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Relativistic Mean Field Theory with the Pion in Finite Nuclei

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Abstract

We study the possible occurrence of finite pion mean field in finite nuclei in the relativistic mean field (RMF) theory. We are led to this conjecture due to the findings of the essential role of the pion in few-body systems and the (p,n) experiments for spin-isospin excitations of medium and heavy nuclei. We calculate explicitly various N=Z closed shell nuclei with finite pion mean field in the RMF theory with the standard parameter set and the pion-nucleon coupling in free space. The finite mean field breaks the parity symmetry of the intrinsic single particle states. We demonstrate the actual occurrence and the property of the finite pion mean field, whose qualitative consequence on various observables is also briefly discussed.

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§1 Introduction

The pion is conjectured by Yukawa as the mediator of nucleon-nucleon interaction¹⁾ It is identified then as the Nambu-Goldstone mode when the underlying symmetry, the chiral symmetry, is spontaneously broken²⁾ It has the property of $J^\pi = 0^-$ spin-parity and $T = 1$ isospin, which provides a rich phenomena in hadron and nuclear physics The pion plays the central role in hadron physics, particularly for the low energy phenomena, and now the chiral perturbation approach with the pion as the essential degree of freedom is the powerful tool for the study of hadron properties and their collisions

On the other hand, since the discovery of the shell model, nuclear structure has been studied in terms of the single particle basis³⁾ There the single particle orbits are obtained under the parity conserved mean field given by the central and spin-orbit potential However, when a virtual pion is emitted and absorbed by a nucleon in a good parity orbit, the nucleon jumps from the single particle state to another with opposite parity and spin-flip due to the pion-nucleon coupling Therefore, to incorporate the effect of the correlation caused by the pion exchange interaction in the parity-conserved single-particle space, we must treat the higher configurations like 2p-2h (2 particle-2 hole) states crossing over major shells To avoid such complications we renormalize the central and the spin-orbit interactions in nuclei to take into account the strong correlations caused by the pion in the restricted model space It means that we assume a model space in which the parity is conserved and then define an effective interaction acting between nucleons in the model space Hence, the pion is not treated explicitly in nuclear physics One of the main motivation of the present study is to expand the model space so that the pion plays its important role explicitly in order to see the effect of the correlation induced by the pion on nuclear structure

As for the importance of the pion, we remind ourselves of the findings of few-body systems, where they are solved rigorously without the restriction of model space with a realistic nucleon-nucleon interaction which has the short-range repulsive core and the tensor force explicitly⁴⁾⁻⁶⁾ The calculations are performed in the non-relativistic framework and hence the pion shows up mainly as the tensor force, which is much stronger than the central force in the one pion exchange interaction There are many extensive calculations using the variational principle and sophisticated numerical techniques and the calculated results compare with experiments very well⁵⁾⁻⁷⁾ The calculated results demonstrate a dominant role of the tensor force in the few-body system Almost a half of the attraction is caused by the tensor force, and hence the pion exchange interaction in the alpha particle⁴⁾ The recent variational calculations of the Argonne group up to $A=8$ system also demonstrate the dominant role of the pion, which amounts to 70-80% in two-body interaction part in $3 \leq A \leq 8$

systems and further the importance of the three-body interaction, which also originates from the pion⁶⁾

There are several experimental data demanding the important role of the pion in medium and heavy nuclei. The (p,n) reactions on medium and heavy nuclei demonstrate that slightly more than a half of the Gamow-Teller ($\sigma\tau$) strength is carried by 1p-1h excitations, while most of the rest is interpreted as carried by 2p-2h excitations due to the coupling with the 1p-1h states by the strong tensor force, which distributes the strength up to at least 50 MeV in excitation energy⁸⁾. A further dramatic result is the ratio of the longitudinal and the transverse spin responses being close to one, while the strong pionic correlations ought to provide a large enhancement in the longitudinal channel^{9), 10)}. This experimental result seems to demand a special improvement on the treatment of the pion in nuclear spin response. In addition, electron elastic scattering with a nucleus and (e,e'p) reactions do not see the full single particle strength^{11), 12)}. A theoretical study claims that a large fraction of the missing strength is caused by the tensor correlation and the short-range repulsion¹³⁾.

