# $\beta$ decay of $^{12}B$ and $^{12}N$†

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The $\beta$ decays of $^{12}B$ and $^{12}N$ to the first excited state of $^{12}C$ have been compared by detecting the $\beta$ particles in $4\pi$ geometry and simultaneously collecting $\beta$-$\gamma$ coincidences. All events were time sorted so that the $^{12}B$, $^{12}N$ components of the $\beta$ and $\beta$-$\gamma$ coincidence spectra could be accurately extracted down to low values of the $\beta$-particle energy. For the ratio of the branching ratios $^{12}N/^{12}B$ we find $1.74 \pm 0.08$. This and other data are analyzed to extract the $\beta$-decay asymmetry $[(ft)^+/(ft)^-]-1=-0.013\pm0.066$. All existing data on asymmetry in Gamow-Teller $\beta$ decay in even-$A$ nuclei of the $1p$ shell are analyzed in the light of the work of Kubodera, Delorme, and Rho that, using appropriate many-body wave functions, relates the asymmetry to two fundamental second class current parameters $\zeta$ and $\lambda$. We find, at the 99% confidence limit, $|\zeta|<4\times10^{-3}$ MeV$^{-1}$; $|\lambda|<1.5\times10^{-2}$. In the model that ascribes all second class current effects to the $G$-parity-violating decay $\omega\to\pi e\nu$ both as an exchange term (which enters into $\lambda$) and as a nucleon decay vertex renormalization term (which determines $\zeta$ and enters into $\lambda$), our results correspond to $\tau_{\omega\to\pi e\nu}>5\times10^{-11}$ sec. This limit corresponds, using SU(3), to the expectation that the asymmetry in $\Sigma\to\Lambda e\nu$ decay is less than 25%.

[NUCLEAR STRUCTURE $^{12}B$, $^{12}N$; measured ratio of branching ratios to $^{12}C(4.44)$; deduced $\beta$-decay asymmetry. Nuclear $1p$-shell $\beta$-decay asymmetry parameters deduced second class current parameters $\zeta$, $\lambda$; calculated limits $\tau(\omega\to\pi e\nu)$, asymmetry $\Sigma\to\Lambda e\nu$.]

## I. INTRODUCTION

When mirror $|T_z|=1$ nuclei, such as $^{12}B$ and $^{12}N$, $\beta$ decay to $T=0$ states of the associated $T_z=0$ nucleus, $^{12}C$ in this case, we naively expect that, to each of the $T_z=0$ states separately, the intrinsic rate of positron emission $(ft)^+$ from the $T_z=-1$ body should be the same as that $(ft)^-$ for negative electron emission from the $T_z=+1$ body. The existence of an asymmetry:

$$\delta=[(ft)^+/(ft)^-]-1\neq0$$

in such mirror Gamow-Teller $\beta$ decay must be due either to a breakdown of mirror symmetry in the nuclear wave functions, as between $^{12}B$ and $^{12}N$ in this example (the "trivial" explanation), or more profoundly, to the existence of "second-class currents" in (presumably the hadronic part of) the fundamental weak interaction.[1] Such second-class terms, which are momentum-transfer dependent and so which appear effectively as second-forbidden corrections to the main (allowed) component of the hadronic current, transform oppositely to the main component under the $G$-parity (or $U$-spin) operator and so add to it for negative electron emission but subtract for positron emission (or vice versa) thereby engendering a finite value for $\delta$. Conservation of the vector current forbids the appearance of such second-class terms in Fermi $\beta$ decay but the axial-vector current is not conserved and so there is no fundamental reason why second-class terms should not arise there, due, perhaps, but not necessarily, to the nucleon-vertex-renormalizing effects of such mesonic decays as $\omega\to\pi e\nu$. It is clearly a matter of interest to investigate this mirror Gamow-Teller asymmetry experimentally as thoroughly as possible and to determine the relative roles of the "trivial" and "fundamental" effects.

It has been known for many years that the decay of $^{12}B$ and $^{12}N$ to the ground state of $^{12}C$ has an asymmetry of $\delta\approx0.12$ but it is only comparatively recently that it has become clear that the phenomenon may be a general one with asymmetries of comparable or larger value showing up elsewhere in several odd-$A$ as well as even-$A$ systems.[2] This observation has touched off a considerable flurry of experimental and theoretical work from which has emerged, among other things, both that much of the earlier experimental data were grossly in error and also that the earlier theoretical expectation[3] for the detailed form of the asymmetry due to a second-class term was oversimplified. It is the consequence of this second point that we now discuss since it, together with suspicion coming from the first, constitutes the motivation for the work that we report here.

For an on-mass-shell nucleon, Lorentz invariance limits the possible form of the axial weak hadronic current; one component of this current is second class and is associated with a single "induced tensor" coupling constant and proportional to the momentum transfer in the decay. If one now just adopts the impulse approximation, uses the on-shell form for the current, simply sandwiches the associated generalized operator between nucleon states as they exist in the complex nucleus, applies the Foldy-Wouthuysen transformation and passes to the limit of large energy release, the latter being pretty well justified in the interesting cases, one finds that the second class current asymmetry $\delta^{\rm SCC}$ should be proportional to an induced tensor coupling constant $\zeta$ and to the sum of the energy releases on the two sides of the mirror $W_0^+ + W_0^-$ but independent of the details of the nuclear wave functions since, in this approximation, the induced tensor and the allowed matrix elements have the same form:

$$\delta^{\rm SCC} = -\frac{4}{3}\frac{\zeta}{g_A}(W_0^+ + W_0^-),$$

where $g_A$ is the axial-vector coupling constant in units of the vector coupling constant. (In this paper we follow the notations and metric of Kubodera, Delorme, and Rho.[4]) The early analyses[2] were carried out under this approximation whose operational importance is that if the asymmetry is of fundamental origin it is determined, apart from the coupling constant, just by $W_0^+ + W_0^-$ so that if, for example, the $\beta$-branching ratios were well known for the positron-emitting body but unknown for the negative-electron-emitting body the data could still be unambiguously interpreted in terms of the value of the induced tensor coupling constant. If, on the other hand, the second class asymmetry were not determined solely by $W_0^+ + W_0^-$ such a procedure would not be available and $\delta$ determinations final state by final state would have to be made, calling therefore for sufficiently accurate branching ratio measurements from both sides of the mirror. In fact, when the off-mass-shell nature of nucleon $\beta$ decay in complex nuclei is considered the definition of the nucleon's induced tensor coupling constant becomes ambiguous and the single coupling constant of the on-shell approximation splits into two, the coupling taking the form:

