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Null Instanton Model on Binding and Excitation Energies of Light and Magnetic Moments of All Nuclei

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The binding and excitation energies of light nuclei are investigated as the eigenvalue problem of a null-instanton induced potential. A study of alpha particle model is done in this approach in order to find an associated physical property from a different aspect. A correction to the intrinsic part of Schmidt magnetic moment has been estimated for odd-A and odd-odd nuclei. It arises through the quarks belonging to the nucleon which is free from a pure photon as well as an associated intranucleon exchange by the similar eigenvalue problem arising from internucleon quark exchange.

A quenching of the magnetic moment can automatically be explained through the compact theoretical formula. It is discussed in relation to the experimental information suitably manipulated, while an anti-quenching requires a delicate examination case by case.

$.1 Introduction

A null-instaton induced potential is assumed to be valid for both parallely and anti-parallely aligned pair of quarks even if the instanon number of the each of associated quark is not zero. The instanton number of an unpaired quark, e.g., the extra one in the nucleon should be zero independently in order to keep the nullity of it in the hadronic states. The simultaneous realization of totally null instanton number of quark pair in addition to unpaired one gives a special role on the quark pairs in the nucleonic state.

We assume the existence of such an induced potential is an exchange type and it is admitted once the quarks and antiquarks are produced in nature,[1] viz., it naturally assumes the existence of the sign inverse potential which induces its partner with the same

quantum number for any state owing to the admittance of the sign inversion of the strength operator of the potential. This seems to be true for the hadronic world as exemplied by the existence of sign inverted states of many known baryonic as well as mesonic ones.

We assume that the same idea can be extended straightforwardly to the quark pair between nucleons by extending the same idea to internucleon quark exchange. Such an extension is allowed by the well known Okubo-Zweig-Iizuka (OZI) rule. Only the difference is an energy scale in between hadronic and nucleic world. In what follows the binding, excitation energies of nuclei and a quenching of Schmidt magnetic moment and possibly its anti-quenching phenomena as an eigenvalue problem based on one and the same potential.

Zinn-Justin[2] has discussed the null-instanton effect extensively and pointed out that it has a strong spatial distance dependence and diminishes quickly with an increase of it.

In this paper the binding[3] as well as the excitation energies[4] of nuclei are studied upto mass number A=4. This mass number region contains spatially extended structure of deuteron and a p-state excitation of a single nucleon so they provide a typical set of examples to start with.

The spin and parity of 4 states among observed10 excited states of the alpha particle are determined experimentally and the remaining 6 states (in which one has two possibilities) less definitely, while the isospins T's are assigned in referring to the charge independent states with A=4. For definiteness we assume T=0 for all unidentified states and the spin and parity 1^+ to the excitaion energy E_x=25.5 MeV in our analysis. The various devices to resolve the difficulties encountered in this restricted mass number choice will involve a typical feature to which we will encounter in future in order to proceed high mass number region. Apart from this problem we have also considered the binding energies of nuclei with A up to 80 by confining them solely as composite states of alpha particles for the purpose to point out a geometrical problem associated with the null instanton model, by eliminating the Coulomb energy. Here the binding of alpha state and their paste part from inter-alpha q-q exchange potential can be separated reasonably from the eigenvalues.

The quenching of magnetic moments from the Schmidt value[5] has a long history. Miyazawa[6] has pointed out that it is caused

from that of meson cloud of the bound nucleon in the nucleus. Arima and Horie[7] have reduced it to the configuration mixing of the adjacent states having the same parity and spin 1 unit different one. However, the observed anti-quenching phenomena[8] has been left untouched. Our model replaces Miyazawa's pion cloud by quark exchange and Arima-Horie's as a fractional nucleon number dependence resulting from multi-quark-pair exchange contribution in the outermost shells. The effect can occur only through the intrinsic part of the Schmidt magnetic moment. The Schmidt formulae for odd-A and odd-odd nuclei admit us to separate the orbital and intrinsic part in them so one can perform a detailed comparison of the experimental and the theoretical results in a general form for individual nucleus without making any approximation. It admits us to explain an occurrence of an gradual increase of quenching in a given orbit with the increase of the outermost nucleon numbers by keeping the same orbit for odd N nuclei where an anomaly of quenching series is explained by the corresponding change of the associated sub-shell structure. This detailed normal quenching series based on the theoretical formula can only be realized vividly in odd N nuclei, but not be realized in odd Z nuclei. It seems that the Coulomb effect obstruct the effect and will not admit such a beautiful picture for the latter.

Our model may be classified as the L-S coupling approach where the spin effect comes into play as an eigenvalue problem of the fermionic operator through the quark-field exchange, while the orbital part behaves as a bosonic parameter. For the magnetic moment of the nucleus the latter notion does not become important.

$.2 is devoted to the binding and excitation energies upto A=4. The alpha particle model is studied in $.3. We spend $.4 and 5 for magnetic moments of odd-A and odd-odd nuclei, respectively. In $.6 we give the conclusion of this work and a future plan of our study.

$.2 Binding and Excitation Energies in Light Nuclei

The s-wave part of the quark exchange potential in the strong interaction is defined as an operator form:

$$c\left[\frac{3+\sigma_i.\sigma_j}{4}+\frac{1-\sigma_i.\sigma_j}{4}\right]P_{ij}\,, \qquad (2.1)$$

where σ_i and σ_j are the Pauli spin operator for the i-th and j-th quark, respectively, and P_{ij} the exchange operator between them. The strength operator c is determined by the internucleon quark pairs as operands in the system. It is allowed to take both signs, since we introduced Eq.(2.1) only when quarks and anti-quarks are present in a formation of hadrons. The reason why one can assume this potential as an internucleon one is based, first of all, on the well-known Okubo-Zweig-Iizuka (OZI) rule, and secondly it automatically reproduces a hard core between nucleons, since only the same colored quark exchange is admitted between nucleons in order to keep the color singlet nature of an individual nucleon.

The main subject of this section is to explain the binding and excitation spectra of light nuclei in terms of this potential so as to show its utility through these examples. The total Hamiltonian is defined as a sum of the mass operator of total nucleons of the system with possible radial and orbital excitation in association and the potential operator defined above should be summed over all the composite internucleon quark pairs. The wave function of the nucleus is defined as the Slater determinant made from quark composite wave function as a member of the SU_2 doublet (if necessary as that of SU_3 octet) baryons[9] with an appropriate spin as well as a color assignment on composites so as to keep color singlet property of the nucleon. This formulation corresponds to L-S coupling scheme. Notice that the wave function of the nucleon takes one and the same form in two flavor groups of interest. Only the relative weight of the quark composite states of nucleons becomes important so one could choose the most convenient form of wave functions in a practical study in a consistent manner to SU_4 (a semi-direct product on flavor and spin group) in a sense that such a choice does not break the internal symmetry of quark flavor combination in the estimate of the matrix elements of interest. One may find that the effective choice consists of two uud (ddu) and udu (dud) combination for p(n) with cyclic spin assignment in mind and an appropriate weight factors for p(n) component wave function arising from flavor group which reduces the labor in the evaluation of matrix elements, since there does not occur the overlapping of the components of wave functions including the color specification. In the problem under study the relative phase π between proton

and neutron is not important. It becomes important only when one treats the transient matrix element, such as, a beta decay.

