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NUCLEAR MATTER EQUATION OF STATE, INCOMPRESSIBILITY AND PROTON RADIOACTIVITY

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A mean field calculation is carried out to obtain the equation of state (EoS) of nuclear matter from a density dependent M3Y interaction (DDM3Y). The constants of density dependence of the effective interaction are obtained by reproducing the saturation energy per nucleon and the saturation density of the symmetric nuclear matter (SNM). In this work, the energy variation of the exchange potential is treated properly in the negative energy domain of nuclear matter in contrast to an earlier work where it was assumed to vary negligibly inside nuclear fluid. The EoS of SNM, thus obtained, is not only free from the superluminosity problem but also provides good estimate of nuclear incompressibility. The DDM3Y, whose density dependence is determined from nuclear matter calculation, provides excellent description for proton radioactivity.

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1. Introduction

High baryonic density behaviour of nuclear matter has recently received new impetus with the new accelerator facility (FAIR) coming up at GSI Darmstadt. The stiffness of a nuclear EoS is characterised by nuclear incompressibility [1] which can be extracted experimentally. Nuclear incompressibility [2,3] also determines the velocity of sound in nuclear medium for predictions of shock wave generation and propagation. The EoS is of fundamental importance in the theories of nucleus–nucleus collisions at energies where the nuclear incompressibility K_0 comes into play as well as in the theories of supernova explosions [4]. A widely used experimental method is

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the determination of the nuclear incompressibility from the observed giant monopole resonances (GMR) [5, 6]. Other recent experimental determinations are based upon the production of hard photons in heavy ion collisions [7] and from isoscalar giant dipole resonances (ISGDR) [8–10]. From the experimental data of isoscalar giant monopole resonance (ISGMR) conclusion can be drawn that $K_0 \approx 240 \pm 20$ MeV [11]. The general theoretical observation is that the non-relativistic [12] and the relativistic [13] mean field models [14] predict for the bulk incompressibility for the SNM, K_0, values which are significantly different from one another, *viz.* $\approx$ 220–235 MeV and $\approx$ 250–270 MeV, respectively. Theoretical EoS for the SNM that predict higher K_0 values $\approx$ 300 MeV are often called "stiff" EoS whereas those EoS which predict smaller K_0 values $\approx$ 200 MeV are termed as "soft" EoS.

In the present work, we show that the theoretical description of nuclear matter using the density dependent M3Y-Reid–Elliott effective interaction [15, 16] gives a value of nuclear incompressibility which is in excellent agreement with values extracted from experiments. The velocity of sound does not become superluminous since the energy dependence is treated properly for the negative energy domain of nuclear matter. The microscopic proton–nucleus interaction potential is obtained by folding the density of the nucleus with DDM3Y effective interaction whose density dependence is determined completely from the nuclear matter calculations. The quantum mechanical tunneling probability is calculated within the WKB framework using these nuclear potentials. These calculations provide reasonable estimates for the observed proton radioactivity lifetimes.

2. The nuclear equation of state for symmetric nuclear matter

In the present work, density dependence of the effective interaction, DDM3Y, is completely determined from nuclear matter calculations. The equilibrium density of the nuclear matter is determined by minimizing the energy per nucleon. In contrast to our earlier calculations for the nuclear EoS where the energy dependence of the zero range potential was treated as fixed at a value corresponding to the equilibrium energy per nucleon ϵ_0 [17] and assumed to vary negligibly with ϵ inside nuclear fluid, in the present calculations the energy variation of the zero range potential is treated more accurately by allowing it to vary freely but only with the kinetic energy part $\epsilon^{\rm kin}$ of the energy per nucleon ϵ over the entire range of ϵ. This is not only more plausible, but also yields excellent result for the incompressibility K_0 of the SNM which no more suffers from the superluminosity problem associated with non-relativistic EoS [1–3, 12, 14, 17].