All these facts seem to indicate the importance of the correlation induced by the pion in the description of nuclear many-body system. Hence, we would like to construct a framework in which the pionic effect is treated explicitly under the mean field approximation for medium and heavy nuclei. To this end, we break the parity, spin, and isospin symmetries of single particle states to treat the pion in the mean field approximation. This trial may oppose to the common sense that pion condensation does not happen in nuclear matter at the saturation density^{14), 15)}. This fact does not mean, however, that the pion mean field vanishes in finite nuclear system. The pion is a pseudoscalar meson, it couples with a nucleon through the $\vec{\sigma} \cdot \vec{\nabla}$ coupling, and therefore the source term of the pion field needs the parity-mixing¹⁶⁾ and the spin-density modulation¹⁷⁾. In infinite matter at a high density region, we have to provide this spin-density modulation (ALS structure) for the pion to work, which costs large energy¹⁷⁾. In finite nuclear system, we have the nuclear surface and automatically there arises the spin-density modulation necessary for the pion to work. Hence, if the pion mean field turns out finite in nuclear system, the pionic correlations utilize the change of the spin-density at the surface.

In this paper, we would like to study the effect of the pion mean field, which incorporates the pionic correlation to the nuclear mean field. We therefore introduce the pion term in the relativistic mean field (RMF) Lagrangian¹⁸⁾ for the calculation of finite nuclei. We shall take the RMF Lagrangian of the sigma, omega, and rho meson parts with a parameter set, with which a systematic study of nuclei in the entire mass region has been performed with success¹⁹⁾. However, the pion was not usually included in this Lagrangian. We would like to study if the pion mean field becomes finite and how the pion mean field develops in the

N=Z closed-shell configuration nuclei with the mass number, when we include the pion term in addition to the present RMF Lagrangian. We shall show how strong is the pion nucleon coupling for the development of the finite pion mean field and what is the consequence of the finite pion mean field on the physical observables. In this paper, we show the RMF formalism with the pion mean field in Section 2 and we present the calculated results for N=Z closed-shell configurations in Section 3. We discuss the structure of the mixed-parity intrinsic wave function and the qualitative consequence on various observables of the finite pion mean field in Section 4. The summary of the present study is given in Section 5 together with the statements for the further work in the direction of the establishment of a mean field theory with the finite pion mean field.

§2 Relativistic mean field theory with the pion

We start with writing the relativistic meson-nucleon Lagrangian density, which naturally includes the pion term, π ,

$$\begin{aligned} \mathcal{L} = & \bar{\psi} [i\gamma_\mu \partial^\mu - M - g_\pi \gamma_5 \gamma_\mu \tau^a \partial^\mu \pi^a - g_\sigma \sigma - g_\omega \gamma_\mu \omega^\mu - g_\rho \gamma_\mu \tau^a \rho^{a\mu} - e\gamma_\mu \frac{(1-\tau_3)}{2} A^\mu] \psi \\ & + \frac{1}{2} \partial_\mu \pi^a \partial^\mu \pi^a - \frac{1}{2} m_\pi^2 \pi^a \pi^a + \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \frac{1}{2} m_\sigma^2 \sigma^2 - \frac{1}{3} g_2 \sigma^3 - \frac{1}{4} g_3 \sigma^4 \\ & - \frac{1}{4} H_{\mu\nu} H^{\mu\nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu + \frac{1}{4} c_3 (\omega_\mu \omega^\mu)^2 \\ & - \frac{1}{4} G_{\mu\nu}^a G^{a\mu\nu} + \frac{1}{2} m_\rho^2 \rho_\mu^a \rho^{a\mu} - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}, \end{aligned} \quad (2.1)$$

where the field tensors H , G and F for the vector fields are defined through¹⁸⁾,

$$H_{\mu\nu} = \partial_\mu \omega_\nu - \partial_\nu \omega_\mu, \quad (2.2a)$$

$$G_{\mu\nu}^a = \partial_\mu \rho_\nu^a - \partial_\nu \rho_\mu^a - 2g_\rho \epsilon^{abc} \rho_\mu^b \rho_\nu^c, \quad (2.2b)$$

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \quad (2.2c)$$

The pion term couples with the nucleon through the pseudovector coupling. We take here all the terms used in Ref. 19). Here, σ denotes the scalar meson, ω the vector meson and ρ the isovector-vector meson. The photon is denoted by A . It should be noted that the coupling of the spin and orbital angular momenta automatically appears in the relativistic formalism.