In the study of the binding and excitation energy there appear totally 9 ${}_AC_2$ expectation values of operator Eq.(2.1), where 9 is the possible number of quark pair states between two nucleons and ${}_AC_2$ the binomial coefficient. This number corresponds to the interaction between quarks belonging the nucleus as internucleon pairs. Why this maximum possibility is allowed in spite of the color whiteness requirement of the bound nucleon stated already? Notice that the color exchange between quarks in the nucleon (intranucleon exchange) occurs more frequently than that between nucleons (internucleon exchange) so that no restriction arises to
keep the color whiteness in the total number of pairs. This is certified by comparing the (absolute) magnitudes of the corresponding strengths found in the intra- and inter-nucleon exchange in our previous[1] and present study, respectively. It becomes a few times an order of magnitude less in the light nuclei under study as will be shown in this section. The higher the mass number becomes, the more natural this statement does, due to the expected smallness of expectation values corresponding to the gradual decrease of the observed level spacing with the increase of mass number. In what follows we can explain the binding energies as well as the assumed level structures of light nuclei, together with the prediction of their sign inverted partners, by taking size variation into account. The property[2] of the null-instanton induced potential provides us a guiding role in the course of analysis performed in this paper.

Let us begin with the binding energy formulae for D, T and α particles which are derived as the expectation values of the potential Eq.(2.1) by neglecting the size effect. They become:

$$B(D) = 4\,a0 + 5\,a1, \qquad (2.2)$$

$$B(T) = 14\,a0 + 13\,a1, \qquad (2.3)$$

and

$$B(\alpha) = 2(14\,a0 + 13\,a1). \qquad (2.4)$$

Here a0 and a1 are the s-state expectation values for anti-parallel and parallel quark pair state, respectively. Notice that we do not require any correction, in addition, such as, a volume term. In order to determine the parameters a0 and a1 we solve the coupled equation by taking the size effect expected from the null-instanton induced potential into account. For this purpose we choose a modified inputs B(T) x relative to B(D) on the left-hand-side (l.h.s.) of Eq.(2.3), where x is an adjustable parameter. This induces the modification of parameters in Eq.(2.3) a0 → c0 and a1→ c1, which will be reflected to the binding parameters of α particle by taking the factor 2 difference in the original equations Eqs.(2.3) and (2.4) into account. They are given explicitly by:

$$b0 = x\,c0\,\frac{B(\alpha)}{2B(D)} \quad \text{and} \quad b1 = x\,c1\,\frac{B(\alpha)}{2B(D)}\,. \qquad (2.5)$$

The s-state binding energies of the ground, first and second excited states of the α particle with isospin T=0, spin-parity 0^+, 1^+ and 2^+ state with the excitation energies E_x = 0, 25.5 and 33.0 MeV, respectively, are given in terms of the new parameters by

$$B(\alpha)_0 = 2(14\,b0 + 13\,b1), \qquad (2.6)$$

$$B(\alpha)_1 = 27(b0 + b1), \qquad (2.7)$$

and

$$B(\alpha)_2 = 6(4\,b0 + 5\,b1), \qquad (2.9)$$

respectively.

Under these constraints we solved coupled equations Eqs.(2.2)-(2.4) by varying x so as to reproduce the B(D), B(T) and $B(\alpha)_0$ as well as its T=0, $J^P = 2^+$ excitation energy at E_x =33.0 MeV. We find the following sets of solutions which reproduce the desired result:

a dimension less parameter: x = -0.2095,

and binding energy parameters in units of MeV:

$$a0 = c0 = -12.8528, \quad a1 = c1 = 10.7272, \quad b0 = 4.49145,$$

$$b1 = -3.74863. \qquad (2.10)$$

These parameters predict the assumed fourth excited T=0, $J^P = 1^+$ at E_x = 25.5 MeV at 8.24 MeV which is far below the observed value. From here after we shall omit energy units MeV unless stated otherwise.

Before proceed further we should like to comment shortly on the D, T and ^3He excitation energies based on our parameters studied above. Our model predicts: (1) the first excitation energy of D at T=0, $E_{xinv}(1^+)$ = 2.22452 and T= 1, 0^+, E_x=19.55 without making any manipulation (see, later discussion), and (2) the T=1/2, $J^P = 1/2^+$ first and second excitation energies of T and ^3He are predicted at 8.482, 16.96 and 7.72, 15.43, respectively.

For the sake of clarity it is convenient to tabulate our plans to explain the α-particle energy spectra in Table I, including sign inverted partners.

Table I. Energy spectra of α particle

				s. i. p.				
	state				state			comment
ex.	T	J^P	E_x	ex.	T	J^P	E_x	
g. s.	0	0^+	0.	v.	0	(0^+)	(56.6)	si 0^+ 20.1, 30.5
1st	0	0^+	20.1	v.	0	0^+	36.5	basip 0^+ 56.6
2nd	0	0^-	21.1	v.	0	0^-	35.5	basip in T=0
3rd	0	2^-	22.1	v.	0	2^-	34.5	basip in T=0
4th	0	1^+	25.5	v.	0	1^+	31.1	basip in T=0
5th	1	2^-	26.4	v.	1	2^-	30.2	basip in T=1
6th	1	1^-	27.5	8th	1	1^-	30.5	basip in T=1

7th	1	0^-	29.5	v.	1	0^-	30.7	basip in T=1
9th	0	1^-	31.0	v.	0	1^-	33.7	basip in T=0
10th	0	2^+	33.0	v.	0	2^+	37.6	basip in T=0

In Table I we have used abbreviated symbols in order to make short the tabulation. They are s. i. p. = sign inverted partner, ex. = excited state, g. s. = ground state, v = virtual state, si = split into, basfv = both are splitted from virtual and basip = both are sign inverted partners. The first, second, third and fourth column represent numbering of excited states, the corresponding isospin T, spin parity J^P and the excitation energy E_x, respectively. The similar applies to the fifth, sixth, seventh and eighth column. The last column gives a short comment on partner states in the same line.

The sign inversion point is set at B = 0 when $\Delta B > 0$ case when no candidate of the partner cannot be found experimentally within positive binding energy, while the partner's E_{xinv} is set at $E_x + 2\Delta B$ when $\Delta B < 0$ by definition. The last choice should be tested from independent data and leave it for future. In this tabulation the 6th and 8th excited states are observed mutual partners arising from the sign change of potential strength.

The explanation of the excitation energies of the first excited states and its partner goes as follows. Let us specify their wave functions by the number corresponding to their excitation energy, viz.,

$$\psi(20.1) = \cos\theta \ \psi_g(0^+) + \sin\theta \ \psi_{ginv}(0^+)$$

and (2.11)

$$\psi(36.5) = \sin\theta \ \psi_g(0^+) - \cos\theta \ \psi_{ginv}(0^+).$$

Here θ is the mixing angle, $\psi_g(0^+)$ and $\psi_{ginv}(0^+)$ are the wave functions for the ground and its partner state, respectively. The corresponding excitation energies are given by

$$E(20.1) = \cos^2\theta\ E_g(0^+) + \sin^2\theta\ E_{ginv}(0^+)$$

and (2.12)

$$E(36.5) = \sin^2\theta\ E_g(0^+) + \cos^2\theta\ E_{ginv}(0^+).$$

The mixing angle becomes $\theta = 36.5814°$. In this mixing the ground state remains unchanged, while its partner's energy is projected into two states: one is at the first excited and the other at its partner's energy.

The similar applies to the fourth excited state at 25.5 and its partner's one at 31.1 as the mixing of predicted 8.2399 and its partner at 48.35. One finds the mixing angle $\theta = 40.993°$.