The constants of density dependence are determined by reproducing the saturation conditions. It is worthwhile to mention here that due to attractive character of the M3Y forces the saturation condition for cold nuclear matter is not fulfilled. However, the realistic description of nuclear matter properties can be obtained with this density dependent M3Y effective interaction. Therefore, the constants of density dependence have been obtained by reproducing the saturation energy per nucleon and the saturation nucleonic density of the cold SNM. Based on the Hartree or mean field assumption and using the DDM3Y interaction, the expression for the energy per nucleon for symmetric nuclear matter ϵ is given by

$$\epsilon = \left[\frac{3\hbar^2 k_{\rm F}^2}{10m}\right] + \left[\frac{\rho J_{v00} C(1-\beta\rho^n)}{2}\right], \tag{1}$$

where Fermi momentum $k_{\rm F} = (1.5\pi^2\rho)^{\frac{1}{3}}$, m is the nucleonic mass equal to $938.91897\,{\rm MeV}/c^2$ and J_{v00} represents the volume integral of the isoscalar part of the M3Y interaction supplemented by the zero-range potential having the form

$$\begin{aligned} J_{v00} &= J_{v00}(\epsilon^{\rm kin}) = \int\int\int t_{00}^{\rm M3Y}(s,\epsilon) d^3s \\ &= 7999\frac{4\pi}{4^3} - 2134\frac{4\pi}{2.5^3} + J_{00}(1-\alpha\epsilon^{\rm kin})\ [{\rm MeV\ fm^3}], \end{aligned} \tag{2}$$

where $J_{00} = -276$ [MeV fm^3] and $\epsilon^{\rm kin} = 3\hbar^2 k_{\rm F}^2/(10m)$ is the kinetic energy part of the energy per nucleon ϵ given by Eq. (1).

The isoscalar $t_{00}^{\rm M3Y}$ and the isovector $t_{01}^{\rm M3Y}$ components of M3Y interaction potentials [16] supplemented by zero range potentials are given by $t_{00}^{\rm M3Y}(s,\epsilon) = 7999\frac{\exp(-4s)}{4s} - 2134\frac{\exp(-2.5s)}{2.5s} - 276(1-\alpha\epsilon)\delta(s)$ and $t_{01}^{\rm M3Y}(s,\epsilon) = -4886\frac{\exp(-4s)}{4s} + 1176\frac{\exp(-2.5s)}{2.5s} + 228(1-\alpha\epsilon)\delta(s)$, respectively, where the energy dependence parameter $\alpha = 0.005/{\rm MeV}$. The DDM3Y effective NN interaction is given by $v_{0i}(s,\rho,\epsilon) = t_{0i}^{\rm M3Y}(s,\epsilon)g(\rho)$ where the density dependence $g(\rho) = C(1-\beta\rho^n)$ and the constants C and β of density dependence have been obtained from the saturation condition $\frac{\partial\epsilon}{\partial\rho} = 0$ at $\rho = \rho_0$ and $\epsilon = \epsilon_0$ where ρ_0 and ϵ_0 are the saturation density and the saturation energy per nucleon, respectively. Eq. (1) can be differentiated with respect to ρ to yield equation

$$\frac{\partial\epsilon}{\partial\rho} = \left[\frac{\hbar^2 k_{\rm F}^2}{5m\rho}\right] + \frac{J_{v00}C}{2}[1-(n+1)\beta\rho^n] - \alpha J_{00}C[1-\beta\rho^n]\left[\frac{\hbar^2 k_{\rm F}^2}{10m}\right]. \tag{3}$$

The equilibrium density of the cold SNM is determined from the saturation condition. Then Eq. (1) and Eq. (3) with the saturation condition $\frac{\partial\epsilon}{\partial\rho} = 0$ can be solved simultaneously for fixed values of the saturation energy per nucleon ϵ_0 and the saturation density ρ_0 of the cold SNM to obtain the values of β and C. The constants of density dependence β and C, thus obtained, are given by

$$\beta = \frac{\left[(1-p) + \left(q - \frac{3q}{p}\right)\right] \rho_0^{-n}}{\left[\left(3n+1\right) - (n+1)p + \left(q - \frac{3q}{p}\right)\right]}, \tag{4}$$

