, \quad c_2 = -A, \quad c_3 = -B, \quad \text{and} \quad c_4(\rho) = \frac{4}{M^2}(G_s^2 + 2A\rho + 3B\rho^2). \tag{16}$$

The nucleon effective mass will now be defined by its standard nonrelativistic as follows

$$M^* = k\left[\frac{\partial H^{class}}{\partial k}\right]^{-1} = M\left[1 + \frac{Mc_4(\rho)\rho}{4}\right]^{-1}, \tag{17}$$

where again we have used $M + S = M$ in Eq. (14).

Correlations

- The energy per particle ($E = \varepsilon/\rho$) expanded as a function of the nuclear density ρ

$$E(x) = E_\infty + \frac{1}{2}K_\infty x^2 + \frac{1}{6}Q_\infty x^3 + \mathcal{O}(x^4), \quad (18)$$

$x = (\rho - \rho_o)/(3\rho_o) \rightarrow$ expansion parameter; $E_\infty \rightarrow$ binding energy at the saturation density ρ_o; $K_\infty \rightarrow$ incompressibility at $\rho = \rho_o$; $Q_\infty \rightarrow$ third derivative (skewness) of the energy per particle at $\rho = \rho_o$.

- The symmetry energy ($\mathcal{S}$) expanded as a function of the nuclear density ρ

$$\mathcal{S}(\rho) = J + Lx + \frac{1}{2}K_{\text{sym}} x^2 + \frac{1}{6}Q_{\text{sym}} x^3 + \mathcal{O}(x^4) \quad (19)$$

$x = (\rho - \rho_o)/(3\rho_o) \rightarrow$ expansion parameter; $J \rightarrow$ symmetry energy at the saturation density ρ_o; $L \rightarrow$ symmetry energy slope at $\rho = \rho_o$; $K_{\text{sym}} \rightarrow$ symmetry energy curvature at $\rho = \rho_o$; $Q_{\text{sym}} \rightarrow$ third derivative (skewness) of symmetry energy $\mathcal{S}$ at $\rho = \rho_o$.

These quantities are defined by

$$K_\infty = 9\frac{\partial P}{\partial \rho}\Big|_{\rho=\rho_o, y=1/2}; \qquad Q_\infty = 27\rho_o^3 \frac{\partial^3 (\varepsilon/\rho)}{\partial \rho^3}\Big|_{\rho=\rho_o, y=1/2}; \tag{20}$$

$$\mathcal{S}(\rho) = \frac{1}{8}\left[\frac{\partial^2(\varepsilon^{(\mathrm{NR})}/\rho)}{\partial y^2}\right]_{y=\frac{1}{2}}; \qquad J = \mathcal{S}(\rho_o); \qquad L = 3\rho_o\left(\frac{\partial \mathcal{S}}{\partial \rho}\right)_{\rho=\rho_o}; \tag{21}$$

$$K_{sym} = 9\rho_o^2\left(\frac{\partial^2 \mathcal{S}}{\partial \rho^2}\right)_{\rho=\rho_o}; \qquad Q_{sym} = 27\rho_o^3\left(\frac{\partial^3 \mathcal{S}}{\partial \rho^3}\right)_{\rho=\rho_o}. \tag{22}$$

The energy density functional at zero temperature for asymmetric nuclear matter is written as

$$\varepsilon(\rho, y) = (G_V^2 - G_S^2)\rho^2 - A\rho^3 - B\rho^4 + G_{TV}^2\rho^2(2y-1)^2 + \frac{3}{10M^*(\rho, y)}\lambda\rho^{\frac{5}{3}}, \tag{23}$$

where the effective mass is

$$M^*(\rho, y) = \frac{M^2}{(M + G_S^2\rho + 2A\rho^2 + 3B\rho^3)H_{\frac{5}{3}}}, \tag{24}$$

with $H_{\frac{5}{3}} = 2^{\frac{2}{3}}[y^{\frac{5}{3}} + (1-y)^{\frac{5}{3}}]$, $\lambda = (3\pi^2/2)^{\frac{2}{3}}$, and $y = \rho_p/\rho$ being the proton fraction of the system. The proton density is ρ_p.

