Superallowed $0^+ \rightarrow 0^+$ nuclear β decays: A new survey with precision tests of the conserved vector current hypothesis and the standard model

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A new critical survey is presented of all half-life, decay-energy, and branching-ratio measurements related to 20 superallowed $0^+ \rightarrow 0^+ \beta$ decays. Compared with our last review, there are numerous improvements: First, we have added 27 recently published measurements and eliminated 9 references, either because they have been superseded by much more precise modern results or because there are now reasons to consider them fatally flawed; of particular importance, the new data include a number of high-precision Penning-trap measurements of decay energies. Second, we have used the recently improved isospin symmetry-breaking corrections, which were motivated by these new Penning-trap results. Third, our calculation of the statistical rate function f now accounts for possible excitation in the daughter atom, a small effect but one that merits inclusion at the present level of experimental precision. Finally, we have re-examined the systematic uncertainty associated with the isospin symmetry-breaking corrections by evaluating the radial-overlap correction using Hartree-Fock radial wave functions and comparing the results with our earlier calculations, which used Saxon-Woods wave functions; the provision for systematic uncertainty has been changed as a consequence. The new "corrected" $\mathcal{F}t$ values are impressively constant and their average, when combined with the muon lifetime, yields the up-down quark-mixing element of the Cabibbo-Kobayashi-Maskawa (CKM) matrix, $V_{ud} = 0.97425 \pm 0.00022$. The unitarity test on the top row of the matrix becomes $|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 0.99995 \pm 0.00061$. Both V_{ud} and the unitarity sum have significantly reduced uncertainties compared with our previous survey, although the new value of V_{ud} is statistically consistent with the old one. From these data we also set limits on the possible existence of scalar interactions, right-hand currents, and extra Z bosons. Finally, we discuss the priorities for future theoretical and experimental work with the goal of making the CKM unitarity test even more definitive.

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I. INTRODUCTION

Precise measurements of the β decay between nuclear analog states of spin $J^\pi = 0^+$ and isospin $T = 1$ provide demanding and fundamental tests of the properties of the electroweak interaction. Collectively, these transitions can sensitively probe the conservation of the vector weak current, set tight limits on the presence of scalar or right-hand currents, and, by providing the most precise value for V_{ud}, the up-down quark-mixing element of the Cabibbo-Kobayashi-Maskawa (CKM) matrix, can contribute to the most demanding available test of the unitarity of that matrix, a property fundamental to the electroweak standard model.

We have published five previous surveys of $0^+ \rightarrow 0^+$ superallowed transitions [1–5], the first having appeared 35 years ago and the most recent four years ago. In each, we published a complete survey of all relevant nuclear data that pertained to these superallowed transitions and used the results to set limits on the weak-interaction parameters that were important at the time. A particularly notable outcome of our analysis four years ago [5] was that the sum of squares of the top-row elements of the CKM matrix—the test of CKM unitarity—remained ambiguous, with the possibility of a significant shortfall in the unitarity sum.

Since our last survey closed in November 2004, there has been a great deal of activity in this field prompted at least in part by the tantalizing possibility that new physics could be revealed by a failure in CKM unitarity. New measurements relating to $0^+ \rightarrow 0^+$ superallowed transitions have appeared in 27 publications, an addition of 20% to the papers accumulated up to 2004. Many of these measurements were of unprecedented precision so they did not merely add more of the same: They palpably improved the results, in some cases by tightening their error bars and, in others, by changing their central values. Penning-trap measurements of decay energies, which only became possible after 2004, have been especially effective in this regard.

In addition to new measurements, there have also been important improvements to the small theoretical corrections that must be applied to the data to extract V_{ud} and the values of other weak-interaction parameters. In the past four years, the radiative [6] and isospin symmetry-breaking corrections [7] have both been subjected to major re-evaluations, which have undoubtedly improved their values and, in the former case, has reduced the uncertainty by a factor of 2.

In parallel with these developments, there has also been considerable activity in the determination of V_{us}, the other matrix element that plays a role in the top-row unitarity test of the CKM matrix. (The third element in the top row, V_{ub}, is very small and contributes a negligible 0.001% to the unitarity sum.) As with the work related to V_{ud}, this activity has encompassed new experiments—precise measurements of kaon branching

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ratios—as well as improved theoretical corrections. However, in contrast with V_{ud}, not only the uncertainty of V_{us} but also its central value have been considerably changed by this recent work (see Ref. [8] for an up-to-date overview of V_{us}).