where $p = [10m\epsilon_0]/[\hbar^2 k_{\mathrm{F}_0}^2]$, $q = 2\alpha\epsilon_0 J_{00}/J_{v00}^0$ with $J_{v00}^0 = J_{v00}(\epsilon_0^{\mathrm{kin}})$ which means J_{v00} at $\epsilon^{\mathrm{kin}} = \epsilon_0^{\mathrm{kin}}$, the kinetic energy part of the saturation energy per nucleon of SNM, $k_{F_0} = [1.5\pi^2\rho_0]^{1/3}$ and

$$C = -\frac{[2\hbar^2 k_{F_0}^2]}{5mJ_{v00}^0\rho_0\left[1 - (n+1)\beta\rho_0^n - \frac{q\hbar^2 k_{F_0}^2(1-\beta\rho_0^n)}{10m\epsilon_0}\right]}, \tag{5}$$

respectively. It is quite obvious that the constants of density dependence C and β obtained by this method depend on the saturation energy per nucleon ϵ_0, the saturation density ρ_0, the index n of the density dependent part and on the strengths of the M3Y interaction through the volume integral J_{v00}^0.

3. The incompressibility of symmetric nuclear matter

The incompressibility or the compression modulus of symmetric nuclear matter, which is a measure of the curvature of an EoS at saturation density and defined as $k_{\mathrm{F}}^2 \frac{\partial^2\epsilon}{\partial k_{\mathrm{F}}^2} |_{k_{\mathrm{F}}=k_{\mathrm{F}_0}}$, measures the stiffness of an EoS. The $\frac{\partial^2\epsilon}{\partial\rho^2}$ is given by

$$\begin{aligned}\frac{\partial^2\epsilon}{\partial\rho^2} &= \left[-\frac{\hbar^2 k_{\mathrm{F}}^2}{15m\rho^2}\right] - \left[\frac{J_{v00}Cn(n+1)\beta\rho^{n-1}}{2}\right] \\ &\quad -\alpha J_{00}C[1-(n+1)\beta\rho^n]\left[\frac{\hbar^2 k_{\mathrm{F}}^2}{5m\rho}\right] + \left[\frac{\alpha J_{00}C(1-\beta\rho^n)\hbar^2 k_{\mathrm{F}}^2}{30m\rho}\right]\end{aligned} \tag{6}$$

and therefore the incompressibility of the cold SNM can be theoretically obtained as

$$\begin{aligned}K_0 = k_{\mathrm{F}}^2 \frac{\partial^2\epsilon}{\partial k_{\mathrm{F}}^2}\bigg|_{k_{\mathrm{F}}=k_{F_0}} &= 9\rho^2 \frac{\partial^2\epsilon}{\partial\rho^2}\bigg|_{\rho=\rho_0} = \left[-\frac{3\hbar^2 k_{F_0}^2}{5m}\right] - \left[\frac{9J_{v00}^0 Cn(n+1)\beta\rho_0^{n+1}}{2}\right] \\ &- 9\alpha J_{00}C[1-(n+1)\beta\rho_0^n]\left[\frac{\rho_0\hbar^2 k_{F_0}^2}{5m}\right] + \left[\frac{3\rho_0\alpha J_{00}C(1-\beta\rho_0^n)\hbar^2 k_{F_0}^2}{10m}\right].\end{aligned} \tag{7}$$

The calculations are performed using the values of the saturation density $\rho_0 = 0.1533\,\mathrm{fm}^{-3}$ [2] and the saturation energy per nucleon $\epsilon_0 = -15.26\,\mathrm{MeV}$ [18] for the SNM obtained from the co-efficient of the volume term of Bethe–Weizsäcker mass formula [19, 20] which is evaluated by fitting the recent experimental and estimated atomic mass excesses from Audi–Wapstra–Thibault atomic mass table [21] by minimizing the mean square deviation incorporating correction for the electronic binding energy [22]. In a similar recent work, including surface symmetry energy term, Wigner term, shell correction and proton form factor correction to Coulomb energy also, a_v turns out to be 15.4496 MeV [23] ($a_v = 14.8497$ MeV when A^0 and $A^{1/3}$ terms are also included). Using the usual values of $\alpha = 0.005\,\mathrm{MeV}^{-1}$ for the parameter of energy dependence of the zero range potential and $n = 2/3$, the values obtained for the constants of density dependence C and β and the SNM incompressibility K_0 are 2.2497, 1.5934 fm^2 and 274.7 MeV, respectively. The saturation energy per nucleon is the volume energy coefficient and the value of -15.26 ± 0.52 MeV covers, more or less, the entire range of values obtained for a_v for which now the values of $C = 2.2497\pm0.0420$, $\beta = 1.5934\pm0.0085\,\mathrm{fm}^2$ and the SNM incompressibility $K_0 = 274.7 \pm 7.4$ MeV.