The pressure is defined by $P(\rho, y) = \rho^2 \frac{\partial(\varepsilon/\rho)}{\partial\rho}$,

$$P(\rho, y) = (G_V^2 - G_S^2)\rho^2 - 2A\rho^3 - 3B\rho^4 + G_{TV}^2\rho^2(2y-1)^2 + \frac{\lambda H_{5/3}}{5M^2}\left(M + \frac{5}{2}G_S^2\rho + 8A\rho^2 + \frac{33}{2}B\rho^3\right)\rho^{\frac{5}{3}} \quad (25)$$

The chemical potential is defined by $\mu = \partial\varepsilon/\partial\rho$,

$$\mu(\rho, y) = 2(G_V^2 - G_S^2)\rho - 3A\rho^2 - 4B\rho^3 + 2G_{TV}^2\rho(2y-1)^2 + \frac{\lambda H_{5/3}}{5M^2}\left(\frac{5}{2}M + 4G_S^2\rho + 11A\rho^2 + 21B\rho^3\right)\rho^{\frac{2}{3}} \quad (26)$$

Thermodynamic consistency: $\mu(\rho, y) = [\varepsilon(\rho, y) + P(\rho, y)]/\rho$

The incompressibility is defined by $K(\rho, y) = 9\frac{\partial P}{\partial \rho}$, is given by

$$\begin{aligned} K(\rho, y) &= 18(G_{\mathrm{V}}^2 - G_{\mathrm{S}}^2)\rho - 54A\rho^2 \\ &-108B\rho^3 + 18G_{\mathrm{TV}}^2\rho(2y-1)^2 \\ &+\frac{3\lambda H_{\frac{5}{3}}}{M^2}\left(M + 4G_{\mathrm{S}}^2\rho + \frac{88}{5}A\rho^2 + \frac{231}{5}B\rho^3\right)\rho^{\frac{2}{3}}, \end{aligned} \tag{27}$$

with $H_{\frac{5}{3}} = 2^{\frac{2}{3}}[y^{\frac{5}{3}} + (1-y)^{\frac{5}{3}}]$, $\lambda = (3\pi^2/2)^{\frac{2}{3}}$, $y = \rho_p/\rho$

We rewrite the coupling constants of the model, namely, G_{S}^2, G_{V}^2, A, and B, in terms of the bulk parameters m^*, ρ_o, B_o, and K_o. This is done by solving a system of four equations, namely, $\varepsilon(\rho_o, 1/2) = -B_o$, $K(\rho_o, 1/2) = K_o$, $P(\rho_o, 1/2) = 0$ (nuclear saturation), and $M^*(\rho_o, 1/2) = M_o^*$.

The symmetry energy is defined by $\mathcal{S}(\rho) = \frac{1}{8}\left[\frac{\partial^2(\varepsilon^{(\mathrm{NR})}/\rho)}{\partial y^2}\right]_{y=\frac{1}{2}}$ and $J = \mathcal{S}(\rho_o)$ is given by

$$J = \frac{\lambda\rho_o^{\frac{2}{3}}}{6M} + \left(G_{\mathrm{S}}^2 + 2A\rho_o + 3B\rho_o^2\right)\frac{\lambda\rho_o^{\frac{5}{3}}}{6M^2} + G_{\mathrm{TV}}^2\rho_o. \tag{28}$$

The symmetry energy $\mathcal{S}(\rho)$ is used again in order to obtain $L = 3\rho_o\left[\frac{\partial\mathcal{S}(\rho)}{\partial\rho}\right]_{\rho=\rho_o}$ and the result is

$$L = \frac{\lambda\rho_o^{\frac{2}{3}}}{3M} + \left(5G_{\mathrm{S}}^2 + 16A\rho_o + 33B\rho_o^2\right)\frac{\lambda\rho_o^{\frac{5}{3}}}{6M^2} + 3G_{\mathrm{TV}}^2\rho_o. \tag{29}$$

We write $L = L(m^*, \rho_o, B_o, K_o)$ and subtracting $3J$ from L, we obtain

$$L = 3J + f(m^*, \rho_o, B_o, K_o), \quad \text{where} \tag{30}$$

$$\begin{aligned} f(m^*, \rho_o, B_o, K_o) &= \frac{5E_F^o}{\left(3M^2 - 19E_F^o M + 18E_F^{o2}\right)} \\ &\times \left\{ \frac{2M}{9m^*}(3M - 14E_F^o) - M\,(M + K_o/9) + E_F^o(5M + 6B_o) \right\} \end{aligned} \tag{31}$$

exhibits a dependence with the inverse of the effective mass, with $E_F^o = 3\lambda\rho_o^{\frac{2}{3}}/10M$.