In the present study, we do not include this contribution of the coupling of the pion with the Δ state for simplicity of discussion. There should be the coupling of the pion with the Δ state and this contribution will act constructively for finite pion mean field. We do not

include here the tensor coupling term of the ρ meson either in order to concentrate on the pion degree of freedom. All these effects will be worked out in the near future²⁰⁾

We then assume that the expectation value of the pion field is finite. In addition, the nuclear system is isospin-singlet due to the choice of $N=Z$ nuclei and hence the self-consistent Hartree-Fock Hamiltonian is invariant under the isospin rotation²¹⁾. We choose therefore the finite value of the pion mean field as the z-component, i.e. $a = 0$, without the loss of generality. We write the finite pion mean field as π without the isospin suffix.

We write here only the equations of motion for the nucleon and the pion explicitly. They are written as,

$$[i\gamma^\mu\partial_\mu - M - g_\pi\vec{\nabla}\pi\gamma_5\vec{\gamma}\tau_0]\psi = 0, \quad (2.3)$$

for nucleon and

$$(\vec{\nabla}^2 - m_\pi^2)\pi = -g_\pi\vec{\nabla}\langle\bar{\psi}\gamma_5\vec{\gamma}\tau_0\psi\rangle, \quad (2.4)$$

for pion. Here, the bracket $\langle \rangle$ above denotes the ground state expectation value. The other mesons follow similar equations of motion as above, whose explicit forms are written in ref. 19)

The equations (2.3) and (2.4) tell the structure of the pion mean field. These equations tell the reason why we have not included the pion mean field until now. If single particle states in a mean field have good-parity, the source term of the pion field becomes zero. Hence, we ought to break parity for the construction of the single particle states. The violation of the parity is caused by the pion term in the above Dirac equation for nucleons. The pion field is finite when the source term breaks the parity. The pion field is enhanced by a spatial change of the source term due to the derivative of the coupling of the pion and the nucleon. When the pion field is finite in (2.3), the nucleon single particle state breaks the parity, which provides again the pion source term finite. The self-consistency condition provides the converged solution to the above equations. If the pion mean field becomes finite, the parity is broken and it is interpreted as the pion condensation occurs in the finite nuclei.

Hence, the single particle state is expressed as

$$\psi_{njm} = \sum_{\kappa} \begin{pmatrix} iG_{n\kappa}\mathcal{Y}_{\kappa m} \\ F_{n\kappa}\mathcal{Y}_{\bar{\kappa} m} \end{pmatrix} = \begin{pmatrix} iG_{n\kappa}\mathcal{Y}_{\kappa m} + iG_{n\bar{\kappa}}\mathcal{Y}_{\bar{\kappa} m} \\ F_{n\kappa}\mathcal{Y}_{\bar{\kappa} m} + F_{n\bar{\kappa}}\mathcal{Y}_{\kappa m} \end{pmatrix} \quad (2.5)$$

Here, $\mathcal{Y}_{\kappa m}$ is the eigenfunction of the total angular momentum $\mathbf{j} = \mathbf{l} + \mathbf{s}$ and $\mathcal{Y}_{\bar{\kappa} m} = \vec{\sigma} \cdot \hat{r} \mathcal{Y}_{\kappa m}$. We assume the spherical symmetry (jm is a good quantum number) for the intrinsic state. G and F are the radial parts of the single particle wave function. We write here the intrinsic state, since the physical state has to be obtained by projection. The summation over κ means the parity mixing, where κ is $\kappa = -(l_\downarrow + 1)$ for $l_\downarrow = j - 1/2$ and $\kappa = l_\uparrow$ for $l_\uparrow = j + 1/2$. The calculational detail will be provided in the forthcoming publication²⁰⁾