The theoretical estimate K_0 of the incompressibility of infinite SNM obtained from present approach using DDM3Y is about 270 MeV. The theoretical estimate of K_0 from the refractive α–nucleus scattering is about 240–270 MeV [24, 25] and that by infinite nuclear matter model (INM) [26] claims a well defined and stable value of $K_0 = 288 \pm 20$ MeV and present theoretical estimate is in reasonably close agreement with the value obtained by INM which rules out any values lower than 200 MeV. Present estimate for the incompressibility K_0 of the infinite SNM is in good agreement with the experimental value of $K_0 = 300\pm25$ MeV obtained from the giant monopole resonance (GMR) [5] and with the recent experimental determination of K_0 based upon the production of hard photons in heavy ion collisions which led to the experimental estimate of $K_0 = 290 \pm 50$ MeV [7]. However, the experimental values of K_0 extracted from the isoscalar giant dipole resonance (ISGDR) are claimed to be smaller [10]. The present non-relativistic mean field model estimate for the nuclear incompressibility K_0 for SNM using DDM3Y interaction is rather close to the theoretical estimates obtained using relativistic mean field models and close to the lower limit of the older experimental values [5] and close to the upper limit of the recent values [6] extracted from experiments.

Considering the status of experimental determination of the SNM incompressibility from data on the compression modes ISGMR and ISGDR of nuclei it can be inferred [11] that due to violations of self consistency in HF–RPA calculations of the strength functions of giant resonances result in shifts in the calculated values of the centroid energies which may be larger in magnitude than the current experimental uncertainties. In fact, the prediction of K_0 lying in the range of 210–220 MeV was due to the use of a not fully self-consistent Skyrme calculations [11]. Correcting for this drawback, Skyrme parmetrizations of SLy4 type predict K_0 values in the range of 230–240 MeV [11]. Moreover, it is possible to build *bona fide* Skyrme forces so that the SNM incompressibility is close to the relativistic value, namely 250–270 MeV. Therefore, from the ISGMR experimental data the conclusion can be drawn that $K_0 \approx 240 \pm 20$ MeV. The ISGDR data tend to point to lower values [8–10] for K_0. However, there is consensus that the extraction of K_0 is in this case more problematic for various reasons. In particular, the maximum cross-section for ISGDR decreases very strongly at high excitation energy and may drop below the current experimental sensitivity for excitation energies [11] above 30 and 26 MeV for ^{116}Sn and ^{208}Pb, respectively. The present value of 274.7 ± 7.4 MeV for the incompressibility K_0 of SNM obtained using DDM3Y interaction is, therefore, an excellent theoretical result.

The constant of density dependence $\beta = 1.5934 \pm 0.0085\,\mathrm{fm}^2$, which has the dimension of cross-section for $n = 2/3$, can be interpreted as the isospin averaged effective nucleon–nucleon interaction cross-section in ground state symmetric nuclear medium. For a nucleon in ground state nuclear matter $k_{\mathrm{F}} \approx 1.3\,\mathrm{fm}^{-1}$ and $q_0 \sim \hbar k_{\mathrm{F}} c \approx 260$ MeV and the present result for the "in medium" effective cross-section is reasonably close to the value obtained from a rigorous Dirac–Brueckner–Hartree–Fock [27] calculations corresponding to these k_{F} and q_0 values, which is ≈ 12 mb. Using the value of the constant of density dependence $\beta = 1.5934 \pm 0.0085\,\mathrm{fm}^2$ corresponding to the standard value of the parameter $n = 2/3$ along with the nucleonic density of $0.1533\,\mathrm{fm}^{-3}$, the value obtained for the nuclear mean free path λ is about 4 fm which is in excellent agreement [28] with that obtained using another method.