Overall, the recent improvements have been numerous enough and their impact on the unitarity test significant enough that this is an opportune time to produce a new and updated survey of the nuclear data used to establish V_{ud}. We incorporate data on a total of 20 superallowed transitions and have continued the practice we began in 1984 [3] of updating all original data to take account of the most modern calibration standards. In addition to including the improved correction terms already referred to, we have also upgraded our calculation of the statistical rate function f to include provision for excitation of the daughter atom, and we have included a more extensive treatment of possible systematic uncertainties associated with the isospin symmetry-breaking corrections.

Superallowed $0^+ \to 0^+$ β decay between $T = 1$ analog states depends uniquely on the vector part of the weak interaction and, according to the conserved vector current (CVC) hypothesis, its experimental ft value should be related to the vector coupling constant, a fundamental constant that is the same for all such transitions. In practice, the expression for ft includes several small ($\sim$1%) correction terms. It is convenient to combine some of these terms with the ft value and define a "corrected" $\mathcal{F}t$ value. Thus, we write [5]

$$\mathcal{F}t \equiv ft(1+\delta'_R)(1+\delta_{\rm NS}-\delta_C) = \frac{K}{2G_V^2\left(1+\Delta_R^V\right)}, \quad (1)$$

where $K/(\hbar c)^6 = 2\pi^3\hbar \ln 2/(m_e c^2)^5 = 8120.2787(11) \times 10^{-10}$ GeV^{-4} s, G_V is the vector coupling constant for semileptonic weak interactions, δ_C is the isospin-symmetry-breaking correction, and Δ_R^V is the transition-independent part of the radiative correction. The terms δ'_R and $\delta_{\rm NS}$ comprise the transition-dependent part of the radiative correction, the former being a function only of the electron's energy and the Z of the daughter nucleus, while the latter, like δ_C, depends in its evaluation on the details of nuclear structure. From this equation, it can be seen that each measured transition establishes an individual value for G_V and, if the CVC assertion is correct that G_V is not renormalized in the nuclear medium, all such values—and all the $\mathcal{F}t$ values themselves—should be identical within uncertainties, regardless of the specific nuclei involved.

Our procedure in this paper is to examine all experimental data related to 20 superallowed transitions, comprising all those that have been well studied, together with others that are now coming under scrutiny after becoming accessible to precision measurement. The methods used in data evaluation are presented in Sec. II. The calculations and corrections required to extract $\mathcal{F}t$ values from these data are described and applied in Sec. III; in the same section, we use the resulting $\mathcal{F}t$ values to test CVC. Finally, in Sec. IV we explore the impact of these results on a number of weak-interaction issues: CKM unitarity as well as the possible existence of scalar interactions, right-hand currents, and extra Z bosons. This is much the same pattern as we followed in our last review [5] so we will not describe the formalism again in detail, referring the reader instead to that earlier work.

II. EXPERIMENTAL DATA

The ft value that characterizes any β transition depends on three measured quantities: the total transition energy, $Q_{\rm EC}$, the half-life, $t_{1/2}$, of the parent state, and the branching ratio, R, for the particular transition of interest. The $Q_{\rm EC}$ value is required to determine the statistical rate function, f, and the half-life and branching ratio combine to yield the partial half-life, t. In Tables I–VII we present the measured values of these three quantities and supporting information for a total of 20 superallowed transitions.

A. Evaluation principles

In our treatment of the data, we considered all measurements formally published before September 2008 and those we knew to be in an advanced state of preparation for publication by that date. We scrutinized all the original experimental reports in detail. Where necessary and possible, we used the information provided there to correct the results for calibration data that have improved since the measurement was made. If corrections were evidently required but insufficient information was provided to make them, the results were rejected. Of the surviving results, only those with (updated) uncertainties that are within a factor of 10 of the most precise measurement for each quantity were retained for averaging in the tables. Each datum appearing in the tables is attributed to its original journal reference via an alphanumeric code comprising the initial two letters of the first author's name and the two last digits of the publication date. These codes are correlated with the actual reference numbers, [9–173], in Table VIII.

The statistical procedures we have followed in analyzing the tabulated data are based on those used by the Particle Data Group in their periodic reviews of particle properties (e.g., Ref. [174]) and adopted by us in our previous surveys. We gave a detailed description of those procedures in our 2004 survey [5] so will not repeat it here.