The energy spectra of odd parity states appearing as 5th, 6th, 7th and 9th excited states are suitable to study in comparison with those for ^4H and ^4Li nuclei. Let us begin with the theoretical expression for one neutron excitation in $p_{3/2}$ and $p_{1/2}$ excited energy belonging to T=1. From $p_{3/2}$ n one gets $2^-, 1^-$ as well as its sign inverted partner combining it with the up down spin in $s_{1/2}$ n while $p_{1/2}$ gives 1^- and 0^- partners. The former is absorbed in Eqs.(2.14a) and (2.14b) by a relative weight 2:1 by the Clebsch-Gordan coupling rule. A sign inverted partner in 0^- state will be observed as an energetically degenerated state, as will be shown numerically below, which is based on the inspection of the energy spectra of the charge independent nuclear spectra. All these belong to T=1 state. One finds their energy eigenvalues as:

$$E_x(2^-) = (14\ b0 + 13\ b1) + 3\ r(p) - (13\ b0(p) + 14\ b1(p)) = 26.4, \quad (2.13a)$$

$$E_x(1^-) = (\text{common as above}) - (14\ b0(p) + 13\ b1(p)) = 27.5, \quad (2.14a)$$

$$E_{xinv}(1^-) = (\text{common as above}) + (14\ b0(p) + 13\ b1(p)) = 30.5, \quad (2.14b)$$

and

$$E_x(0^-) = (\text{common as above}) - (14\ b0'(p) + 13\ b1'(p)) = 29.5, \quad (2.15a)$$

$$E_{xinv}(0^-) = (\text{common as above}) + (14\ b0'(p) + 13\ b1'(p)) = 29.5. \quad (2.15b)$$

Here r(p) represents the orbital excitation energy of n relative to 3 s-state nucleons, the eigenvalues for two quark pairs in different spin states are specified with and without prime, and energy eigenvalues are shown for clarity. We have admitted the sign inversion of eigenvalues only for the p-state n relevant terms, because the common term requires relatively big excitation energy and the r(p) in it is the bosonic parameter.

The excitation energy for the states belonging to T=0, $p_{1/2}$ n become:

$$E_x(0^-) = (\text{common as above}) - (14\ b0'(p) + 13\ b1'(p)) = 21.1, \quad (2.16a)$$

$$E_{xinv}(0^-) = (\text{common as above}) + (14\ b0')p) + 13\ b1'(p)) = 35.5, \quad (2.16b)$$

$$E_x(1^-) = (\text{common as above}) - (13\ bo'(p) + 14\ b1'(p)) = 31.0. \quad (2.17)$$

The numerical work reduces to solve a simple coupled equation by subtracting suitably the neighboring energy each other. The same expressions apply to the corresponding T= 1 ^4H and ^4Li states. We perform the numerical study for these two nuclei beforehand for convenience. One finds:

$$b0(p) = -0.774 \quad \text{and} \quad b1(p) = 0.926 \quad (2.18)$$

from ^4H spectra and

$$b0(p) = -0.607 \quad \text{and} \quad b1(p) = 0.793 \quad (2.19)$$

from ^4Li spectra. This nucleus admits to estimate the combined parameters 14 b0'+ 13 b1' from

$$14\ c0 + 13\ c1 = 1.8 - 14\ b0' - 13\ b1' \qquad (2.20)$$

Here the number 1.8 is the energy difference between 3rd and 2nd excited states in ^4Li. One finds

$$14\ b0' + 13\ b1' = -0.000011. \qquad (2.20a)$$

This result proves numerically that $E_x(0^-)$ and $E_{xinv}(0^-)$ are degenerated energetically. We think that the main cause of the approximate energy-specification sign found experimentally for the 2nd excited 0^- state of ^4Li is counted as a reflection of the degeneracy of the two states predicted here. The mixing of T=1 0^- partners for excitation energies 3.2 and 3.1 for ^4Li and ^4He, respectively, have been done corresponding to the sets of theoretical excitation energies 2.1 and 4.3 as well as 1.65 and 3.55. The former set gives the mixing angle 45°, while the latter set gives the same mixing angle but predicts $E_x(0^-) = 2.6$ which is somewhat smaller than the observed 4.1. This, however, may not negate completely the possibility of the degeneracy of interest in that state in ^4He.

Finally, the r(p) is estimated from the T=1, 2^- ground state energy in ^4He, viz.,

$$14\ b0 + 13\ b1 + 3\ r(p) - 13\ b0(p) - 14\ b1(p) = 26.4$$

$$r(p) = 0.3346 \qquad (2.11)$$

$. 3 α-particle Model in Null-Instanton Approach

The purpose of this study is to look for the role of geometrical problem inherent in the null-instanton model from completely different objects. The interaction potential is the same operator given in Eq.(2.1). The mass range is confined from A = 4 to 80.

It is useful to classify the eigenvalues into α-particle formation and their paste parameters. They are specified by 2 n (14 b0 + 13 b1) and 36n(n-1)(p0 + p1), where n is the number of α-particles bound in s-state. We can eliminate the Coulomb effect by making use of the binding energy data for charge

independent triplet nuclei and its Z(Z-1) dependence. This is possible for nuclei of interest except for α where we adopted A=5 nucleus data for the same purpose. The result becomes: $B(\alpha)$= 28.7996, $B(^8Be)$= 59.5333, $B(^{12}C)$= 99.7414, $B(^{16}O)$= 140.283, $B(^{20}Ne)$= 183.9714, $B(^{24}Mg)$= 227.1721, $B(^{28}Si)$= 272.8139, $B(^{32}S)$= 322.3035, $B(^{36}Ar)$= 369.4479 and $B(^{40}Ca)$= 411.2566 (MeV). One can estimate the paste energy part of each nucleus by taking the difference

$$\text{paste} = \frac{1}{n(n-1)}\,[B(^{4n}_{2n}X) - n\,B(\alpha)] = 36(p0 + p1), \quad (n = 2 - 10). \qquad (3.1)$$

Here the use is made of a shortened notation X for a nucleus.

One may find:

paste = 1.9341, 4.44753, 4.18078, 3.99734, 3.62497, 3.39127, 3.28238, 3.06254, 2.73912, for n=2-10, respectively. (3.2)

We have treated that all alpha particles are bound in s-state in this numerical work.

The smallness of n=2 paste value is expected from the beginning, since ^8Be is unstable for 2α decay. The gradual degrease of bond strength of paste may be reduced to three reasons:
(1) a purely geometrical reason , (2) the shell structure effect, and (3) the alpha particle formation in the nucleus is more restricted with the increase of mass number.

The third one is difficult to answer for the time being and left for future. The second one is an orbital excitation problem by assuming that the number of alphas are raised to their orbit specified by the assignment of shell structure. This is done by adding the orbit dependent parameters to the binding energy formula. The additional term becomes:

-r(p), -2 r(p), -3 r(p), -3r(p)-r(d)-3 r(p,d), -3r(p)-2r(d)-6r(p,d),

-3r(p)-3r(d)-r(2s)-9r(p,d)-3r(p,2s)-3r(d,2s), -3r(p)-4r(d)-r(2s)-12r(p,d)-3r(p,s)-4r(d,s), -3r(p)-5r(d0=r(2s),-15r(p,d)-3r(p,2s)-5r(d,2s), for n= 2-10, respectively. (3.3)

Here the single and relative orbit dependence are specified by one and two indices, and finds but it cannot help to remove the discrepancy found in the paste term shown above.

We think that the first one is an attractive problem to pursuit from the null-instanton model point of view, because the alpha particles can make a close contact up to n=4, while it is difficult to make it above n=4. We cannot yet solve this geometrical problem in the writing stage of the paper. It is preferable to get a help of a mathematician in order to clarify the spatial dependence of the strength parameter.