$$\pm i[g_T\sigma_{\lambda\mu}(p-p')_\mu\gamma_5 + g_T' i(p+p')_\lambda\gamma_5]\tau_\pm ,$$

where $p, p'$ are the initial, final nucleon momenta. The constant $\zeta$ of the on-shell approximation is replaced by the two constants $g_T$ and $g_T'$. Setting $\zeta = g_T + g_T'$ and going on shell we gain the above expression for $\delta$ which is not correct in the context of a complex nucleus. The correct expression[4] for $\delta$ recognizes not only the off-shell point just made about the nucleon decay but also the fact that mesonic exchange currents must play an explicit and, perhaps, important role[5] through, for example, the $G$-parity-violating decay $\omega \to \pi e \nu$. Kubodera, Delorme, and Rho[4] find:

$$\delta^{\rm SCC} = -4\frac{\lambda}{g_A}J + \frac{2}{3g_A}(\lambda L - 2\zeta)(W_0^+ + W_0^-). \qquad (1)$$

Here $J$ and $L$ are matrix elements of complicated two-nucleon operators and require explicit many-body nuclear wave functions for their evaluation. The constant $\lambda$ involves $g_T'$ and also properties of the exchange meson (its strong coupling to the nucleon and its intrinsic $G$-parity-violating weak decay rate) whose second-class decay is contributing to the overall asymmetry. $(W_0^+ + W_0^-)$ proportionality of $\delta$ would be recovered only if the induced tensor current were to be conserved, in which case $\lambda = 0$.[6] This eventuality is denied by the observation[7] that in the $A = 8$ system $\delta$, although large ($\approx 0.11$) is essentially independent of $W_0^+ + W_0^-$; the which possibility is offered by the condition $\lambda L = 2\zeta$ in the above expression for $\delta^{\rm SCC}$.

We are therefore, experimentally and theoretically, put into the position where we have no simple nuclear-structure-model-independent expectation as to the $(W_0^+ + W_0^-)$ dependence of $\delta$. This means that $\beta$ branches cannot simply all be lumped together since one of low $W_0^+ + W_0^-$ may have a large $\delta$ and vice versa: each branch is of interest in its own right.

The same conclusion as to the importance of the individual branches is reached by consideration of the possible trivial explanations of a finite $\delta$. Of these trivial explanations the most obvious is that the Coulomb force results in different binding energies for the "decaying proton" in the $T_Z = -1$ body and its mirror "decaying neutron" in the $T_Z = +1$ body so that their wave functions are different with different overlap with the respective product nucleons in the $T_Z = 0$ body hence $(ft)^+ \neq (ft)^-$: $\delta \neq 0$. When it is taken into account that the parentage of the various initial and final states is not unique, viz., that there are several components to the overall $\beta$-decay amplitude with different associated nucleon separation energies and hence different asymptotic single-nucleon wave functions and overlap integrals depending on the operative parent states, we find that this "binding energy" asymmetry is expected to fluctuate widely from case to case, $A$ value to $A$ value and, within a single $A$ value, final state to final state even changing sign from one to the other.[8]

We now come immediately to the case in point. The $A=12$ system, in addition to providing a well-determined asymmetry $\delta \approx 0.12$ for the $\beta$ decay to the ground state of $^{12}$C, shows a significant decay to the first excited state of $^{12}$C at about 4.44 MeV. It therefore offers a valuable opportunity for the determination of a second $\delta$ within one $A$ system. The $(W_0^+ + W_0^-)$ value for the excited state decay is about 0.7 times that for the ground state so the simple $(W_0^+ + W_0^-)$ proportionality of the naive approximation would lead, if the ground state $\delta$ were entirely due to a fundamental second class effect, to an expectation of $\delta \approx 0.08$ for the decays to the excited state. While, as remarked above, we cannot expect to find so simple a proportionality in practice, the large value of $\delta$ which it would imply shows that the momentum transfer involved in the transition to the excited state is enough to give the possibility of large effects [see Eq. (1)].

Kubodera, Delorme, and Rho[4] have calculated the values of the relevant two-nucleon matrix elements $J$ and $L$ of Eq. (1) for the cases of $1p$-shell $\beta$ decay using the well-founded wave functions of Barker[9] for $A=8$ and 9 and of Cohen and Kurath[10] for $A=12$ and 13. As will be discussed later, the odd-$A$ cases are not as yet sufficiently well-studied from the nuclear structure point of view of estimating the trivial aspects of the asymmetry and must be set on one side. This means that only $A=8$ and 12 are at present available for quantitative analysis in an attempt to extract information separately about the two second class elements $\zeta$ and $\lambda$ of Eq. (1). $A=8$ gives two pieces of information because we know both the overall $\delta$ and also its $(W_0^+ + W_0^-)$ dependence.[7] The ground state decay of $A=12$ is a third datum but the value of the first excited state decay of $A=12$ (to the 4.44 MeV state of $^{12}$C) is greater than the mere adding of a fourth constraint on $\zeta$ and $\lambda$ since, as will be seen, the dependence of $\delta$ on $\zeta$ and $\lambda$ in that case is markedly different from the others.

Nuclear structure studies, experimental allied with theoretical, offer then the promise of determinations of the two second class coupling constants $\zeta$ and $\lambda$ independently. $\zeta$ is also accessible from elementary particle studies but $\lambda$ appears to be uniquely the province of nuclear structure since the possibilities for setting significant limits directly upon the $G$-parity-violating weak decays of the relevant mesons by direct measurement are extremely remote as will appear in our discussion.

To return to the specific case of the first excited state decay in $A=12$, from the purely experimental point of view the present situation is scarcely satisfactory. For reasons that will become obvious it is usual to discuss the situation in terms of the ratio $R$:

$$R = \frac{[I(\beta_1)/I(\beta)]_{^{12}\mathrm{N}}}{[I(\beta_1)/I(\beta)]_{^{12}\mathrm{B}}},$$

where $I(\beta_1)$ and $I(\beta)$ are the intensity of the $\beta$ transition to the first excited state and the total $\beta$-decay intensity, respectively, i.e., $R$ is the ratio of the first excited state branching ratios. Three measurements of $R$ are available at present, apart from the one we report here: Two older measurements are $R=1.72\pm0.15$[11] and $R=1.84\pm0.10$[12] which are concordant but which lead to $\delta=-0.117\pm0.041$ which is so far out of line with the systematics of the asymmetry phenomenon that a remeasurement was obviously needed[8]; the third, more recent, measurement[13] gives $R=1.52\pm0.06$ which corresponds to the more comfortable value $\delta=0.06\pm0.04$. The wide disparity between the older and the newer measurements calls for a further redetermination, particularly in view of the special importance of this transition to which we have alluded and which we shall enlarge upon later.