Our evaluation principles and associated statistical procedures constitute a very conservative approach to the data. Unless there is a clearly identifiable reason to reject a result, we include it in our data base even if it deviates significantly from other measurements of the same quantity, the consequent nonstatistical spread in results being reflected in an increased uncertainty assigned to the average. Wherever this occurs, the factor by which the uncertainty has been increased is listed in the "scale" column of a table. There are a few examples, though, of a single publication that includes a number of measurements—a set of half-lives or $Q_{\rm EC}$ values for example—most or all of which deviate substantially from other accepted measurements of the same quantities. In such cases, we consider that some systematic problem has been revealed, and we exclude all the results from that publication. If any measurement with an acceptable uncertainty is nevertheless

TABLE I. Measured results from which the decay transition energies, Q_{EC}, have been derived for superallowed β decays. The lines giving the average superallowed Q_{EC} values themselves are in bold print. (See Table VIII for the correlation between the alphanumeric reference code used in this table and the actual reference numbers.)

Parent/daughter nuclei		Property[a]	Measured energies used to determine Q_{EC} (keV)			Average value	
			1	2	3	Energy (keV)	Scale
$T_z = -1$							
^{10}C	^{10}B	Q_{EC}(gs)	3647.84 ± 0.34 [Ba84]	3647.95 ± 0.12 [Ba98]		3647.94 ± 0.11	1.0
		E_x(d0$^+$)	1740.15 ± 0.17 [Aj88]	1740.07 ± 0.02[b]		1740.07 ± 0.02	1.0
		Q_{EC}(sa)				**1907.87 ± 0.11**	
^{14}O	^{14}N	Q_{EC}(gs)	5143.35 ± 0.60 [Bu61]	5145.09 ± 0.46 [Ba62]	5145.57 ± 0.48 [Ro70]		
			5143.43 ± 0.37 [Wh77]	5144.34 ± 0.17 [To03]		5144.33 ± 0.29	2.1
		E_x(d0$^+$)	2312.798 ± 0.011 [Aj91]			2312.798 ± 0.011	
		Q_{EC}(sa)				**2831.24 ± 0.23**[c]	**2.3**
^{18}Ne	^{18}F	ME(p)	5316.8 ± 1.5 [Ma94]	5317.63 ± 0.36 [Bl04b]		5317.58 ± 0.35	1.0
		ME(d)	873.31 ± 0.94 [Bo64]	875.5 ± 2.2 [Ho64]	876.5 ± 2.8 [Pr67]		
			877.2 ± 3.0 [Se73]	873.96 ± 0.61 [Ro75]		874.02 ± 0.48	1.0
		Q_{EC}(gs)	4438 ± 9 [Fr63]			4443.54 ± 0.60	1.0
		E_x(d0$^+$)	1041.55 ± 0.08 [Ti95]			1041.55 ± 0.08	
		Q_{EC}(sa)				**3401.99 ± 0.60**	
^{22}Mg	^{22}Na	ME(p)	−401.2 ± 3.0 [Ha74c]	−400.4 ± 1.3[d]	−400.5 ± 1.0 [Pa05]	−400.5 ± 0.8	1.0
		ME(d)	−5184.3 ± 1.5 [We68]	−5182.5 ± 0.5 [Be68]	−5181.3 ± 1.7 [An70]		
			−5183.2 ± 1.0 [Gi72]	−5181.56 ± 0.16 [Mu04]	−5181.08 ± 0.30 [Sa04]	−5181.58 ± 0.19	1.7
		Q_{EC}(gs)	4781.64 ± 0.28 [Mu04]	4781.40 ± 0.67 [Sa04]		4781.55 ± 0.25	1.0