$. 4 Magnetic Moments of Odd A Nuclei

A normal definition of u and d quark magnetic moments (abbreviated m.m. hereafter) are defined as

$$\mu_u = \frac{2}{3}\mu \quad \text{and} \quad \mu_d = -\frac{1}{3}\mu, \qquad (4.1)$$

by neglecting their mass difference. Here, μ is one-half of the gyromagnetic ratio for the quark. It is well known fact that these choices explain the observed ratio of the proton and neutron m.m. in certain degree in SU_4 limit, where this group is introduced as a semidirect product of flavor SU_2 and spin SU_2 group. The p and n are classified as members of flavor SU_2 doublet. We shall at first show that the p and n m.m. in SU_4 limit are modified under the presence of null-instanton induced potential Eq.(2.1), by taking its effect to the p and n under interaction with a photon through their composite quarks. The modification of the m.m. in pure SU_4 limit occurs when this photon associated quark has additional interaction with that belonging to the same nucleon. We call this the intranucleon q-q exchange effect. The estimate of the effect can be done as the eigenvalue problem as we have already experienced through previous two sections. Only the difference is that the photon associated quark can have two-bonds with remaining two quarks free from the photon. The eigenvalues should be estimated under a cyclic permutation of the photon for essentially independent two components of wave

functions. In the final stage one should do in addition to the multiplication by the square of relative weight of the component wave function and the summation over the number of same matrix element arising from the number of SU_4 components with similar spin structure which were skipped at first stage for ease of estimate. After doing this procedure for each of initially classified essentially two independent-component wave functions, the net result should be normalized by the corresponding factor of the total wave function. We recommend novice to refer a precious work of Gibson and Pollard on this procedure[9] in the pure SU_6 approach. Repeating one and the same procedure the correction term enters to the magnetic moments of proton and neutron as additional terms in the pure SU_4 approach. The result is given by:

$$\mu_p = (1 - a1 - a2)\,\mu \quad \text{and} \quad \mu_n = -\frac{2}{3}(1 - \frac{3}{2}a0 - \frac{1}{2}a1)\,\mu, \qquad (4.2)$$

for p and n, respectively. They are considered as the observed magnetic moments of free p and n, respectively, including the strong interaction effect explicitly. We shall use the same notations as those used in $.1 for the correction term. Here they are measured in units of μ. We hope that this convention will not induce confusion. The eigenvalues a0 and a1 represent those for anti-parallel and parallel spin q-q pair states in the nucleon, respectively. The physical meaning of the parametrization in our previous work[10] has already been discussed[1] in connection with the present result.

When the nucleon is bound to the nucleus its m.m. is expressed by the so-called Schmidt value.[5] The observed m.m. of many odd-A nuclei fall in between two Schmidt lines which is known as the quenching phenomenon. This phenomenon was taken seriously by Miyazawa as well as Arima and Horie in early days. The former author proposed that it is caused from the quenching of the pion cloud around the m.m. relevant nucleon in the nucleus. The latter authors have shown that the configuration mixing of neighboring nucleons with assumed quantum number explains the phenomena in right ways semi-quantitatively by a perturbative approach. Some of the m.m. of odd-A nuclei show anti-

qeunching, viz., their measured m.m. fall outside of the Schmit lines. This problem is left untouched for many years.

The Schmidt formula admits us to separate the orbital and intrinsic m.m. . In that formula one substitues the observed free m.m., viz., it is totally called a Schmit value. We find that the null-instanton induced potential modifies solely affect the intrinsic part of the Schmidt formula. Let us call it the intrinsic m.m. after eliminating orbital part in order to discriminate it from the free one. The modification of the intrinsic m.m. occurs through the operation of the null-instanton induced potential between quarks belonging to different nucleons. Let us name this the internucleon exchange. The calculation of the eigenvalues of the operator is classified into two types, owing to the kind of quarks joining to the exchange process. They should be members of the intrinsic m.m. relevant nucleon, viz., these two types are: (i) quarks free from the photon vertex and (ii) that free from the intranucleon exchange. The number of the former is two, while the latter is one. In the practical work there appear a few set of eigenvalues for each sub-wave function. It should be repeated for the cyclic permutation of a photon and intranucleon exchange vertices in vertex free manner. Usually the positive strength sign of the potential causes a quanching of m.m. under the operation of the Pauli exclusion principle. One can understand qualitatively that the machinery under study will give us a correct sign of quenching by noticing that the sign of the m.m. of p and n is realized in quark model, owing to the fact that total sign of them is due to the dominance of charge of u and d involved in them. If the eigenvalues of the internucleon potential strength take appropriate sign or the sign inverted partner's state contributes dominantly they cause the anti-quenching of the m.m. This depends on the structure of the nuclear state under study. One may find that the anti-quenching of the m.m. of odd-odd nuclei may also occur by definition.

It is not difficult to calculate the m.m. of odd A nuclei under the presence of the null-instanton induced potential, because all the nucleons are coupled to spin zero except for the one containing the photon vertex. Let us call the last one the host nucleon for convenience sake.

In the course of study we find that the quenching behavior may be explained by taking the contribution from nucleons belonging to the outermost shells into account. We shall present first our result on this limited choice in order to compare

it with the experimental data. The result is easily generalized to include the contribution from full nucleons. The constraint on the choice of number of shells will be decided by the property of the null-instanton induced potential. It may also depends upon the value of m.m. under study.

The static problem under study will be extended to a dynamical problem where the eigenvalue problem under study contributes part of it. We shall postpone a discussion on such an example at the end section of this paper.

Let us present the m.m. formulae for ^3H and ^3He and generalize them for odd Z and N nuclei in a form applicable to any odd A nuclei. One finds that intrinsic parts of them are given by

$$\mu\,(^3H) = [1 - a0 - a1 - 6(c0 + c1) + 3(a0+a1)(c0 + c1)]\,\mu \tag{4.3}$$

and

$$\mu\,(^3He) = -\frac{2}{3}[1 - \frac{3}{2}\,a0 - \frac{1}{2}\,a1 - 6(c0+c1) + \frac{3}{2}(3a0+a1)(c0+c1)]\,\mu\,, \tag{4.4}$$

respectively. Comparing these equations with those in Eq.(4.2) one sees that the eigenvalues c0 and c1 are those came from internucleon q-q pair spin states. The application of twice the Pauli exclusion principle has resulted sign of the last term in the square bracket. Usually a0+a1 and 3a0+a1 is less than 2 and 4, respectively, so that the quenching and anti-quenching are controlled by the sign of (c0+c1) term, a pure photon correction term. These results can easily be generalized in the convenient forms compared to the experimental m.m. for odd Z and N nuclei, viz.,

$$(\mu_p)_{\mathrm{intr}} = [\,1 - \frac{3(k+n)(c0+c1)(2-a0-a1)}{2\mu_p}]\,\mu_p\,, \tag{4.5}$$

and

$$(\mu_n)_{\rm intr} = [\,1 + \frac{3(k+n)(c0+c1)(4-3a0-a1)}{4\mu_n}\,]\,\mu_n\,, \qquad (4.6)$$

where the subscript intr means the intrinsic part of the m.m. in the Schmidt formula. Here k and n are number of protons and neutrons belonging to the outermost shell except for the host nucleon, viz., k and n take even numbers 0, 2, 4,... . In these expressions we did not discriminate shell dependence of eigenvalues c0 and c1. They may take different numerical values by the difference of shells to which host p and n belong., viz., it is more suitable to write:

$$(k+n)(c0+c1) \rightarrow (k_i + n_j)(co(i)+c1(j)), \qquad (4.7)$$

where i and j specify the relevant proton and neutron shell, respectively. The summation over all i and j admits us to include the full contribution from nucleons belonging to the given nucleus.