## II. EXPERIMENT

We measure in our experiment the relative intensity of the $\beta$ branching to the $^{12}$C first-excited state ($E_x=4.44$ MeV) in terms of the total number of $\beta$ particles emitted $I(\beta_1)/I(\beta)$ for the $T_Z=-1$ and the $T_Z=+1$ parents, $^{12}$N and $^{12}$B, respectively. From these quantities we deduce the ratio $(ft)^+(^{12}\mathrm{N}\to{}^{12}\mathrm{C}_{4.44})/(ft)^-(^{12}\mathrm{B}\to{}^{12}\mathrm{C}_{4.44})$. The experiment is greatly simplified in that the various $\beta$ branches do not have to be counted with known absolute efficiencies; the essential requirement is that the efficiencies be the same for both $^{12}$B and $^{12}$N since all we need is the ratio $R=[I(\beta_1)/I(\beta)]_{^{12}\mathrm{N}}/[I(\beta_1)/I(\beta)]_{^{12}\mathrm{B}}$.

Thus the most convenient experimental arrangement from our point of view is one in which both $\beta$-ray events and $\beta$-$\gamma$ (4.44) coincidence events may be measured simultaneously; the $\beta$-$\gamma$ coincidence events are the signature of branching to the 4.44-MeV level, while the $\beta$-ray counting provides a measure of the total activity. The various $\beta$ activities have different end-point energies, and corrections due to counting efficiency dependence on $\beta$-ray energy need to be carefully accounted for. The counting geometry will be described first, then the main features of the counting pro-

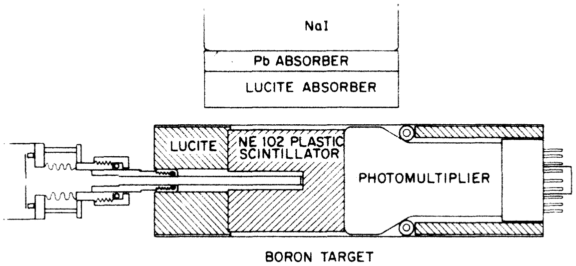


FIG. 1. The experimental arrangement (to scale).

cedure. The $^{12}$N and $^{12}$B activities were produced with the $^{10}$B($^3$He, $n$)$^{12}$N ($Q$ = 1.57 MeV) and $^{11}$B($d, p$)-$^{12}$B ($Q$ = 1.14 MeV) reactions. The incident ion beams at the bombarding energies $E_{^3\mathrm{He}}$ = 2.4 MeV and $E_d$ = 0.40 MeV were provided by the Lockheed Palo Alto Research Laboratory 3 MV Van de Graaff generator. The $\beta$ rays were counted in "4$\pi$" geometry obtained with a well-type plastic scintillator, located over the end of the accelerator beam tube (see Fig. 1). The scintillator was a right cylinder of NE 102,[14] 6.99 cm diam by 7.94 cm long. A cylindrical hole, 1.27 cm diam by 5.08 cm long, was drilled axially into the cylinder. A 7.62 cm diam photomultiplier tube viewed one end of the cylinder. This tube was especially selected and tested to exhibit low-gain shift with fluctuating counting rates. To ensure mechanical stability, the other end of the scintillator was attached to a cylinder of Lucite, 7.46 cm long with a 1.59 cm diam bore. The entire package was wrapped in netic and conetic foil to provide magnetic shielding[15] and made up into a rigid package readily mounted and positioned about the target beam tube. The target tube consisted of stainless steel, 0.625 cm diam, of wall thickness 0.015 cm, with an end plug 0.005 cm thick. A NaI(Tl) detector 12.7 cm diam by 15.24 cm long, located at 90° to the incident beam direction and 8.89 cm from the reaction site, was used to detect $\gamma$ radiation. Both a 2.54 cm thick Lucite plate and 1.35 cm thick Pb shield covered the front face of this detector. Rigid mounting of detectors together with optical alignment of target tube and collimator ensured a constant source to detector(s) distance throughout the experiment. The 4$\pi$ geometry is rather convenient for the present experiment in that it minimized systematic bias between counting $\beta^-$ and $\beta^+$ particles from $^{12}$B and $^{12}$N, respectively. This bias is due to the sensitivity of the plastic scintillator to $\gamma$ radiation as well as $\beta$ rays; thus positrons with attendant annihilation radiation have a greater chance of being counted than negative electrons. In 4$\pi$ geometry $\beta^+$ particles reaching the scintillator give rise to only one count, whether or not the associated annihilation radiation is detected. (The increase in count rate due to detecting annihilation radiation from positrons that do not reach the scintillator is estimated to be <1%.) The targets were produced using separated isotopes of $^{10}$B and $^{11}$B. A measured amount of the isotope was packed into the beam tube; typically, target thickness amounted to 175 mg/cm$^2$. Thick targets have the disadvantage of increasing the mass at the target spot, thus increasing the amount of absorbing material between radioactive body and the counter; however, they provide a convenient contaminant-free beam stop and they could be readily fabricated in this geometry. We calculate that $\beta$'s with $E_\beta > 0.7$ MeV reach the scintillator.