4. Proton radioactivity lifetimes using effective interaction whose density dependence is determined from nuclear matter calculation

Microscopic proton–nucleus interaction potentials $V_{\mathrm{N}}(R)$ are obtained by single folding the density of the nucleus with M3Y effective interaction supplemented by a zero-range potential for exchange along with the density dependence:

$$V_{\mathrm{N}}(R) = \int \rho(\vec{r}) v_{00}[|\vec{r} - \vec{R}|]\, d^3r \,, \tag{8}$$

where $\vec{R}$ and $\vec{r}$ are, respectively, the co-ordinates of the emitted proton and a nucleon belonging to the residual daughter nucleus with respect to its centre. The density distribution function ρ used for the daughter nucleus, has been chosen to be of the spherically symmetric form given by

$$\rho(r) = \rho_0 / [1 + \exp((r - c)/a)] \,, \tag{9}$$

where $c = r_\rho(1-\pi^2a^2/3r_\rho^2)$, $r_\rho = 1.13A_\mathrm{d}^{1/3}$, $a = 0.54$ fm and the value of ρ_0 is fixed by equating the volume integral of the density distribution function to the mass number A_d of the residual daughter nucleus. The distance between a nucleon belonging to the residual daughter nucleus and the emitted proton is $s = |\vec{r} - \vec{R}|$ while the interaction potential between any such two nucleons $v_{00}(s)$ appearing in Eq. (8) is given by the DDM3Y effective interaction. The total interaction energy between the proton and the residual daughter nucleus $E(R) = V_\mathrm{N}(R)+V_\mathrm{C}(R)+\hbar^2l(l+1)/(2\mu R^2)$, the sum of the nuclear interaction energy, the Coulomb interaction energy and the centrifugal barrier. Here $\mu = M_pM_\mathrm{d}/M_A$ is the reduced mass, M_p, M_d and M_A are the masses of the proton, the daughter nucleus and the parent nucleus, respectively, all measured in the units of MeV$/c^2$. Assuming spherical charge distribution for the residual daughter nucleus, the proton–nucleus Coulomb interaction potential $V_\mathrm{C}(R)=(Z_\mathrm{d}e^2/2R_\mathrm{c})[3 - (R/R_\mathrm{c})^2]$ for $R \leq R_\mathrm{c}$, $= Z_\mathrm{d}e^2/R$ otherwise, where Z_d is the atomic number of the daughter nucleus. The touching radial separation R_c between the proton and the daughter nucleus is given by $R_\mathrm{c} = c_p + c_d$ where c_p and c_d have been obtained using expression for c provided after Eq. (9). The energetics allow spontaneous emission of protons only if the released energy $Q = [M_A - (M_p + M_\mathrm{d})]c^2$ is a positive quantity.

In the present work, the half life of the parent nucleus decaying via proton emission is calculated using the WKB barrier penetration probability. The assault frequency ν is obtained from the zero point vibration energy $E_v = \frac{1}{2}\hbar\omega = \frac{1}{2}h\nu$. The decay half life T of the parent nucleus (A, Z) into a proton and a daughter $(A_\mathrm{d}, Z_\mathrm{d})$ is given by

$$T = [(h \ln 2)/(2E_v)][1 + \exp(K)]\,, \tag{10}$$

where the action integral K within the WKB approximation is given by $K = (2/\hbar)\int_{R_a}^{R_b}[2\mu(E(R) - E_v - Q)]^{1/2}dR$ where R_a and R_b are its 2nd and 3rd turning points determined from the equations $E(R_a) = Q+E_v = E(R_b)$. The isovector or the symmetry component of the DDM3Y folded potential $V_\mathrm{N}^\mathrm{Lane}(R)$ [29] has been added to the isoscalar part of the folded potential whose results have already been presented in Table I. The nuclear potential $V_\mathrm{N}(R)$ of Eq. (8), therefore, has been replaced by $V_\mathrm{N}(R)+V_\mathrm{N}^\mathrm{Lane}(R)$ [30] where