§3 Numerical results with finite pion field

We show here the numerical results. We take the TM1 parameter set of Ref. 19) for all the parameters except for the pion-nucleon coupling. As for this we take the value of Bonn A potential²²⁾, which corresponds to taking $g_\pi = f_\pi/m_\pi$ and $f_\pi \sim 1$. We stress again here that we use all the terms given in the Lagrangian (2.1). This means that the saturation property is guaranteed and the bulk part of the nucleus tends to have the saturation density. Since we are especially interested in the occurrence of finite pion mean field and want to see its effect under the simplest condition, we neglect the Coulomb term. We calculate the N=Z closed shell nuclei as ^{12}C , ^{16}O , ^{40}Ca , ^{56}Ni , ^{80}Zr , ^{100}Sn and ^{164}Pb .

We show the results in Table I with and without pion mean field. With the present choice of the parameters, we find the actual occurrence of finite pion mean field for all the nuclei studied here. This is not a trivial result that the pion mean field becomes finite. As we will show later, the pion mean field become finite only when the pion-nucleon coupling constant is larger than the critical value. There we will see the dependence of the pion mean field on the pion-nucleon coupling strength together with a critical strength for the on-set of the pion mean field. Table I shows the total binding energy (BE), the kinetic energy (KE), the sigma energy (V_σ), omega energy (V_ω) and pion energy (V_π). The rho meson mean field vanishes for N=Z nuclei. The case with finite pion mean field provides larger total binding energy than the one without pion mean field. The pion energy is attractive, while all the other terms as the kinetic energy, the combined sigma and omega energy do not favor finite pion mean field. It means that a single particle state is brought up to the next major shell by the pionic correlation and, then, the kinetic energy becomes larger.

We show the mass number dependence of the pion energy per nucleon in Fig. 1. We mention here that the kinetic energy and the sigma and omega energies per particle are almost constant of the mass number and hence they are volume like as easily recognized from the numbers in Table I. We see, on the other hand, a peculiar behavior in the pion energy. The magnitudes of the pion energy are clearly separated into two groups. One group is large and the common feature is jj-closed shell nuclei, the magic number nuclei due to a larger spin-orbit partner (j-upper) being filled. The other group is small and they are LS-closed shell nuclei. The pion energy per nucleon for jj-closed shell nuclei decreases monotonically with the mass number. The rate of the decrease follow more strongly than $A^{-1/3}$. This means that the pion mean field energy behaves in proportion with the nuclear surface or even stronger than that. Hence, we use the word of surface pion condensation. Concerning the separation into two groups for the pion energy, LS-closed and jj-closed shell cases, we shall mention the possible reason later.

Table I The contributions of kinetic energy (KE), sigma (V_σ), omega (V_ω) and pion (V_π) energies to the total binding energy (BE) given in MeV. Many N=Z closed shell nuclei are calculated with and without finite pion mean field. ^{12}C , ^{56}Ni , ^{100}Sn and ^{164}Pb are the jj-closed shell nuclei and ^{16}O , ^{40}Ca and ^{80}Zr are the LS-closed shell nuclei.

	With π mean field					Without π mean field			
	BE	KE	V_σ	V_ω	V_π	BE	KE	V_σ	V_ω
^{12}C	116	228	-1387	1196	-108	98	197	-1376	1128
^{16}O	148	227	-1794	1490	-14	148	223	-1811	1498
^{40}Ca	431	576	-4914	4127	-81	427	556	-5003	4164
^{56}Ni	650	955	-7431	6316	-281	626	877	-7622	6340
^{80}Zr	931	1118	-10458	8751	-43	931	1106	-10519	8784
^{100}Sn	1231	1599	-13742	11629	-326	1214	1509	-14006	11689
^{164}Pb	2109	2520	-22913	19352	-409	2094	2409	-23352	19533

Fig. 1 The pion energy per nucleon as a function of the mass number in the log-log plot. There are two groups; one is for the jj-closed shell nuclei denoted by open circle and the other is for the LS-closed shell nuclei denoted by closed circle. The pion energy per nucleon for the jj-closed shell nuclei decreases monotonically and follows more steeply than $A^{-1/3}$, which is shown by solid line.