The formula given above allow us to make a detailed comparison with the experimental data by eliminating the orbital part in Schmidt formula and by taking out the intrinsic m.m. with a suitable manner. In the following study the shell structure of the outer most nucleons becomes important, so we shall refer the known shell structure[5] and the corresponding filled-shell nucleon numbers up to that shell for convenience. They are:

$$s_{1/2}(2), p_{3/2}(6), p_{1/2}(8), d_{5/2}(14), s_{1/2}(16), d_{3/2}(20), f_{7/2}(28), p_{3/2}(32),$$
$$f_{5/2}(38), p_{1/2}(40), g_{9/2}(50), g_{7/2}(58), d_{5/2}(64), d_{3/2}(68), s_{1/2}(70),$$
$$h_{11/2}(82), h_{9/2}(92), f_{7/2}(100), f_{5/2}(106), p_{3/2}(110), p_{1/2}(112), i_{13/2}(126),$$
$$\ldots . \qquad (4.8)$$

We shall arrange the intrinsic m.m. extracted from the experimental data will be compared easily with the second terms in the squared bracket in Eqs.(4.5) and (4.6), respectively, for odd-Z and N nuclei. In the actual study one may find only a fraction of (k+n) has a meaning. Even such a proportionality should apply to odd Z nucleus, while this does not seem to be realized apparently at high Z region as one sees from the data scan made in this work. The physical meaning will be discussed after studying them. The numerical results are presented in units of

μ_p and μ_n for odd-Z and N, respectively, where we have multiplied inverse of orbit dependent factor when it comes in from Schmidt formula so as to give the desired $(\mu_{p(n)})_{\rm intr} / \mu_{p(n)}$.

First of all, we shall tabulate in Table II the m.m. of charge symmetric nuclei for make it easy to compare.

Table II. Ratio of intrinsic m.m. and its free one

nucleus	$(\mu_p)_{\rm intr} / \mu_p$	$(\mu_n)_{\rm intr} / \mu_n$	nucleus
^3H**	1.0660	1.11213	^3He
^{11}B	0.604630	0.5368	^{11}C
^{15}N	1.020	(1.127)	^{15}O**
^{17}F	0.97470	0.989894	^{17}O
^{19}F	0.9411016	0.9859	^{19}Ne**
^{21}Na**	-0.349786	-0.57653	^{21}Ne
^{29}P**	(0.44217)	0.29024	^{29}Si
^{35}Cl	0.58374	0.543	^{35}Ar**
^{41}Sc**	(0.870)	0.833562	^{41}Ca**

In the first and the last column nucleus symbol for odd-Z and its charge symmetric one is listed. The corresponding ratios for intrinsic and free nucleon m.m. are given in the third and fourth column, respectively, for ease of comparison. The significant numbers are given based on their experimental accuracy. The nucleus symbol with double star superscript means that it is an artificial radioactive one. The round bracketed number means that its sign has not yet been determined experimentally, where we have chosen the sign so as to coincide with Schmidt value and to become closer to its charge symmetric partner's one.

One can learn physics involved in each datum. The absolute value less than 1 indicates that its m.m. lies within Schmidt lines.

When it exceeds 1 means that it is an indication of anti-quenching. We learn in general that the quenching and anti-quenching occur commonly to the charge symmetric nuclei where the lowest $s_{1/2}$ and $p_{1/2}$ orbits appear as common for anti-quenching ones. Turning our attention to individual pair.

Let us start from A=3 pair where k=0, n=2 and k=2, n=0 for odd Z and odd N, respectively. We have seen the rapid vaiation of the potential parameters occurs including their sign in A=2 to 4. Actually we found a0=-12.85, a1 = 10.73 and b0= 4.491 and b1= -3.749 (MeV) from A= 2 to 3 and 4 in $.1. The parameters c0 and c1 in this section correspond to those numbers with the same 0 and 1 indices. If the propotionality holds between energy and m.m. parameters the sign of c0+c1 becomes negative (positive) for the first (second) set, respectively. It is a delicate problem which state of sign inverted partners contibutes mainly the state under investigation which also concerns the sign of (c0+c1). Let us assume that our formula satisfies the anti-quenching condition observed here both for $s_{1/2}$ and $p_{1/2}$ orbit and turn our attention to next pair. The closeness of the ratios 1 in A=17 pair where k=n=0 means that inner shell effect is negligible. The core nucleons form the fully closed shell so they ditribute spherically symmetric manner on the core surface, where the only a very limited fraction of nucleons enters into the potential range so as to reduce the intrinsic m.m. In this connection a slight reduction of the ratios found for A=19 pair relative to A=17 admits us to estimate an additional phenomenological-reduction factor for each one which is multiplied to k=0, n=2 and k=2, n=0 choice, respectively. One sees that k and n does not appear as the reduction factor themselves. They should be taken as the convenient indices for pointing a quenching direction(see, later discussion). We shall leave the determination of individual parameter with a new multiplicative reduction factor for each nucleus and turn our attention to the data survey of many other nuclei from the similar point of view.

Let us consider next the m.m. of odd Z nuclei in general by classifying them by their host proton orbital states in arranging them with the increasing order of mass number and shell structuture. It is convenient to repeat this by taking ratio of intrinsic m.m. and free proton one corresponding to Eq.(4.5) as already done. In order to make clear the physical meaning of the correction term we shall list $1-(\mu_p)_{intr}/\mu_p$. One may expect a

gradual increase with (k+n) dependence with some positive factor for the particular subshell to which the host proton belongs. More precisely (i) the nearby closed shell, closed sub-shell, the change of the associated sub-shell structure should appear if the full outermost sub-shell nucleon contribute effectively. This expectation should be found from the suitable manipulation suggested above. Let us look at the experimental information for odd Z nuclei listed in Table III.

Table III. $1-(\mu_p)_{intr}/\mu_p$ for odd-Z nuclei

$p_{3/2}$	$^{7}_{3}Li$	$^{11}_{5}B$	$p_{1/2}$	$^{13}_{7}N$**	$^{15}_{7}N$
	0.192118	0.39537		-0.07188	-0.0202
$d_{5/2}$	$^{17}_{9}F$**	$^{27}_{13}Al$	$d_{3/2}$	$^{21}_{11}Na$**	$^{23}_{11}Na$
	0.02530	(0.41228)		1.34979	1.24916
$d_{3/2}$	$^{35}_{17}Cl$	$^{37}_{17}Cl$	$^{39}_{19}K$	$^{41}_{19}K$	$^{43}_{19}K$
	0.41627	0.33407	0.15941	0.05413	0.0221
$f_{7/2}$	$^{41}_{21}Sc$**	$^{43}_{21}Sc$**	$^{45}_{21}Sc$	$^{47}_{21}Sc$	$^{51}_{23}V$
	(0.130)	0.420	0.371132	0.162	0.2309
$f_{7/2}$	$^{53}_{25}Mn$	$^{57}_{27}Co$	$^{59}_{27}Co$	$f_{5/2}$ $^{51}_{25}Mn$**	$^{55}_{25}Mn$
	(0.267)	0.3834	0.4239	(1.3639)	1.30632
$p_{3/2}$	$^{61}_{29}Cu$**	$^{63}_{29}Cu$	$^{65}_{29}Cu$	$^{67}_{31}Ga$**	$^{69}_{31}Ga$
	0.592	0.56095	0.50409	0.70	0.636167
$p_{3/2}$	$^{71}_{31}Ga$	$g_{9/2}$ $^{93}_{41}Nb$	$^{99}_{43}Tc$**	$^{109}_{49}In$**	$^{113}_{49}In$
	0.440804	0.22370	0.39800	0.449	0.45436
$p_{3/2}$	$^{75}_{33}As$	$^{79}_{35}Br$	$^{81}_{35}Br$	$^{81}_{37}Rb$**	
	0.8234	0.60406	0.54540	0.624	
$f_{5/2}$	$^{83}_{37}Rb$**	$^{85}_{37}Rb$	$p_{1/2}$ $^{105}_{47}Ag$**	$^{107}_{47}Ag$	$^{109}_{47}Ag$
	0.280	0.245947	(0.1752)	0.16190	0.14365
$p_{1/2}$	$^{111}_{47}Ag$**	$^{113}_{47}Ag$**	$g_{9/2}$ $^{109}_{49}In$**	$^{111}_{49}In$**	$^{113}_{49}In$