- Usually, in nuclear mean-field models, the binding energy and the saturation density are well established close around the values of $B_o = 16$ MeV and $\rho_o = 0.15$ fm^{-3}.

- The same assumption does not apply to the incompressibility and effective mass.

- Thus, we analyze how the function varies with the incompressibility for a fixed value of the effective mass.

- And we analyze how the function varies with the effective mass for a fixed value of the incompressibility.

- For a fixed value of m^*, the variation in f will be given by

$$(\Delta f)_{K_o} = -\frac{5ME_F^o}{9(3M^2 - 19E_{\mathrm{F}}^o M + 18E_{\mathrm{F}}^{o2})}\Delta K_o. \tag{32}$$

For the range of $250 \leq K_o \leq 315$ MeV, we verify that $|(\Delta f)_{K_o}| = 0.32$ MeV.

- For two different models with same incompressibility K_o but with two different effective masses m_1^* and m_2^*, the f variation (with $\Delta m^* = m_2^* - m_1^*$) can be inferred by

$$(\Delta f)_{m^*} = \frac{5ME_F^o(3M - 14E_{\mathrm{F}}^o)}{9(3M^2 - 19E_{\mathrm{F}}^o M + 18E_{\mathrm{F}}^{o2})}\frac{\Delta m^*}{m_1^* m_2^*}, \tag{33}$$

For the range of $0.50 \leq m^* \leq 0.80$, we verify that $|(\Delta f)_{m^*}| = 18$ MeV.

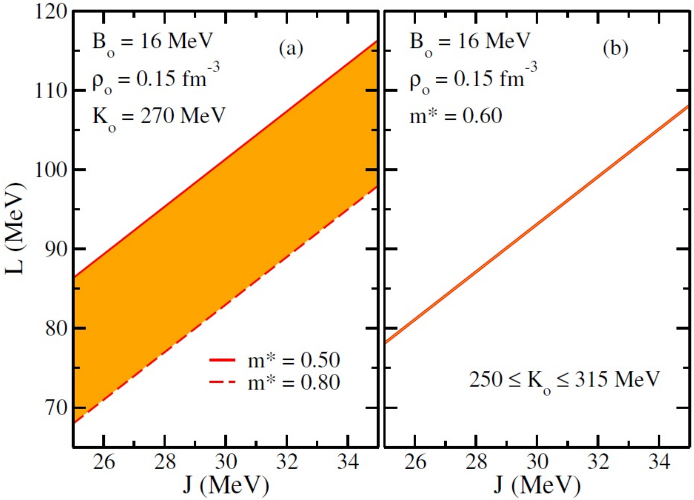


Figure: Effect of Δf in the $L - J$ correlation of Eq. (30) for (a) $0.50 \leq m^* \leq 0.80$, and (b) $250 \leq K_o \leq 315$ MeV. Ref. [B. M. Santos, M. Dutra, O. Lourenço, and A. Delfino, Phys. Rev. C **90**, 035203 (2014)].

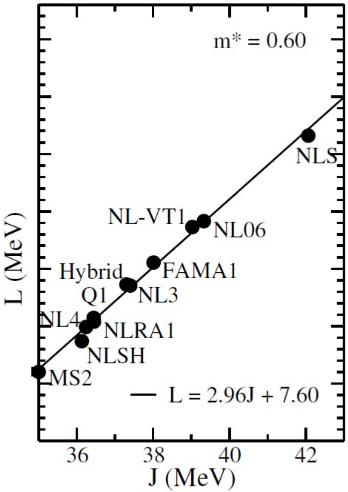
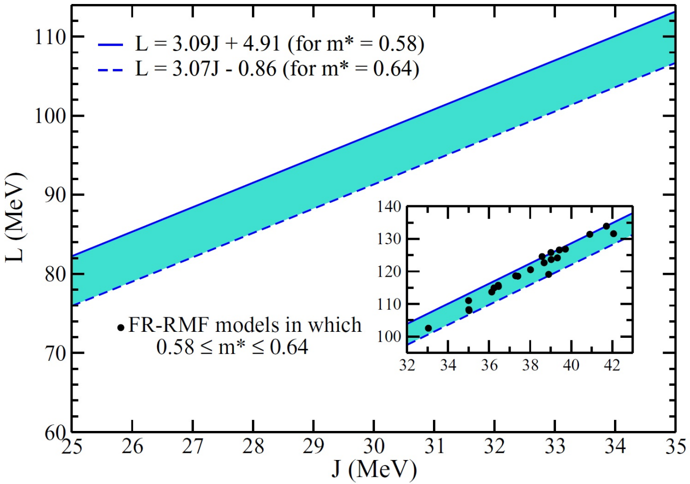