Fig 2 Energy differences per nucleon of each term for the cases with and without the pion mean field as a function of the mass number Here shown is the case of the jj-closed shell nuclei

We discuss here the behaviors of the contributions of various terms when the pion mean field becomes finite as a function of the nucleon mass. Figure 2 shows the energy differences per nucleon of each term for the cases with and without the pion mean field for the case of the jj-closed shell nuclei. The pion energy per nucleon is denoted by the dotted curve and the energy gain per nucleon decreases with the nucleon mass. The kinetic energy difference (long dashed curve) and also the potential energy difference (short dashed curve) are positive. The potential energy includes the contribution from the sigma and omega mesons. This indicates that the energy gain due to the pion mean field compensates largely with the kinetic energy and the potential energy. Both the kinetic energy and the potential energy act against the occurrence of finite pion mean field. The net result is a slight gain in energy in favor of surface pion condensation.

We show here the results of the finite pion mean field by showing the pion energy per nucleon for various nuclei as a function of the pion nucleon coupling constant (f_π) in Fig 3. If the pion nucleon coupling constant is decreased by 5 %, already the onset of the finite pion mean field disappears except ^{12}C . This fact indicates that the symmetry breaking mean field is fragile and the various effects as delta isobar contributions, the rho meson tensor coupling effects, the effects of the Fock (exchange) term and several others would influence the occurrence of the finite pion mean field in finite nuclei. This quantitative aspect will be

Fig 3 The pion energy per nucleon as a function of the pion-nucleon coupling constant square
In this systematic study, we have fixed the set of the filled single particle states with the same spins as in the closed shell configurations

worked out in the subsequent works²⁰⁾

§4 Intrinsic state and parity projection

We discuss here the role of the pion by analyzing a symmetry-broken intrinsic state and performing the parity projection from the symmetry broken intrinsic state. We write the single particle state with mixed parity (2.5) in a simpler form as

$$|\bar{j}\bar{m}\rangle = \alpha_j |jm\rangle + \beta_j |\widetilde{j}\bar{m}\rangle \quad (4.1)$$

Here, $|\bar{j}\bar{m}\rangle$ denotes a parity mixed single particle state expressed as a linear combination of $|jm\rangle$, some parity state (we call it as a normal parity state) and $|\widetilde{j}\bar{m}\rangle$, the opposite parity state (abnormal parity state). If we write the angular part of $|jm\rangle$ as $\mathcal{Y}_{jm}$ and the one of $|\widetilde{j}\bar{m}\rangle$ as $\mathcal{Y}_{\widetilde{j}\bar{m}}$, they are related as $\mathcal{Y}_{\widetilde{j}\bar{m}} = (\vec{\sigma} \cdot \hat{r}) \mathcal{Y}_{jm}$. Hence, the pionic correlation is the source of the admixture of the two states. We write the intrinsic state with these single particle states up to the Fermi surface and with all the magnetic sub-shells filled as

$$\begin{aligned} \Psi &= \prod_{jm} (\alpha_j |jm\rangle + \beta_j |\widetilde{j}\bar{m}\rangle) \\ &= \prod_{jm} \alpha_j |jm\rangle + \sum_{j_1 m_1} \prod_{jm \neq j_1 m_1} \alpha_j \beta_{j_1} |jm\rangle |\widetilde{j_1 \bar{m}_1}\rangle \\ &\quad + \sum_{j_1 m_1 j_2 m_2} \prod_{jm \neq j_1 m_1 j_2 m_2} \alpha_j \beta_{j_1} \beta_{j_2} |jm\rangle |\widetilde{j_1 \bar{m}_1}\rangle |\widetilde{j_2 \bar{m}_2}\rangle + \end{aligned} \quad (4.2)$$