	0.127	0.113		0.449	0.452	0.45436				
$g_{9/2}$	$^{115}_{49}In$**		$p_{1/2}$	$^{113m}_{49}In$**	$^{115m}_{49}In$**	$^{117m}_{49}In$				
	0.45013			0.0578	0.02202	(0.0143)				
$d_{5/2}$	$^{121}_{51}Sb$	$^{125}_{53}I$**		$^{127}_{53}I$	$^{129m}_{53}I$	$^{151m}_{63}Eu$				
	0.51365	(0.64)		0.7102	0.70	0.79				
$g_{7/2}$	$^{123}_{51}Sb$	$^{129}_{53}I$		$^{131}_{53}I$	$^{133}_{53}I$**	$s_{1/2}$	$^{127}_{55}Cs$**			
	0.38220	0.3684		0.4702	(0.5158)		0.448			
$s_{1/2}$	$^{129}_{55}Cs$**	$g_{7/2}$	$^{133}_{55}Cs$	$^{135}_{55}Cs$**	$^{137}_{55}Cs$**	$^{139}_{57}La$				
	0.470		0.3970	0.46613	0.51640	0.48873				
$g_{7/2}$	$^{175}_{71}Lu$	$^{177}_{71}Lu$	$^{181}_{73}Ta$	$f_{7/2}$	$^{165}_{67}Ho$	$p_{1/2}$	$^{169}_{69}Tm$			
	0.236	0.239	0.292		0.638		0.03597			
$d_{5/2}$	$^{185}_{75}Re$	$^{185}_{75}Re$	$d_{3/2}$	$^{191}_{77}Ir$	$^{193}_{77}Ir$	$^{191}_{79}Au$**				
	0.58061	0.56897		0.012889	0.020635	0.00758				
$d_{3/2}$	$^{193}_{79}Au$**	$^{195}_{79}Au$**	$^{197}_{79}Au$	$^{199}_{79}Au$**	$s_{1/2}$	$^{199}_{81}Tl$**				
	(0.0088)	(0.0136)	0.01288	0.08683		0.434				
$s_{1/2}$	$^{201}_{81}Tl$**	$^{203}_{81}Tl$	$^{205}_{81}Tl$	$h_{9/2}$	$^{203}_{83}Bi$**	$^{205}_{83}Bi$**	$^{209}_{83}Bi$			
	0.431	0.423058	0.417383		0.860	1.26	0.6367			
$p_{3/2}$	$^{227}_{89}Ac$	$^{231}_{91}Pa$*	$^{233}_{91}Pa$**	$d_{5/2}$	$^{237}_{93}Np$**					
	0.96	0.65	0.14		-0.43					
$f_{5/2}$	$^{241}_{95}Am$	$^{243}_{95}Am$**								
	0.36	0.36								

In this table m after A on left superscript attached to the nucleus symbol means a metastable, and single * and double ** on right superscript show the natural and artificial radioactive nucleus, respectively.

The data are arranged from light to heavy nuclei by classifying them in their orbital dependence with an increasing

order of shell structure. There does not seem to exists strong (k+n) dependence in general except for the first $p_{3/2}$ nuclei. The expected (k+n) dependence seems to disappear all other orbital states. This is understood as follows: first of all, (i) the effective nucleon number contributing to the correction not necessarily be integer as we saw already, because the null-instanton effect is very sensitive to the relative distance between the correction-term relevant nucleons, and secondly, (ii) the Coulomb effect obstruct the expected number dependence. This second effect forces to distribute protons on the surface of the core equally. As we shall see in comparison with the intrinsic m.m. of odd N nuclei the null-instanton effect and the Coulomb effect compete each other on the surface host proton binding as to arrange the outermost proton distribution on the core surface. Here the latter effect is somewhat over the former one. One may find a delicate dependence on the host-proton associated sub-shell change by examining the table. The new anti-quenching of m.m. will be found but it is restricted to a limited number of nuclei.

The similar quantity is extracted from the available data for odd N nuclei and given in Table IV.

Table IV. $1 - (\mu_n)_{intr} / \mu_n$ for odd-N nuclei

$p_{3/2}$	$^{9}_{4}Be$	$^{12}_{6}C$**		$p_{1/2}$	$^{13}_{6}C$	$^{15}_{8}O$**	$^{57}_{26}Fe$			
	0.38452	0.463159			-0.101466	(-0.1273)	0.8579			
$d_{5/2}$	$^{17}_{8}O$	$^{25}_{12}Mg$		$d_{3/2}$	$^{21}_{10}Ne$	$^{33}_{16}S$	$^{35}_{16}S$			
	0.010106	0.55301			1.58124	0.43	(0.129)			
$f_{7/2}$	$^{41}_{20}Ca$	$^{43}_{20}Ca$	$^{45}_{22}Ti$**	$^{49}_{22}Ti$	$f_{5/2}$	$^{47}_{22}Ti$	$^{49}_{24}Cr$**			
	0.166438	0.31146	(0.9503)	0.42312		1.5768	(1.348)			
$f_{5/2}$	$^{61m}_{28}Ni$	$^{65}_{30}Zn$**	$^{67}_{30}Zn$	$p_{3/2}$	$^{53}_{24}Cr$	$^{57m}_{26}Fe$	$^{61}_{28}Ni$			
	1.5479	0.4374	0.35948		0.75204	0.91917	0.60864			
$g_{9/2}$	$^{73}_{32}Ge$	$^{83}_{36}Kr$	$^{85}_{36}Kr$**	$^{87}_{38}Sr$	$d_{5/2}$	$^{91}_{40}Zr$	$^{95}_{42}Mo$			
	1.75226	0.49282	0.4752	0.42855		0.31883	0.5225			
$d_{5/2}$	$^{97}_{42}Mo$	$^{99}_{44}Ru$	$^{101}_{44}Ru$	$^{105}_{46}Pd$	$^{107}_{48}Cd$**	$^{109}_{48}Cd$**				
	0.5125	0.674	0.64	0.6782	0.6778	0.5668				

$h_{11/2}$	$^{113m}_{48}Cd$	$^{115m}_{48}Cd$	$s_{1/2}$	$^{111}_{48}Cd$	$^{113}_{48}Cd$ *	$^{115}_{48}Cd$ **	$^{115}_{50}Sn$
	0.43498	0.45438		0.6980	0.67461	0.6618	0.52019