The counting procedure is straightforward. With the aid of a mechanical shutter, the incident beam is switched onto the target for a fixed bombardment interval after which the beam is interrupted and the counting interval begun. The pulse-height distribution of both $\beta$ rays and $\beta$-$\gamma$ coincidence events were recorded together with the time after bombardment; following the counting interval, the cycle was repeated. The counting system was composed of commercially available electronics, with analog-to-digital converters and scalers interfaced to an SEL 810A computer. Clock time was obtained from the computer clock. Full advantage of the flexibility of the system was taken in data collection: recorded data consisted of both the matrix of $\beta$-ray pulse height [$N(E_\beta)$] vs time after bombardment ($t$) recorded in computer core memory, and also the matrix of coincident $\beta$-$\gamma$ pulse heights $N(E_\beta)$, $N(E_\gamma)$, and $t$, recorded event by event onto magnetic tape. The former spectrum resembles a traditional "multiscale" spectrum except that $\beta$-ray pulse height is also recorded. In this experiment dead time corrections are important for subsequent manipulations of the data matrix. Primary dead time correction was done in the following way: every linear pulse presented to the analog-to-digital converter (ADC) was accompanied by a logic pulse. These logic pulses were counted with negligible counting loss and recorded as a part of the multiscaling process. Thus a record was kept of every pulse presented to the ADC during each time interval, and so dead time corrections become straightforward. This procedure is necessary since dead time corrections based upon storing a constant rate pulser along with the data are inadequate when the count rate is fluctuating rapidly. It remains to ensure that the efficiency of recording coincidence events and the efficiency of recording the $\beta$ spectrum are identical. The computer program for this experiment was developed with this point in mind; as a check, linear pulses (injected at the preamplifiers with a constant rate) were stored along with events due to reaction-induced activities. Within experimental error, the loss of coincidence pulses matched the loss of single pulses. The beam pulsing system was composed of a mechanical shutter together with an optical readout system. While no doubt an electrostatic deflection system would also have sufficed, use of a mechanical beam shutter eliminates any possibility of stray beam striking the target during any part of the counting interval. The state of the shutter was read out directly using a light source on one side of the shutter and a photosensitive transistor on the other. A pulse derived from this system was included in the counting system logic: no counting was done with beam on target either at the beginning or at the end of a count cycle.

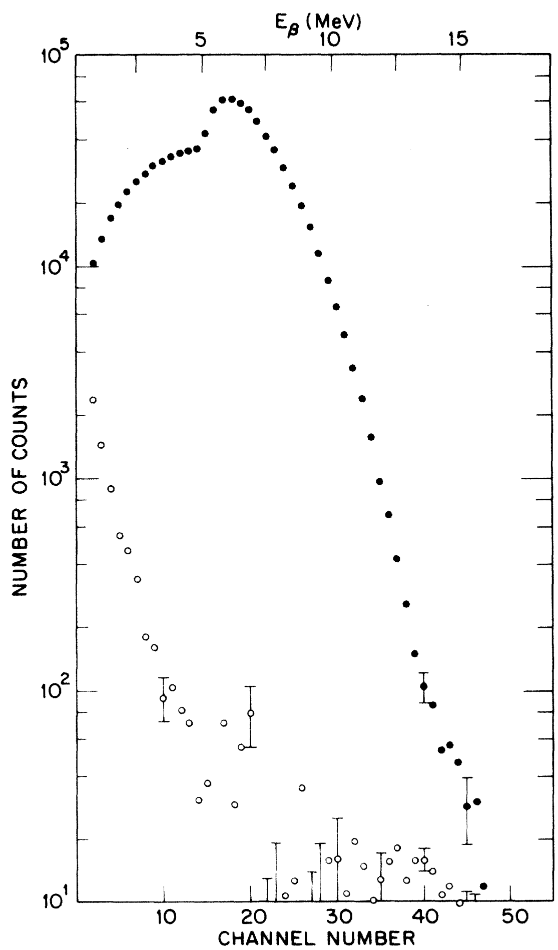


FIG. 2. The β-ray spectrum collected in the run labelled in Table I as $^{12}$B, Run 1. The spectrum is decomposed into the contribution due to $^{12}$B (solid points) and long-lived background (open points). (The peak around $E_\beta \approx 7$ MeV is composed of "folded back" pulses due to β particles that pass through the plastic scintillator.)

Extraction of the ratio of the number of β-γ coincidences, $N_{\beta-\gamma}$, to the number of β's, $N_\beta$, for both $^{12}$B and $^{12}$N is described next, beginning with a description of the extraction of the raw experimental number followed by a description of the various corrections due to target thickness and detection bias. Remember that both the two-parameter matrix, $N(E_\beta)$ vs $t$, and the three-parameter matrix $N(E_\beta)$, $N(E_\gamma)$, and $t$, were collected simultaneously and with the same dead time; thus to deduce $N_{\beta-\gamma}$ it is most convenient to deduce $N_\beta$ from the number of counts in the two-parameter matrix due to $^{12}$N ($^{12}$B) decays only and $N_{\beta-\gamma}$ from the intensity of the 4.44-MeV γ rays in the three-parameter matrix. To begin with, the matrix $N(E_\beta)$ vs $t$ was analyzed using a multicomponent half-life fitting program; it was found that above $E_\beta = 1005$ (1340) keV for $^{12}$N ($^{12}$B) the decay curve was well represented by a single exponential decay together with a constant background component, after dead time corrections were made. (By "constant" we mean long-lived compared with the $^{12}$B and $^{12}$N half-lives of 20 and 11 msec, respectively.) A typical β-ray spectrum for $^{12}$B is illustrated in Fig. 2 (in terms of electron energy at decay). Using this description, the amount of constant background in the one-line data matrix was calculated and subsequently $N_\beta$ due to $^{12}$N ($^{12}$B) alone was extracted from the two-

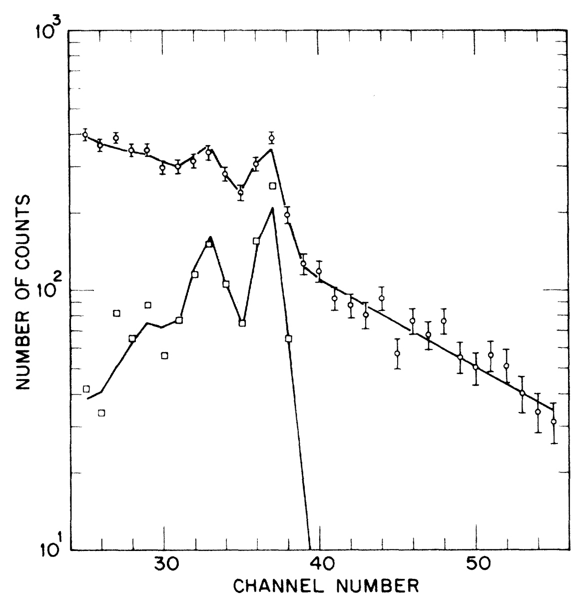


FIG. 3. The γ-ray spectrum (circular points) collected with the β-γ coincidence condition imposed collected in the run $^{12}$N, Run 3 of Table I. After the contribution due to bremsstrahlung and annihilation in flight is subtracted, the data represented by square symbols remain. The lines represent the computer fit to these data (point connected) derived as described in the text.