This intrinsic state has the total spin 0 because all the magnetic sub-shells are filled, but the parity is mixed. The first term,

$$|[0]\rangle = \prod_{jm} \alpha_j |jm\rangle = N_0 \Phi_0 \quad (4.3)$$

in (4.2) has the positive parity and corresponds to the ground state without the pion. Here, $\Phi_0 = \prod_{jm} |jm\rangle$ is the zeroth order normalized ground state and $N_0 = \prod_{jm} \alpha_j$ is the normalization factor. The second term has the negative parity, since each normal parity state, $|j_1 m_1\rangle$, is replaced by an abnormal parity state, $|\widetilde{j_1 m_1}\rangle$. The second term can be written as a coherent sum of 1p-1h states over the occupied single particle states $(j_1 m_1)$,

$$|[1p-1h]\rangle = N_0 \sum_{j_1 m_1} \frac{\beta_{j_1}}{\alpha_{j_1}} |\widetilde{j_1 m_1}\rangle |(j_1 m_1)^{-1}\rangle \quad (4.4)$$

In the above expression $|(j_1 m_1)^{-1}\rangle$ represents one hole state, $\prod_{jm \neq j_1 m_1} |jm\rangle$. The third term in Eq. (4.2) can be seen in the same manner as

$$|[2p-2h]\rangle = N_0 \sum_{j_1 m_1 j_2 m_2} \frac{\beta_{j_1}}{\alpha_{j_1}} \frac{\beta_{j_2}}{\alpha_{j_2}} |\widetilde{j_1 m_1}\rangle |\widetilde{j_2 m_2}\rangle |(j_1 m_1)^{-1} (j_2 m_2)^{-1}\rangle \quad (4.5)$$

with

$$|(j_1 m_1)^{-1} (j_2 m_2)^{-1}\rangle = \prod_{jm \neq j_1 m_1 j_2 m_2} |jm\rangle \quad (4.6)$$

It is a coherent sum of 2p-2h states on the non-perturbative ground state Φ_0 . Hence, if we say the first term as the 0p-0h state, then the second term is a coherent 1p-1h state with 0^- spin parity. The third term consists of 2p-2h states with a pair of 1p-1h states with 0^- spin parity and therefore has 0^+ spin parity. The next term has three 1p-1h pairs with 0^- spin parity and therefore has 0^- spin parity and so on.

By projecting out the good parity states from the mixed-parity intrinsic wave function (4.2), we obtain the states of good-parity, positive and negative parity states,

$$P_+ \Psi = |[0]\rangle + |[2p-2h]\rangle + |[4p-4h]\rangle + \dots, \quad (4.7)$$

$$P_- \Psi = |[1p-1h]\rangle + |[3p-3h]\rangle + \dots \quad (4.8)$$

The operators, P_+ and P_- are the projection operators of positive parity and negative parity respectively. Hence, the positive parity state consists of even number of 1p-1h pairs with 0^- spin parity. This means that the positive parity projection provides 2p-2h states as the major correction terms. Hence, surface pion condensation together with parity projection provides the 2p-2h admixture due to the pion exchange interaction. This admixture corresponds to

the D-state admixture in the α particle. The D-state probability in α particle is known to be around 15%^{4), 5), 7)}. The negative parity state has odd number of 1p-1h pairs with 0^- spin parity. This is the brother state having the quantum number of 0^- to the 0^+ ground state. Therefore, the ground state consists of highly correlated particle-hole states and the 0^- state is a coherent 1p-1h state compiled on the highly correlated 0^+ ground state.