$s_{1/2}$	$^{117}_{50}Sn$	$^{119}_{50}Sn$	$^{129}_{50}Xe$	$f_{7/2}$	$^{141}_{58}Ce$ **	$^{143}_{60}Nd$	$^{145}_{60}Nd$
	0.47731	0.54307	0.59390		0.494	0.4438	0.6585

$f_{7/2}$	$^{177}_{72}Ho$	$g_{7/2}$	$^{167}_{68}Er$	$p_{1/2}$	$^{169}_{68}Er$ **	$^{171}_{70}Yb$	$f_{5/2}$ $^{173}_{70}Yb$
	1.32		1.3795		0.196	0.2277	1.4973

$p_{1/2}$	$^{187}_{76}Os$	$^{195}_{78}Pt$	$^{195}_{80}Hg$ **	$^{197}_{80}Hg$ **	$^{199}_{80}Hg$	$p_{3/2}$	$^{189}_{76}Os$
	0.89915	0.049823	0.15638	0.17835	0.2118		1.34283

$p_{3/2}$	$^{201}_{80}Hg$	$i_{13/2}$	$^{195m}_{80}Hg$ **	$^{197m}_{80}Hg$ **	$f_{5/2}$	$^{203}_{80}Hg$ **	$p_{1/2}$ $^{207}_{82}Pb$
	0.70906		0.45171	0.46086		0.385	0.07571

$d_{5/2}$	$^{229}_{90}Th$	$^{233}_{92}U$ **	$^{241}_{94}Pu$ **	$f_{7/2}$	$^{235}_{92}U$ *	$^{239}_{94}Pu$ **
	1.18	0.72	0.64		0.82	1.105

From this tabulation one finds in general an expected increase with the increase of (k+n) by taking a fraction of this integer, since only a fractional number of outermost sub-shell nucleons contribute to the correction of the m.m. of the host neutron as we have seen already. One may also learn from the tabulation that the intrinsic m.m. get increase in oppose to (k+n) where outermost shell approach to close the shell. The delicate orbit dependence and (k+n) dependence should be taken into account case by case. Our tabulation can be used for such a purpose if it becomes necessary. The nucleons try to distribute nuclear surface in a cooperative manner with the central force and compete with the null-instanton induced force. The latter will admit the local concentration of neutron distribution on core surface. If the nuclear deformation is induced by the interaction between outermost nucleons and inner nucleons its influence should be reflected to the m.m. under study. The big deformation seems to occur for odd N nucleus where outermost neutrons belong to an orbit far from the closed shell will play an essential role as far as our m.m. analysis tells. We do not wish to spend a more time on this problem in this paper. We shall leave an estimate of a fractional multiplicative number to (k+n)

term for individual data, which may be done first for odd N then for odd Z nuclei by taking Coulomb effect in mind.
In this manner one can find the role of the Coulomb effect in each nucleus from a null-instanton-induced-model point of view.

$.5 Magnetic Moments of Odd-Odd Nuclei

We shall present first the Schmidt magnetic moment (abbreviated m.m. hereafter) formulae for odd-odd nuclei in order to extract the intrinsic m.m. from them. Owing to its simplicity of derivation we shall explain shortly in words and then give the result explicitly. Let us specify the spin operators $\vec{j}_p$ and $\vec{j}_n$ for unpaired single proton and neutron part, respectively. The total m.m. operator becomes the sum of Schmidt operators for p and n by replacing the total spin of the corresponding single particle operators by these spin operators for the proton and neutron part, respectively. Firstly one takes scaler products of this operator with the dyads $\vec{j}_p\vec{j}_p/\vec{j}_p^{\,2}$ and $\vec{j}_n\vec{j}_n/\vec{j}_n^{\,2}$ for the proton and neutron part, respectively, and then take that operator with the dyad $\vec{J}\vec{J}/\vec{J}^{\,2}$ made from the total spin $\vec{J}$. The coeeficient of the total spin $\vec{J}$ thus calculated gives the desired m.m. Specifying the orbital angular momentum of unpaired proton and neutron by $\vec{l}_p$ and $\vec{l}_n$, respectively, the results are summarized as follows for the sake of reader.

For simplicity let us define first the four quantities by

$$x(l_p-)=\frac{\mu_p+j_p-\frac{1}{2}}{j_p}(J(J+1)+j_p(j_p+1)-j_n(j_n+1)), \qquad (5.1a)$$

$$X(l_n-)=\frac{\mu_n}{j_n}(J(J+1)+j_n(j_n+1)-j_p(j_p+1)), \qquad (5.1b)$$

$$X(l_p+)=\frac{-\mu_p+j_p+\frac{3}{2}}{j_p+1}(J(J+1)+j_p(j_p+1)-j_n(j_n+1)), \qquad (5.1c)$$

$$X(l_n-)=\frac{\mu_n}{j_n+1}(J(J+1)+j_n(j_n+1)-j_p(j_p+1)), \qquad (5.1d)$$

where μ_p and μ_n should be understood as the intrinsic m.m. of proton and neutron, respectively, although we avoided their explicit specification for simplicity, and $l_p- = j_p -1/2$, $l_n- = j_n -1/2$, $l_p+ = j_p +1/2$ and $l_n+ = j_n +1/2$. Making use of these notations one finds four possible m.m. of interest as

$$\mu(l_p-,l_n-) = \frac{1}{2(J+1)}[X(l_p-) + X(l_n-)], \qquad (5.2a)$$

$$\mu(l_p-,l_n+) = \frac{1}{2(J+1)}[X(l_p-) - X(l_n+)], \qquad (5.2b)$$

$$\mu(l_p+,l_n-) = \frac{1}{2(J+1)}[X(l_p+) + X(l_n-)], \qquad (5.2c)$$

$$\mu(l_p+,l_n+) = \frac{1}{2(J+1)}[X(l_p+) - X(l_n+)]. \qquad (5.2d)$$

The proton and neutron part of these expressions coincide with the Schmidt formula by making an appropriate substitution as well be checked easily.

The derivation of the intrinsic magnetic moment of odd-odd nucleus including the correction term is a little bit complicated, especially the internucleon effect between unpaired single proton and neutron. It is impossible to find out experimentally the intrinsic m.m. of proton and neutron separately, while the theory does. It will be instructive to present them separately and combine them suitably in referring to the Schmidt formulae, Eqs.(5.2a)-(5.2d). They are given by

$$(\mu_p)_{intr} = [1- \frac{2[(19c0+15c1)-(15a0+19a1)c1]\mu}{27\mu_p}]\ \mu_p\ , \qquad (5.3a)$$

$$(\mu_n)_{intr} = [1+ \frac{2[(2c0+15c1)-(15a0+2a1)c1]\mu}{27\mu_n}]\ \mu_n\ . \qquad (5.3b)$$

We used the same notation as that for odd A nuclei. One sees a slight difference in a pure photon and intranucleon q-q pair term in these expressions for the last single proton and neutron, which is caused from the u and d charge effect in the combined states.