parameter data matrix. Next, the three-parameter matrix $N(E_\beta)$, $N(E_\gamma)$, and $t$ was integrated over the same $\beta$-ray energy interval and time interval as the two-parameter matrix. The resulting $\gamma$-ray pulse-height distribution is illustrated in Fig. 3. We see that the spectrum for one of the $^{12}N$ runs (the $^{12}B$ $\gamma$-ray coincidence spectra are distinctly superior) is dominated by bremsstrahlung and positron-annihilation-in-flight radiation, rather than the contribution due to the 4.44 MeV $\gamma$ ray. This spectrum illustrates the compromises made in our arrangement. Firstly, the use of a $4\pi$ $\beta$ counter means relatively higher bremsstrahlung background as compared with, e.g., a $2\pi$ arrangement having the NaI(Tl) counter and $\beta$ counter at 180° with respect to each other. Secondly, digitizing and storing the matrix $N(E_\beta)$ vs $t$ with our computer system is a relatively slow process compared with the standard multiscale procedure. With the additional self-imposed constraint of keeping instantaneous system counting rate losses to <20%, the data collection process is quite tedious (the spectrum of Fig. 2 represents 12 h of running time). As may be obvious, the chief experimental error will be due to the uncertainty in the intensity of the 4.44 MeV $\gamma$ ray. The yield of the 4.44 MeV $\gamma$ ray was obtained by the following procedure: a computer program was developed based on a nonlinear least-squares fit of an arbitrary line shape plus an exponential background. The line shape chosen represents the response of the NaI(Tl) detector to 4.44 MeV $\gamma$ radiation, while the exponential component represents the bremsstrahlung, etc. Using a four-point interpolation scheme both the position and amplitude of the standard line shape were allowed to vary as well as both amplitude and slope of the exponential background. The 4.44-MeV $\gamma$-ray line shape was obtained in our experimental arrangement rather simply, using the $^{10}B(^3He, p)$-$^{12}C$ reaction, and measuring $p$-$\gamma$ coincidences. With this description normalized $\chi^2$ values $=1.5 \pm 0.5$ resulted typically from the least-squares fit. Thus, the parametrization works reasonably well; the fitted $\gamma$-ray spectrum after the exponential is subtracted is illustrated by the curve in the lower portion of Fig. 3. ($\gamma$-ray yields deduced from the raw data after simple subtraction of an exponential background agreed to within 2% with the yields arrived at using the calculated line shape.) These data are summarized in Table I.

## III. ANALYSIS

The quantity $R' = [N_{\beta-\gamma}/N_\beta]_{^{12}N}/[N_{\beta-\gamma}/N_\beta]_{^{12}B} = 1.79 \pm 0.08$ obtained directly from the data of Table I must be corrected for counting losses due to electronic bias and target absorption effects in order to yield the desired $R = [I(\beta_1)/I(\beta)]_{^{12}N}/[I(\beta_1)/I(\beta)]_{^{12}B}$. The following procedure was used: the losses sustained by the $^{12}B$ and $^{12}N$ $\beta$-ray spectrum were estimated by comparing the measured (life-time-corrected) spectra above channel 3 (4) for $^{12}N$ ($^{12}B$) with theoretically-computed spectra allowing for $\beta$-ray energy loss in the target assembly etc. An empirical correction factor, as a function of channel number, by which the observed counts must be multiplied to recover the theoretical spectra was deduced; this factor changed slowly and smoothly with channel number. Counts below the cutoff channel were deduced by extrapolation of this correction factor and use of the theoretical computed spectra, hence the counting losses below the cutoff channel were accurately estimated. An extrapolation of these loss factors gave the corresponding factors for the $^{12}B(4.44)$ and $^{12}N(4.44)$ spectra, and hence the final overall correction: $R/R' = 0.970$. In a sec-

TABLE I. Summary of experimental results.

| Body | Run | $N_\beta$ $(\ )\times 10^6 \pm (\ )\times 10^4$ | | $N_{\beta-\gamma}$ $(\ )\times 10^3 \pm (\ )\times 10^0$ | | | $N_{\beta-\gamma}/N_\beta$ $\times 10^{-3}$ |
|---|---|---|---|---|---|---|---|
| $^{12}N$ | 1 | 5.710 | 1.41 | 1.244 | 96 | | 0.217 ± 0.017 |
| | 2 | 4.046 | 2.28 | 0.827 | 58 | | 0.204 ± 0.014 |
| | 3 | 4.737 | 1.38 | 1.044 | 62 | | 0.220 ± 0.013 |
| | 4 | 2.928 | 0.98 | 0.588 | 49 | | 0.200 ± 0.016 |
| | 5 | 4.522 | 1.31 | 0.859 | 62 | | 0.189 ± 0.013 |
| | | | | | | Average | 0.206 ± 0.006 |
| $^{12}B$ | 1 | 4.658 | 0.99 | 0.541 | 53 | | 0.116 ± 0.011 |
| | 2 | 4.980 | 1.10 | 0.592 | 57 | | 0.118 ± 0.011 |
| | 3 | 3.127 | 0.81 | 0.353 | 41 | | 0.112 ± 0.013 |
| | 4 | 5.227 | 1.11 | 0.590 | 38 | | 0.112 ± 0.007 |
| | 5 | 6.845 | 1.21 | 0.812 | 53 | | 0.118 ± 0.007 |
| | | | | | | Average | 0.115 ± 0.004 |

ond approach a purely empirical polynomial fitting to the observed $\beta$ spectra was employed to extrapolate the observed spectra back to low-channel numbers. This leads to a correction factor 0.969. The individual loss factors are somewhat different for the two procedures but we arrive at almost the same overall correction factor since the actual loss factor differences cancel out to a high degree. The agreement between final results of the two very different approaches to this important problem gives us a good confidence in the reliability of the $R' \to R$ correction. (Other corrections due to absorption and coincidence processes are negligible in their effect on $R$.)

In this fashion we arrive at $R = 1.74 \pm 0.08$. Comparing this number with other measurements of $R$ we find rather poor agreement with the recent result of Alburger[13] ($R = 1.52 \pm 0.06$) but apparently better consistency with the earlier and less accurate values of $1.72 \pm 0.15$ [11] and $1.84 \pm 0.10$.[12]

As compared with Alburger's measurement[13] we have the advantage that our time sorting of the $\beta$ spectra and $\beta$-$\gamma$ coincidence spectra enables us to extract with confidence the components due to $^{12}B$ or $^{12}N$ down to very low $E_\beta$ values and so to correct for other induced activities. On the other hand our system is intrinsically rather slow so that our limit is in the statistics rather than in the analysis of the data; in Alburger's case the statistics were relatively excellent and the limit lay in the data analysis. We find it difficult to comment on the discrepancy between the two careful recent experiments and feel that the best present value for $R$ must split the difference between them and increase the error to include both; we therefore use:

$$R = 1.63 \pm 0.11$$

in the subsequent analysis.