This demonstration gives us the hint of providing the two groups of the strong pion energy case, the jj-closed shell, and the weak pion energy case, the LS-closed shell. In the jj-closed shell case, this higher spin-orbit partner state is now filled and is used for the 0^- particle-hole excitations as a hole state. In this case, the particle state is $j_< = (l_{\max} + 1) - 1/2$ and the hole state is $j_> = l_{\max} + 1/2$. On the other hand, in the LS-closed shell case, the higher spin-orbit partner $j_> = l_{\max} + 1/2$ in the upper shell above the Fermi surface is not used for the 0^- particle-hole excitations. Hence, the pion energy is gained largely for the jj-closed shell nuclei, because the number of states to be mixed by the pionic correlations increases and these states contribute coherently to the pion energy. More attraction is expected for jj-closed shell nuclei than those of the LS-closed shell nuclei. We are, however, not yet clear about the peculiar mass dependence of the LS-closed shell nuclei. This distinct feature of the difference between the LS-closed and the jj-closed shell nuclei can be the consequence of using only the pion correlations.

We discuss here the examples of the qualitative consequence of finite pion mean field. First, we discuss the Gamow-Teller (GT) transitions. Without pion condensation, there exist no transitions for LS-closed shell nuclei as ^{16}O , ^{40}Ca and only two transitions, for example, for ^{90}Zr . However, the mixing of parity in the intrinsic state allows transitions of 2p-2h states (see Eq. (4.7)). This makes the spectrum of the GT transitions with some GT strengths above the two dominant peaks in ^{90}Zr ⁸⁾. Hence, naturally we have strengths in the region of the higher excitation energy in the mean field theory with the pion as the experiment demands. Second, the longitudinal spin response functions, which are caused by the pionic correlations, should be largely modified due to pion condensation. Since the large pionic strength is used up to construct the nuclear ground state, the pionic fluctuation oughts to be reduced largely. This should make the spin response in the pion channel weak. This remains to be demonstrated in the future work.

Until now the phenomena of involving high momentum components have been assigned as nuclear correlations induced by the repulsive core and the tensor force. The surface pion condensation provides us with the possibility to describe the part of the correlation effect induced by the tensor force. This fact indicates that we are now able to separate the short range correlation phenomena into the one due to the pion and the other as due to the repulsive core. In fact, as seen above in the discussion of the parity projection, the surface

pion condensation provides a large amount of the 2p-2h excitations in the nuclear ground state automatically. There should be many other consequences of surface pion condensation in nuclear phenomena. The pairing correlations and the spin-orbit couplings are all surface phenomena and surface pion condensation would couple with these correlations and provide rich phenomena.

§5 Conclusion

We have discussed the possible occurrence of the finite pion mean field in finite nuclei by introducing the pion field in the relativistic mean field (RMF) theory. We have extended the RMF theory by introducing the parity-mixed single particle basis to accommodate the finite pion mean field. We have taken the TM1 parameter set in the RMF theory and introduce the pion field in the pseudovector coupling with the nucleon. With the use of the pion-nucleon coupling constant in free space, we have made calculations for $N=Z$ closed shell nuclei and demonstrated the actual occurrence of finite pion mean field. We have demonstrated that the potential energy associated with the pion behaves as proportional to or even stronger than the nuclear surface. Hence, we name the on-set of the finite pion mean field as surface pion condensation. We have made qualitative discussions on the consequence of surface pion condensation on the Gamow-Teller strengths, the spin response functions and the short range correlations. The large difference of the pion energies in jj -closed shell and LS -closed shell nuclei has been found and may be connected with the mechanism of a part of spin-orbit splitting as due to the tensor force discussed long time ago by Takagi *et al.* and Terasawa²³⁾

We would like to stress here that we are at the initial stage of our investigation on the role of the pion in finite nuclei. We have to do various studies in order to establish the mean field theory with the inclusion of the pion, pursued in this paper for medium and heavy nuclei. We have to definitely introduce the rho meson tensor coupling term, which acts against the pion finite mean field. We should include also the delta isobar pion coupling terms, which favor finite pion mean field. We have to study carefully the effect of the short range repulsions, so called g' term in the non-relativistic framework. We would like also to work out the exchange terms, the Fock terms in the relativistic many-body theory. We then have to work out the parameter search for the coupling constants in the Lagrangian. We should further perform the parity projection. There are many studies to be done for the establishment of the occurrence of surface pion condensation. We are at the gate of exploring both theoretically and experimentally the phenomena caused by the most essential boson, the pion, in nuclear physics.

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