Figure: L versus J for FR parametrizations in which m^* is the same. Ref. [B. M. Santos, M. Dutra, O. Lourenço, and A. Delfino, Phys. Rev. C **90**, 035203 (2014)].



FRS Constraint:
Spin-Orbit Splitting

Ref. [R.J. Furnstahl, J.J. Rusnak and B.D. Serot, Nucl. Phys. **A632**, 607 (1998)].

Figure: Graphic constraint in the L versus J plane, Ref. [B. M. Santos, M. Dutra, O. Lourenço, and A. Delfino, Phys. Rev. C **90**, 035203 (2014)].

The symmetry energy curvature, defined as $K_{\text{sym}} = 9\rho_o^2 \left(\frac{\partial^2 S}{\partial \rho^2}\right)_{\rho=\rho_o}$, is given by

$$K_{\text{sym}} = \left(\frac{1}{m^*} - 1\right) s(\rho_o) + r(\rho_o, B_o, K_o), \quad \text{where} \tag{34}$$

$$s(\rho_o) = \frac{5\lambda\rho_o^{\frac{2}{3}}}{3M}\left[1 + \frac{4E_{\text{F}}^o}{(M - 2E_{\text{F}}^o)} - \frac{E_{\text{F}}^o\,(M - 10E_{\text{F}}^o)\,(19M - 18E_{\text{F}}^o)}{5\,(M - 2E_{\text{F}}^o)\left(3M^2 - 19E_{\text{F}}^o M + 18E_{\text{F}}^{o2}\right)}\right], \tag{35}$$

$$r(\rho_o, B_o, K_o) = -\frac{\lambda\rho_o^{\frac{2}{3}}}{3M}\left[1 + \frac{K_o\,(19M - 18E_{\text{F}}^o)}{2\left(3M^2 - 19E_{\text{F}}^o M + 18E_{\text{F}}^{o2}\right)} - \frac{(81B_oM + 8E_{\text{F}}^o M + 18E_{\text{F}}^{o2})}{3M^2 - 19E_{\text{F}}^o M + 18E_{\text{F}}^{o2}}\right]. \tag{36}$$

By rearranging these equations, we find a simplified form for K_{sym}, namely,

$$K_{\text{sym}} = [L - 3J]\, p(\rho_o) + q(\rho_o, B_o, K_o), \quad \text{where} \tag{37}$$

$$p(\rho_o) = \frac{s(\rho_o)}{g(\rho_o)}, \tag{38}$$

$$\begin{aligned} q(\rho_o, B_o, K_o) &= -h(\rho_o, B_o, K_o) p(\rho_o) + r(\rho_o, B_o, K_o) \\ &= \frac{\lambda \rho_o^{\frac{2}{3}}}{3M} \Bigg\{ \frac{[p(\rho_o) - 2]}{2} + \frac{M E_{\mathrm{F}}^{o} [p(\rho_o) + 8]}{\left(3M^2 - 19 E_{\mathrm{F}}^{o} M + 18 E_{\mathrm{F}}^{o2}\right)} \\ &- \frac{9 E_{\mathrm{F}}^{o2} [p(\rho_o) - 2] + 27 B_o [E_{\mathrm{F}}^{o} p(\rho_o) - 3M]}{\left(3M^2 - 19 E_{\mathrm{F}}^{o} M + 18 E_{\mathrm{F}}^{o2}\right)} + \frac{M\,[p(\rho_o) - 19] + 18 E_{\mathrm{F}}^{o}}{2\left(3M^2 - 19 E_{\mathrm{F}}^{o} M + 18 E_{\mathrm{F}}^{o2}\right)} K_o \Bigg\}. \end{aligned} \tag{39}$$