The selection of data to be included in $R$ may be criticized as arbitrary. We justify our procedure by pointing out (i) that the recent measurements were designed and optimized to measure only the quantity $R$. (ii) The corrections to the $\beta$ spectrum to account for counting losses due to electronic cutoff effects, target thickness, etc., were smaller as a result of (i) than in earlier work. (iii) Recent improvements in electronic equipment in terms of stability with respect to large fluctuations in counting rate (such as encountered in these experiments) are just those required to make experiments of this type more reliable. (iv) Finally, we note that had we taken the weighted average of all four data points, $R_{\rm ave} = 1.65 \pm 0.04$, we should in any case have been forced to arbitrarily increase the error in $R_{\rm ave}$ since for these data points $\chi^2$/(degrees of freedom) = 3.3.

We now wish to compute $\delta$ for $^{12}N$ and $^{12}B$ transitions to the first excited state of $^{12}C$ using this value of $R$. We also need, but are not sensitively dependent on, the absolute branching ratios to this and other, higher, states of $^{12}C$. For the $^{12}B$ branch to the 4.44 MeV state we combine Alburger's recent figure[13] of $(1.27 \pm 0.06)\%$ and that recomended by Ajzenberg-Selove and Lauritsen[16] namely $(1.33 \pm 0.09)\%$ finding $(1.29 \pm 0.05)\%$. For the other $\beta$ branches we use the recommended[16] values as also for the half-lives, $20.41 \pm 0.06$ and $10.97 \pm 0.04$ msec for $^{12}B$ and $^{12}N$, respectively. The maximum $\beta$-particle kinetic energies in the ground state decays, $13\,370.4 \pm 1.3$ and $16\,322 \pm 5$ keV for $^{12}B$ and $^{12}N$, respectively, we take from the standard mass tables.[17] For the excitation of the first excited state of $^{12}C$ we use 4439.2 $\pm 0.3$ keV.[16]

For the $f$ values themselves we use a recent parametrization[18] that takes into account finite nuclear size effects, including the convolution of the nucleon and lepton wave functions through the nuclear volume, screening, the "outer" (energy-dependent) radiative correction of order $\alpha$ and the finite mass effect. This parametrization has an intrinsic uncertainty of less than 0.1% and,

TABLE II. $ft$ values for the $\beta$ decay of $^{12}B$ and $^{12}N$ to the ground and first-excited (4.44 MeV) states of $^{12}C$.

| Decaying body | Final state of $^{12}C$ | $E_\beta$(max) [a] (keV) | $f$ | Branch (%) | $t_{1/2}$ (msec) | $ft$ (sec) |
|---|---|---|---|---|---|---|
| $^{12}B$ | Ground | $13\,370.4 \pm 1.3$ | $(5.6113 \pm 0.0026) \times 10^5$ | $97.13 \pm 0.31$ | $21.01 \pm 0.09$ | $(1.1789 \pm 0.0051) \times 10^4$ |
| $^{12}B$ | 4.44 | $8931.2 \pm 1.3$ | $(8.1739 \pm 0.0056) \times 10^4$ | $1.29 \pm 0.05$ | $1582 \pm 60$ | $(1.293 \pm 0.049) \times 10^5$ |
| $^{12}N$ | Ground | $16\,322 \pm 5$ | $(1.1327 \pm 0.0017) \times 10^6$ | $94.45 \pm 0.48$ | $11.61 \pm 0.07$ | $(1.3151 \pm 0.0082) \times 10^4$ |
| $^{12}N$ | 4.44 | $11\,881 \pm 5$ | $(2.4452 \pm 0.0050) \times 10^5$ | $2.10 \pm 0.16$ | $522 \pm 40$ | $(1.276 \pm 0.098) \times 10^5$ |

[a] This quantity is not strictly the maximum kinetic energy of the electron since it neglects the nuclear recoil. This omission is allowed for in the computation of $f$ which it lowers by the factor $1 - 3W_0/(2A)$ where $W_0$ is the total energy release in $mc^2$ units and $A$ the mass value in the same units. This correction ranges up to 0.22% in the present examples and exceeds the error due to mass uncertainty.

together with the above-quoted energetics, gives the $f$ values of Table II. The above-quoted data, together with the $R$ value now recommended, lead to the branching ratios, and, with the adopted half-lives, the $t$ values and hence the $ft$ values of Table II.

For the asymmetry in the ground state decay we cannot do better than use the separate $ft$ values of Table II finding:

$$\delta_{g.s.} = 0.1155 \pm 0.0085 .$$

However, for the decay to the first excited state we are only weakly sensitive to the absolute $\beta$-branching ratios since the decays are predominantly to the $^{12}C$ ground state; we may therefore use $R$ almost directly to find:

$$\delta_{4.44} = -0.013 \pm 0.066 .$$

## IV. DISCUSSION

Our objective is now to analyze the four pieces of experimental data reliably available in the $1p$ shell in order to extract values for the two fundamental second-class quantities $\zeta$ and $\lambda$ that appear in Eq. (1) for the asymmetry $\delta$, using the computations of Kubodera, Delorme, and Rho[4] that link these quantities with the nuclear structure. These data are:

$$\delta_8 = 0.107 \pm 0.011,$$ [7]

$$d\delta_8/d(W_0^+ + W_0^-) = (0 \pm 6) \times 10^{-4} \text{ MeV}^{-1},$$ [19]

$$\delta_{12} = 0.1155 \pm 0.0085 ,$$

$$\delta_{12^*} = -0.013 \pm 0.066 .$$

The situation seems well overdetermined since each datum depends only on $\zeta$ and $\lambda$; however, the impediment to an immediate extraction of these second class quantities is the "trivial" nuclear structure effects already mentioned in the Introduction. These effects give rise to asymmetries that we call $\delta^{\text{nucl}}$ so that the fundamental second class current asymmetry $\delta^{\text{SCC}}$ of Eq. (1) is given by $\delta^{\text{SCC}} = \delta^{\text{exp}} - \delta^{\text{nucl}}$. Before confronting the above experimental $\delta$ with Eq. (1) we must cope with $\delta^{\text{nucl}}$.