It may be convenient to give the full expression including the correction terms arising from the outermost p and n associated nucleons, as we experienced already, viz., an expected fractional (k+n) effect. We shall again show them intrinsic p and n part, separately, in spite of a lengthy equation. One finds

$$(\mu_p)_{intr} = [1- \frac{\mu}{54\mu_p}\{4(19c0+15c1)+162(k+n)(c0+c1)-4(15a0+18a1)c1$$
$$- 81(k+n)(c0+c1)(a0+a1)\}]\ \mu_p\ , \qquad (5.4a)$$

$$(\mu_n)_{intr} = [1+ \frac{\mu}{54\mu_d}\{4(2c0+15c1)+108(k+n)(c0+c1)-4(15a0+2a1)c1$$
$$-81(k+n)(c0+c1)(3a0+a1)\}]\ \mu_n\ . \qquad (5.4b)$$

One may find three types of combination of intrinsic p and n m.m. from Schmidt formula which will directly be compared with the experimental information, viz., type (1) a pure p+n, (2) a pure p-n and (3) an isolated p intrinsic m.m. with some numerical coefficient depending upon to which orbit they belong. In the last example orbital effect cancels for intrinsic n m.m. We shall choose them from the available experimental information and give them in tabular form after multiplying inverse of the orbital coefficient by dividing it by the corresponding combined free m.m. The result is tabulated in Table V.

Table V. $(Schmidt)_{intr} / (m.m.)_{free}$ for odd-odd Nuclei

nucleus	shell assignment	J^P	type	$(Schmidt)_{intr} / (m.m.)_{free}$
${}^{2}_{1}H$	$(s_{1/2}, s_{1/2})$	1^+	(1)	0.974545
${}^{10}_{5}B$	$(p_{3/2}, p_{3/2})$	3^+	(1)	0.91007
${}^{14}_{7}N$	$(p_{1/2}, p_{1/2})$	1^+	(1)	0.89715
${}^{38}_{19}K^{**}$	$(d_{3/2}, d_{3/2})$	3^+	(1)	1.3735
${}^{60}_{27}Co^{**}$	$(f_{7/2}, p_{3/2})$	5^+	(1)	0.898

$^{6}_{3}Li$	$(p_{3/2}, p_{3/2})$	1^+	(1)	1.66631
$^{8}_{3}Li$**	$(p_{3/2}, p_{3/2})$	2^+	(1)	1.6820
$^{20}_{9}F$**	$(d_{5/2}, d_{5/2})$	2^+	(1)	3.745
$^{22}_{11}Na$**	$(d_{5/2}, d_{5/2})$	3^+	(1)	1.03
$^{24}_{11}Na$**	$(d_{5/2}, d_{5/2})$	4^+	(1)	0.12815
$^{32}_{15}P$**	$(s_{1/2}, d_{3/2})$	1^+	(2)	0.1072
$^{36}_{17}Cl$**	$(d_{3/2}, d_{3/2})$	2^+	(1)	0.24261
$^{44m}_{21}Sc$	$(f_{7/2}, f_{7/2})$	6^+	(1)	1.74
$^{46}_{21}Sc$**	$(f_{7/2}, f_{7/2})$	4^+	(1)	2.63
$^{50}_{23}V$*	$(f_{7/2}, f_{7/2})$	6^+	(1)	1.028
$^{52}_{25}Mn$*	$(f_{7/2}, f_{7/2})$	6^+	(1)	0.6678
$^{54}_{25}Mn$**	$(f_{7/2}, p_{3/2})$	3^+	(3)	0.305
$^{56}_{25}Mn$**	$(f_{7/2}, p_{3/2})$	3^+	(3)	0.27949
$^{64}_{29}Cu$	$(p_{3/2}, f_{5/2})$	1^+	(1)	3.26

In this tabulation the first, second and the third column specify the nucleus, the assumed configuration for the host p and n, and the observed spin and parity, the fourth column is the corresponding type of combined intrinsic m.m. and the last column is the ratio of the intrinsic m.m. divided by the corresponding combined free p and n as well as p m.m. The deviation from the Schmidt limit will be defined as the excess from 1 in term of the newly appeared ones. Such a big

correction is almost ubiquitous phenomena in the mass range studies.

$.6 Conclusion and Discussion

We have studied in this paper the binding and excitation energies of light nuclei, the alpha-particle model, the magnetic moments of odd-A and odd-odd nuclei as the eigenvalue problem of the null-instanton induced potential. The null-instanton model suggests the strong relative distance dependence between the quarks concerned in their exchange machinery and it diminishes quickly with increase of it. The binding as well as the excitation energies of light nuclei were explained by taking the nucleon distribution in their parent nucleus into account, where the existence of the sign inverted partner of any state is suggested based on the property of the assumed potential. An quick increase of the paste quark pairs in the alpha particle model suggests a necessity to solve the spatial distribution of alpha particles in A > 16 in order to explain the associated binding energy.

The general formulae which include the corrections to the intrinsic part of the magnetic moments in Schmidt model have been derived for odd-A and odd-odd nuclei. These formulae explain naturally the well known quenching of magnetic moments of odd-A nuclei. A numerical detail depends on how much fraction of the outermost nucleons contribute effectively to the mechanism. The clear evidence is found from odd N data with outermost sub-shell structure variation dependence in certain orbital state analyzed individually. The effect becomes obscured for odd Z nucleus due to the more or less equal inner nuclear surface distribution of protons by the Coulomb effect. The numerical value can be compared with the theoretical parameters are tabulated for each nucleus but a detail is left untouched. The explanation of the anti-quenching is also reduced to the choice of the sign of the theoretical parameters, we hope that it may also be realized by the parameters introduced in the particular mass number region of interest. The correction to the intrinsic part of the Schmidt limit for the odd-odd nuclei is tabulated in units of the corresponding combination of free nucleons. An apparently big anti-quenching found in this table may not become important if it is done in units of the free proton magnetic moment. Our presentation of magnetic moment data

will be related intimately to the nuclear structure and it should be pursued further in connection with an independent information on the properties of individual nucleus.

The null-instanton model presented in this paper can be applied wide range of problems. We have concerned so far on the static problem on nuclei in this paper. Our model can be extended to the dynamic problem, such as, a weak decay. There is a fundamental question on the universality of the Fermi coupling constant in any weak processes. The CVC hypothesis does not admit the weak vertex correction from strong interaction.[11] This does not admit any vertex correction for the leptonic decay of a meson. However the corrections becomes important for both semi-leptonic and hadronic decays through spectator quarks: (i) quark exchange effect between the parent quark and the spectator quark, and (ii) final state interaction in between daughter quark and associated spectator quark(s) and initial spectator ones. They are exemplified at first by the semi-leptonic decay of mesons and baryons. The former correction to nucleon beta decay is the same as the one obtained magnetic moment study for free as well as the bound nucleon, while the latter correction enters newly as a radiative correction for any semi-leptonic and hadronic decays as a similar eigenvalue problem. We hope that the universality of Fermi constant is certified for widely distributed ft value in nuclear beta decay having the same forbidness by an introduction of expected large correction arising from the quark exchange effect between the spectator quark as well as the nucleons involved in the parent nucleus and the daughter quark, spectator quarks in daughter nucleon as well as daughter nucleons in the nucleus. The large correction is expected with the increase of mass number as we experienced already through alpha particle model. Our formulation admits us to include all these corrections as an eigenvalue problem of one and the same potential in associated strong interaction without worrying about the renormalization problem inherent in quantum field theory. We did not use any QCD language in our approach. For the time being our method will be considered as a substitute of the QCD field theory for many-body problem in a low energy region. We should like to concentrate on weak interactions which will affect to find a reliable answer for the determination of the famous Cabibbo-Kobayashi-Maskawa (CKM) weak mixing angles. One learns that

the experimental values of them still distributed over some extent depending upon from which process they are derived.

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Remark: Doubly printed line 5 in page 13 should be read:

and two characters, respectively, in order to shown their origin.

One finds easily that this additional modification cannot help to