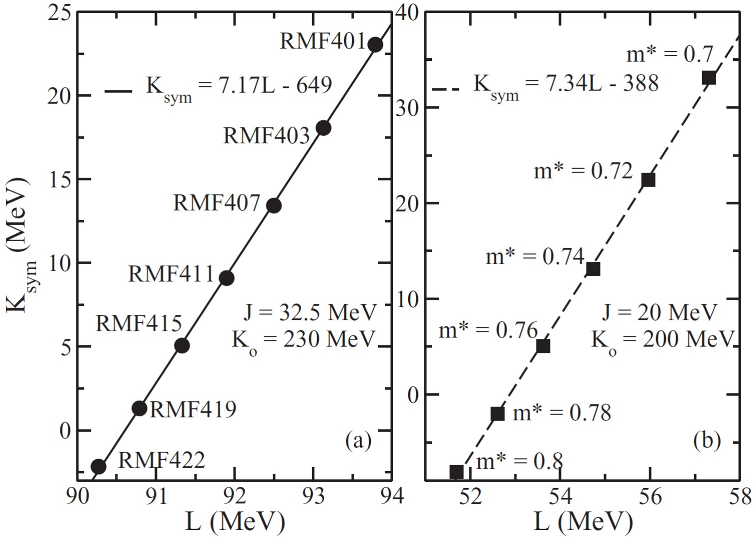


Figure: Correlation between K_{sym} and L plane, Ref. [B. M. Santos, M. Dutra, O. Lourenço, and A. Delfino, Phys. Rev. C **90**, 035203 (2014)].

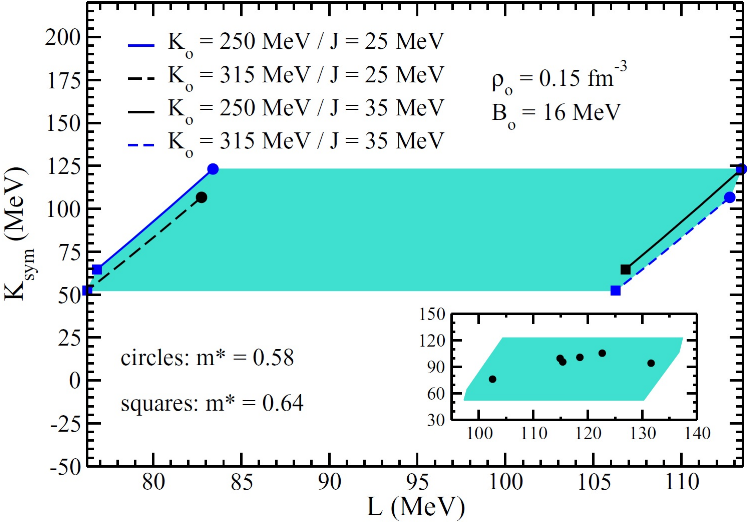


Figure: Graphic constraint in the K_{sym} versus L plane, Ref. [B. M. Santos, M. Dutra, O. Lourenço, and A. Delfino, Phys. Rev. C **90**, 035203 (2014)].

Correlations

Is there correlation between Q_∞ and K_∞? Yes.

PHYSICAL REVIEW C **88**, 034319 (2013)

Determination of the density dependence of the nuclear incompressibility

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The incompressibility at the saturation density ρ_o, defined as $K_\infty = 9\frac{\partial P}{\partial \rho}\Big|_{\rho_o}$, is given by

$$K_\infty = 18(G_{\mathrm{V}}^2 - G_{\mathrm{S}}^2)\rho_o - 54A\rho_o^2 - 108B\rho_o^3 + \frac{3\lambda}{M^2}\left(M + 4G_{\mathrm{S}}^2\rho_o + \frac{88}{5}A\rho_o^2 + \frac{231}{5}B\rho_o^3\right)\rho_o^{\frac{2}{3}}. \quad (40)$$

The skewness parameter at the saturation density ρ_o, defined as $Q_\infty = 27\rho_o^3\frac{\partial^3(\varepsilon/\rho)}{\partial \rho^3}\Big|_{\rho_o}$, is given by

$$Q_\infty = -162(G_{\mathrm{V}}^2 - G_{\mathrm{S}}^2)\rho_o + 324A\rho_o^2 + 324B\rho_o^3 - \frac{3\lambda}{M^2}\left(10M + 28G_{\mathrm{S}}^2\rho_o + \frac{352}{5}A\rho_o^2 + \frac{231}{5}B\rho_o^3\right)\rho_o^{\frac{2}{3}}. \quad (41)$$