The most important component of $\delta^{\text{nucl}}$ comes from the binding energy effect referred to in the Introduction; unfortunately, its computation is very delicate and highly sensitive to the details of the procedure adopted and to the parameters of the potentials used to generate the single-nucleon wave functions. This is illustrated in Table III where the earlier computations of Wilkinson[8] in two versions, W-A and W-B, are given together with recent computations of Towner[20] also in two versions T-I and T-II. Both authors used the

TABLE III. Computed contribution of the binding energy effect to $\delta^{\text{nucl}}$ in $A = 8$ and 12 according to Wilkinson (Ref. 8) (W-A and W-B) and to Towner (Ref. 20) (T-I and T-II).

| Decay | W-A | W-B | T-I | T-II |
|---|---|---|---|---|
| $A = 8$ | 0.040 | 0.048 | 0.115 | 0.129 |
| $A = 12$ | 0.146 | 0.098 | 0.144 | 0.171 |
| $A = 12^*$ | 0.063 | 0.048 | 0.134 | 0.093 |

same wave functions [Cohen and Kurath[10] in the (6–16)2BME version for $A = 8$ and in the (8–16) POT version for $A = 12$]. The differences therefore arise entirely from differences of procedure and potentials.[21] [The binding-energy effect on the $(W_0^+ + W_0^-)$ dependence of $\delta_8$ should be small since it is due to the change of binding of the strongly-bound final-state nucleons in $^8Be$ where the nucleon tail effects that give rise to the asymmetry are slight.] Towner[20] has furthermore extensively investigated the sensitivity of the $A = 12$ binding energy effect to the many parameters that enter into his computation; although his numbers quoted in Table III are his "best values" it is clear that substantially different values could result from not-unreasonable changes in the parameters.

It seems that the best present attitude to the binding energy effect is to use as "limits" for each $A$ value the largest and smallest values reported in Table III but recognize that the real figure may be anywhere in between or, indeed, somewhat outside so that we interpret the limits as standard deviation points. If we identify $\delta^{\text{nucl}}$[22] in this way and combine it with the experimental uncertainties where these are significant we gain the following values for $\delta^{\text{SCC}}$:

$$\delta_8^{\text{SCC}} = 0.025 \pm 0.045,$$

$$\delta_{12}^{\text{SCC}} = -0.02 \pm 0.04,$$

$$\delta_{12^*}^{\text{SCC}} = -0.10 \pm 0.11 .$$

These ranges we now combine with Kubodera, Delorme, and Rho's nuclear structure computations[4] of $J$ and $L$ which we insert into Eq. (1) to define, for each datum, a band on a $\zeta$-$\lambda$ plot. This is shown in Fig. 4. From these data we derive, at the 99% confidence limit:

$$|\zeta| < 4 \times 10^{-3} \text{ MeV}^{-1},$$ [23]

$$|\lambda| < 1.5 \times 10^{-2} .$$

The merit of the $A = 12^*$ datum is apparent from Fig. 4: it cuts through the plot at a large angle (owing to a large negative value of $L$) and so is particularly useful in restricting the range of $\lambda$ just as the $A = 8$ slope measurement is powerful

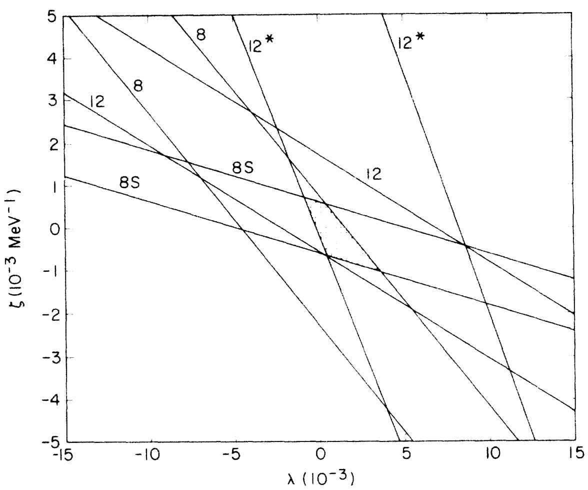


FIG. 4. The $\zeta$-$\lambda$ plot for the four data discussed in the text. The lines are the one standard deviation limits to the allowed $\zeta$-$\lambda$ bands for the various $A$ values that label them (12* means the transitions to the first excited state of $^{12}$C; 8S is the band deriving from the $(W_0^+ + W_0^-)$ dependence of $\delta$ in $A=8$); the cross-hatched area is that common to all four data.

in restricting the range of $\zeta$.

Before continuing our discussion we pause to examine the question of the odd-$A$ members of the $1p$ shell, $A=9$ and 13, for which asymmetry measurements have been made[8]: $\delta_9 = 0.188 \pm 0.030$; $\delta_{13} = 0.166 \pm 0.026$. These asymmetries are large and well-determined; superficially they appear to demand large and significant values of $\delta^{\mathrm{SCC}}$ since the range of binding energy corrections (defined as above for the even-$A$ cases) is only $-0.01$ to 0.06 for $A=9$ and 0.03 to 0.08 for $A=13$. The reason for not using these data at present concerns the assumption that is made about mirror symmetry of the final states. As has been pointed out before[24] the even-$A$ cases (including those in the $2s$, $1d$ shell) show little or no residual $\delta^{\mathrm{SCC}}$ asymmetry after correcting for the binding energy effect in the initial state whereas the odd-$A$ cases (those just quoted plus $A=17$ and 25) still show large asymmetries after such correction. However, in the odd-$A$ cases the *a priori* expectation $\delta=0$ rests on identity of the two final states as well as of the two initial states; but in fact the odd-$A$ final states are far from good mirrors of each other because the nucleon binding energies are low—in most cases negative for the proton in the $T_Z = -\frac{1}{2}$ bodies—and significantly different for the two decays: for the ground state decays the final state neutron (proton) binding energies are 1.67 ($-0.19$) and 4.95 (1.94) MeV for $A=9$ and 13, respectively. The qualitative sense of this final state binding energy effect will be to give a positive $\delta$ as observed since the overlap with the less-tightly-bound or continuum state proton in the $T_Z = -\frac{1}{2}$ final state will be poorer than with the more-tightly-bound mirror neutron in the $T_Z = +\frac{1}{2}$ final state. This final state effect, absent for the even-$A$ cases discussed here, has not yet been reliably estimated but it seems unlikely that it should be less important than the initial state effect and so may well account for the residual asymmetry.