$$V_\mathrm{N}^\mathrm{Lane}(R)=\int\int[\rho_{1n}(\vec{r_1})-\rho_{1p}(\vec{r_1})][\rho_{2n}(\vec{r_2})-\rho_{2p}(\vec{r_2})]v_{01}[|\vec{r_2}-\vec{r_1}+\vec{R}|]d^3r_1d^3r_2\,, \tag{11}$$

where the subscripts 1 and 2 denote the daughter and the emitted nuclei, respectively, while the subscripts n and p denote neutron and proton densities, respectively. With simple assumption that $\rho_{1p}=[\frac{Z_\mathrm{d}}{A_\mathrm{d}}]\rho$ and $\rho_{1n}=[\frac{(A_\mathrm{d}-Z_\mathrm{d})}{A_\mathrm{d}}]\rho$, and for the emitted particle being proton $\rho_{2n}(\vec{r_2})-\rho_{2p}(\vec{r_2})=-\rho_2(\vec{r_2})=-\delta(\vec{r_2})$,

TABLE I

Comparison between experimentally measured and theoretically calculated half lives of spherical proton emitters. The experimental Q values, half lives and l values are from Ref. [35]. The results of the present calculations using the isoscalar and isovector components of DDM3Y folded potentials are compared with the experimental values. Experimental errors in Q values [35] and corresponding errors in calculated half lives are inside parentheses.

Parent ^AZ	l $\hbar$	Q MeV	1^{st} tpt R_1[fm]	2^{nd} tpt R_a[fm]	3^{rd} tpt R_b[fm]	Exp. $\log_{10} T(s)$	this work $\log_{10} T(s)$
^{105}Sb	2	0.491(15)	1.43	6.69	134.30	$2.049^{+0.058}_{-0.067}$	1.90(45)
^{109}I	2	0.829(3)	1.42	6.78	83.29	$-3.987^{+0.020}_{-0.022}$	−4.31(5)
^{112}Cs	2	0.824(7)	1.44	6.81	88.61	$-3.301^{+0.079}_{-0.097}$	−3.21(11)
^{113}Cs	2	0.978(3)	1.44	6.84	73.45	$-4.777^{+0.018}_{-0.019}$	−5.61(4)
^{145}Tm	5	1.753(10)	3.20	6.63	56.27	$-5.409^{+0.109}_{-0.146}$	−5.28(7)
^{147}Tm	5	1.071(3)	3.18	6.63	88.65	$0.591^{+0.125}_{-0.175}$	0.83(4)
^{147}Tm*	2	1.139(5)	1.44	7.28	78.97	$-3.444^{+0.046}_{-0.051}$	−3.46(6)
^{150}Lu	5	1.283(4)	3.21	6.67	78.23	$-1.180^{+0.055}_{-0.064}$	−0.74(4)
^{150}Lu*	2	1.317(15)	1.45	7.33	71.79	$-4.523^{+0.620}_{-0.301}$	−4.46(15)
^{151}Lu	5	1.255(3)	3.21	6.69	78.41	$-0.896^{+0.011}_{-0.012}$	−0.82(4)
^{151}Lu*	2	1.332(10)	1.46	7.35	69.63	$-4.796^{+0.026}_{-0.027}$	−4.96(10)
^{155}Ta	5	1.791(10)	3.21	6.78	57.83	$-4.921^{+0.125}_{-0.125}$	−4.80(7)
^{156}Ta	2	1.028(5)	1.47	7.37	94.18	$-0.620^{+0.082}_{-0.101}$	−0.47(8)
^{156}Ta*	5	1.130(8)	3.21	6.76	90.30	$0.949^{+0.100}_{-0.129}$	1.50(10)
^{157}Ta	0	0.947(7)	0.00	7.55	98.95	$-0.523^{+0.135}_{-0.198}$	−0.51(12)
^{160}Re	2	1.284(6)	1.45	7.43	77.67	$-3.046^{+0.075}_{-0.056}$	−3.08(7)
^{161}Re	0	1.214(6)	0.00	7.62	79.33	$-3.432^{+0.045}_{-0.049}$	−3.53(7)
^{161}Re*	5	1.338(7)	3.22	6.84	77.47	$-0.488^{+0.056}_{-0.065}$	−0.75(8)
^{164}Ir	5	1.844(9)	3.20	6.91	59.97	$-3.959^{+0.190}_{-0.139}$	−4.08(6)
^{165}Ir*	5	1.733(7)	3.21	6.93	62.35	$-3.469^{+0.082}_{-0.100}$	−3.67(5)
^{166}Ir	2	1.168(8)	1.47	7.49	87.51	$-0.824^{+0.166}_{-0.273}$	−1.19(10)
^{166}Ir*	5	1.340(8)	3.22	6.91	80.67	$-0.076^{+0.125}_{-0.176}$	0.06(9)
^{167}Ir	0	1.086(6)	0.00	7.68	91.08	$-0.959^{+0.024}_{-0.025}$	−1.35(8)
^{167}Ir*	5	1.261(7)	3.22	6.92	83.82	$0.875^{+0.098}_{-0.127}$	0.54(8)
^{171}Au	0	1.469(17)	0.00	7.74	69.09	$-4.770^{+0.185}_{-0.151}$	−5.10(16)
^{171}Au*	5	1.718(6)	3.21	7.01	64.25	$-2.654^{+0.054}_{-0.060}$	−3.19(5)
^{177}Tl	0	1.180(20)	0.00	7.76	88.25	$-1.174^{+0.191}_{-0.349}$	−1.44(26)
^{177}Tl*	5	1.986(10)	3.22	7.10	57.43	$-3.347^{+0.095}_{-0.122}$	−4.64(6)
^{185}Bi	0	1.624(16)	0.00	7.91	65.71	$-4.229^{+0.068}_{-0.081}$	−5.53(14)