We rewrite the coupling constants of the model, namely, G_{S}^2, G_{V}^2, A, and B, in terms of the bulk parameters m^*, ρ_o, B_o, and K_o. After we write $Q_\infty = Q_\infty(m^*, \rho_o, B_o, K_o)$. By doing so, and subtracting $9K_\infty$ from Q_∞, we obtain

$$Q_\infty = 9K_\infty - \frac{202}{5}(L - 3J_o) + K_{\mathrm{sym}} + \frac{49}{75}Q_{sym} - \frac{2}{3}\left(243E_o - \frac{688E_F^o}{9m^*} + \frac{774E_F^o}{5}\right) \tag{42}$$

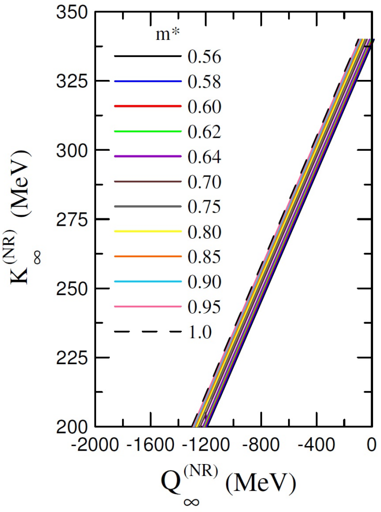

Figure: Incompressibility (K_∞) versus the skewness parameter (Q_∞) at the saturation density. Fixed the parameters $\rho_0 = 0.15\,\mathrm{fm}^{-3}$ and $B_0 = 16$ MeV. The ranger $0.56 < m^* < 1.00$ have been used for ratio $m^* = M_o^*/M$.

Conclusions

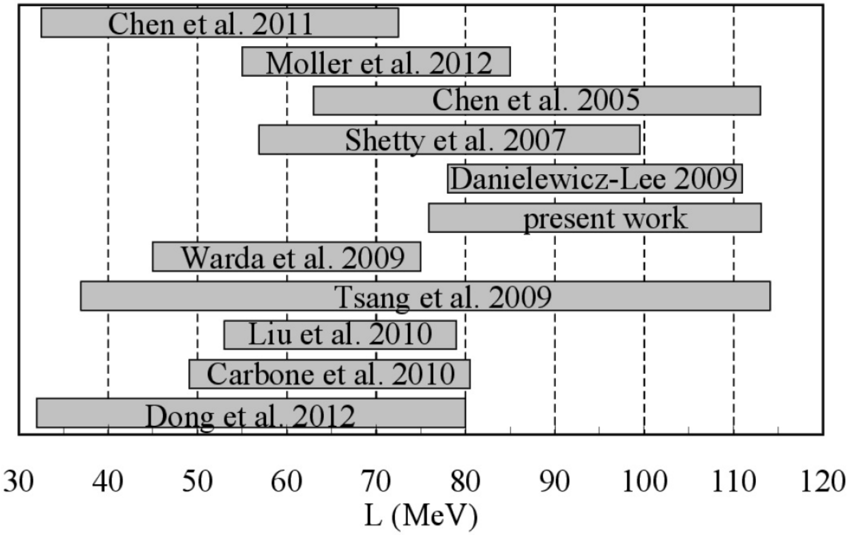

Figure: Comparison between the limits of L obtained in this work and others authors; See reference [J. Dong, W. Zuo, J. Gu, and U. Lombardo, Phys. Rev. C **85**, 034308 (2012).]

Conclusions

- A more systematic study regarding $(B_n, \rho_o) \times B_t$ is underway:

$$V(r) = V_A \frac{e^{-\mu_A r}}{r} + V_R \frac{e^{-\mu_R r}}{r} \qquad (MT5) \tag{43}$$

  Vary V_A, V_R, μ_A and μ_R. So that $B_2 = 2.2$ MeV.

  Calculate B_t, B_N and ρ_o. Look for the correlations.

- Use the NR procedure to find other bulk nuclear matter parameters, looking for new correlations

- Extend the NR procedure to study asymmetric nuclear matter.

Thank you!