We now return to our discussion of $\zeta$ and $\lambda$. Owing to the extreme delicacy of the binding energy correction it seems most unlikely that our confidence in the limits to $|\zeta|$ and $|\lambda|$ that we have quoted above will improve very much and that we must conclude that limits of this order are as close as can be set on second class currents from the asymmetry phenomenon itself among the nuclei we discuss here. Indeed, when we consider that we have not yet questioned the reliability of the nuclear wave functions on which the estimates of the binding energy correction and the construction of Fig. 4 depend, one may feel that the limits on $\zeta$ and $\lambda$ presented above are too sharp. The sensitivity of our limits to possible changes in the wave functions ought to be checked and in the measure that this has not yet been done our conclusions must remain tentative.

As has been pointed out[4] many of the nuclear structure uncertainties drop away when correlation experiments, as opposed to total decay rate experiments, are performed; it is in this direction that one should go for a sharper probing of the existence of second class currents from the $\beta$ decay of complex nuclei.

It is finally interesting to examine our limits in the light of the model[4] that takes the bare second class nucleon vertex as negligible and that associates the real nucleon terms $g_T$ and $g_T'$ with the nucleon vertex renormalization coming from the $\omega$-meson decay $\omega \to \pi e \nu$; in this case both the nucleon contribution $\zeta$ and the explicit meson exchange current contribution contained in $\lambda$ are due to the $\omega$ meson.[25] Kubodera, Delorme, and Rho,[4] in a simplified calculation, find that $\zeta$ has the divergence-free form:

$$\zeta \approx 1.6 \frac{F_\omega g_{\omega NN} g_r}{16\pi^2 m_\omega^{\,2}} ,$$

where $m_\omega$ is the mass of the $\omega$ meson (784 MeV), $g_{\omega NN}$ the $\omega NN$ coupling constant (6.8), $g_r$ the renormalized $\pi NN$ coupling constant (13.5), and $F_\omega$ the unknown form factor for $\omega \to \pi e \nu$ decay. Our limit on $\zeta$ therefore corresponds to:

$$|F_\omega|_\zeta < 2.8 M ,$$

where $M$ is the nucleon mass. Unfortunately the individual $g_T$, $g_T'$ show logarithmic divergences that prevent the immediate extraction of a limit on $F_\omega$ from $\lambda$:

$$\lambda = \frac{m_\pi{}^3 g_r{}^2}{24\pi M^2}\,\kappa[f(\Lambda) - 16\pi^2/g_r{}^2]\ , \qquad (2)$$

where $f(\Lambda)$ is a cutoff function such that $g_T' = \kappa f(\Lambda)$. However, $f$ is a slowly changing function of $\Lambda$ and for moderate values of a few GeV for the latter we find that our limit on $\lambda$ corresponds to:

$$|F_\omega|_\lambda < 3.5M\ .$$

It therefore seems, on the basis of this simple model that $|F_\omega|$ is likely to be less than about $3M$ to which would correspond a free $\omega$-meson partial lifetime:

$$\tau_{\omega\to\pi e\nu} > 5\times 10^{-11}\ \text{sec}\ .$$

This limit corresponds to a branching ratio for this mode in free $\omega$ decay of less than 1 part in $10^{12}$ (the width of the $\omega$ meson is 10 MeV) which explains our earlier remark that direct experimentation is unlikely to compete with our present result which, although surrounded in many uncertainties, can presumably not be wrong by a large factor.

To set our limit on $\omega\to\pi e\nu$ in perspective we may consider the decay through the vector interaction of a hypothetical pion-like particle $\Pi$, $\Pi\to\pi e\nu$, of mass equal to that of the $\omega$ meson, i.e., assume a tremendous mass splitting between $\pi^\pm$ (viz. $\Pi^\pm$) and $\pi^0$. We have, using standard results: ($R_0$ allows for the fact that, in $\pi$ decay, the mass of the electron is finite although $\pi^0$ recoil can be neglected; $R_2$ allows for the fact that, in $\Pi$ decay, the $\pi^0$ recoil cannot be neglected although the electron mass can.)

$$\frac{\tau_{\pi\to\pi e\nu}}{\tau_{\Pi\to\pi e\nu}} = \frac{15}{32}\left(\frac{m_\Pi}{m_{\pi^+}-m_{\pi^0}}\right)^5 \frac{I(\eta)}{R(\xi)}\ , \qquad (3)$$

where

$$\eta = 1-(m_\pi/m_\Pi)^2,\quad \xi = m_e/(m_{\pi^+}-m_{\pi^0})\ ,$$

$$I(\eta) = -(1-\eta)^2\ln(1-\eta) - \eta + \tfrac{3}{2}\eta^2 - \tfrac{1}{3}\eta^3 - \tfrac{1}{12}\eta^4\ ,$$

$$R(\xi) = (1-\xi^2)^{1/2}(1-\tfrac{9}{2}\xi^2+4\xi^4) + \tfrac{15}{2}\xi^4\ln\left(\frac{1+(1-\xi^2)^{1/2}}{\xi}\right)\ .$$

This factor is about $4.8\times 10^{-9}$ so, with $\tau_{\pi\to\pi e\nu}\approx 2.5$ sec we gain $\tau_{\Pi\to\pi e\nu}\approx 5\times 10^{-10}$ sec. It therefore seems that the intrinsic strength of the second class interaction cannot yet be said to be less than that of the first. We finally note, with Kubodera, Delorme, and Rho,[4] that the $\omega$-meson vertex renormalization model just used, if extended to the mirror hyperon decay $\Sigma\to\Lambda e\nu$ through the use of SU(3), means that our present limit on $F_\omega$ implies an asymmetry in those decays of < 0.25.

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[22]Other contributions to $\delta^{\text{nucl}}$ are due to: (i) the effect of Coulomb and other charge-dependent forces in changing the fractional parentage coefficients as between the two sides of the mirror; (ii) second-forbidden terms

in the axial current; (iii) weak magnetism; (iv) the induced pseudoscalar term; (v) radiative corrections. Towner has computed all these corrections for the cases in point and finds their total to be small in relation to the uncertainty in the binding energy effect. We therefore neglect them here.

[23]Note that earlier analyses (e.g., Refs. 2 and 24) have used natural units in which we should then say $|\zeta| < 2\times10^{-3}$.

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[25]It seems reasonable to single out the $\omega$ meson for this role since it is by far the lightest meson of appropriate quantum numbers that is strongly coupled to the nucleon.