(*) in the experimental Q values denotes the isomeric state.

the Lane potential becomes $V_N^{\rm Lane}(R) = -[\frac{(A_{\rm d}-2Z_{\rm d})}{A_{\rm d}}] \int \rho(\vec{r}) v_{01}\ [|\vec{r}-\vec{R}|] d^3r$ where $v_{01}(s) = t_{01}^{\rm M3Y}(s,\epsilon) g(\rho)$ with the isovector part $t_{01}^{\rm M3Y} = -[4886\frac{e^{-4s}}{4s} - 1176\frac{e^{-2.5s}}{2.5s}] + 228(1-0.005Qm/\mu)\,\delta(s)$. The inclusion of this Lane potential causes insignificant changes in the lifetimes.

Thus half lives of the decays of spherical nuclei away from proton drip line by proton emissions are estimated theoretically. The half life of a parent nucleus decaying via proton emission is calculated using the WKB barrier penetration probability. The WKB method is found to be quite satisfactory for the α decay half life calculations of heavy [31] and superheavy elements [32] and somewhat better than the S-matrix method [33]. For the present calculations, the zero point vibration energies used here are given by Eq. (5) of Ref. [34] extended to protons and the experimental Q values [35] are used. Present calculations for half lives have been performed using $C = 2.2497$ and $\beta = 1.5934\,{\rm fm}^2$ obtained here from the nuclear matter calculations assuming kinetic energy dependence of zero range potential and presented in Table I. The agreement of the present calculations with a wide range of experimental data for the proton radioactivity lifetimes is reasonable.

5. Summary and conclusion

A mean field calculation is carried out to obtain the equation of state of nuclear matter from a DDM3Y interaction. The microscopic proton–nucleus potential is obtained by folding the DDM3Y effective interaction with the density of interacting nucleus. The energy per nucleon is minimized to obtain ground state of the symmetric nuclear matter. The constants of density dependence of the effective interaction are obtained by reproducing the saturation energy per nucleon and the saturation density of SNM. In this work the energy variation of the exchange potential is treated properly in the negative energy domain of nuclear matter. The EoS of SNM, thus obtained, is free from the superluminosity problem encountered in some previous prescriptions. Moreover, the result of the present calculation for the compression modulus for the infinite symmetric nuclear matter is in better agreement with that extracted from experiments. The results of the present calculations using single folded microscopic potentials for the proton-radioactivity lifetimes are in good agreement over a wide range of experimental data. With the energies and interaction rates foreseen at FAIR, the compressed baryonic matter will create highest baryon densities in nucleus–nucleus collisions to explore the properties of superdense baryonic matter and the in-medium modifications of hadrons.

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