		E_x(d0$^+$)	657.00 ± 0.14 [En98]			657.00 ± 0.14	
		Q_{EC}(sa)				**4124.55 ± 0.28**	
^{26}Si	^{26}Al	ME(p)	−7145.4 ± 3.0 [Ha74c]	−7139.5 ± 1.0 [Pa05]		−7140.1 ± 1.8	1.9
		ME(d0$^+$)	−11981.99 ± 0.26[e]			−11981.99 ± 0.26	
		Q_{EC}(sa)	4850 ± 13 [Fr63]			**4842.0 ± 1.8**	**1.0**
^{30}S	^{30}P	ME(p)	−14060 ± 15 [Mi67]	−14054 ± 25 [Mc67]	−14068 ± 30 [Ha68]		
			−14063.4 ± 3.0 [Ha74c]			−14063.1 ± 2.9	1.0
		ME(d)	−20203 ± 3 [Ha67]	−20200.58 ± 0.40 [Re85]		−20200.62 ± 0.40	1.0
		Q_{EC}(gs)				6137.5 ± 2.9	
		E_x(d0$^+$)	677.29 ± 0.07 [En98]			677.29 ± 0.07	
		Q_{EC}(sa)	5437 ± 17 [Fr63]			**5459.5 ± 3.9**	**1.3**
^{34}Ar	^{34}Cl	ME(p)	−18380.2 ± 3.0 [Ha74c]	−18378.4 ± 3.5 [He01]	−18377.10 ± 0.41 [He02]	−18377.17 ± 0.40	1.0
		ME(d)	−24440.15 ± 0.26[e]			−24440.15 ± 0.26	
		Q_{EC}(sa)				**6062.98 ± 0.48**	
^{38}Ca	^{38}K	ME(p)	−22058.53 ± 0.28 [Ri07]	−22058.01 ± 0.65 [Ge07]		−22058.45 ± 0.26	1.0
		ME(d0$^+$)	−28670.20 ± 0.32[e]			−28670.20 ± 0.32	
		Q_{EC}(sa)				**6611.75 ± 0.41**	
^{42}Ti	^{42}Sc	ME(p)	−25121 ± 6 [Mi67]	−25086 ± 30 [Ha68]	−25124 ± 13 [Zi72]	−25120.7 ± 5.3	1.0
		ME(d)	−32121.12 ± 0.29[e]			−32121.12 ± 0.29	
		Q_{EC}(sa)				**7000.5 ± 5.4**	
$T_z = 0$							
^{26}Al^m	^{26}Mg	Q_{EC}(gs)	4004.79 ± 0.55 [De69]	4004.41 ± 0.10[f]	4004.40 ± 0.22 [Ge08]	4004.42 ± 0.09	1.0
		E_x(p0$^+$)	228.305 ± 0.013 [En98]			228.305 ± 0.013	
		Q_{EC}(sa)	4232.19 ± 0.12 [Br94]	4232.83 ± 0.13 [Er06b]		**4232.66 ± 0.12**[c]	**2.1**
^{34}Cl	^{34}S	**Q_{EC}(sa)**	5490.3 ± 1.9 [Ry73a]	5491.6 ± 2.3 [Ha74d]	5491.71 ± 0.54 [Ba77c]		
			5491.65 ± 0.26[g]			**5491.64 ± 0.23**	**1.0**
^{38}K^m	^{38}Ar	Q_{EC}(gs)	5914.82 ± 0.61 [Ja78]			5914.82 ± 0.61	
		E_x(p0$^+$)	130.4 ± 0.3 [En98]			130.4 ± 0.3	
		Q_{EC}(sa)	6044.6 ± 1.5 [Bu79]	6044.38 ± 0.12 [Ha98]		**6044.40 ± 0.11**	**1.0**
^{42}Sc	^{42}Ca	**Q_{EC}(sa)**	6425.84 ± 0.17[h]	6426.13 ± 0.21 [Er06b]		**6426.28 ± 0.30**[c]	**3.0**
^{46}V	^{46}Ti	**Q_{EC}(sa)**	7053.3 ± 1.8 [Sq76]	7052.90 ± 0.40 [Sa05]	7052.72 ± 0.31 [Er06b]		
			7052.11 ± 0.27 [Fa09]			**7052.49 ± 0.16**	**1.3**
^{50}Mn	^{50}Cr	**Q_{EC}(sa)**	7634.48 ± 0.07 [Er08]			**7634.45 ± 0.07**[c]	**1.0**
^{54}Co	^{54}Fe	**Q_{EC}(sa)**	8244.54 ± 0.10 [Er08]			**8244.37 ± 0.28**[c]	**3.4**
^{62}Ga	^{62}Zn	**Q_{EC}(sa)**	9181.07 ± 0.54 [Er06a]			**9181.07 ± 0.54**	

TABLE I. (*Continued.*)

Parent/daughter nuclei		Property[a]	Measured energies used to determine Q_{EC} (keV)			Average value	
			1	2	3	Energy (keV)	Scale
^{66}As	^{66}Ge	ME(p)	−52018 ± 30 [Sc07]			52018 ± 30	
		ME(d)	−61607.0 ± 2.4 [Sc07]			−61607.0 ± 2.4	
		Q_{EC}(sa)	9550 ± 50 [Da80]			**9579 ± 26**	**1.0**
^{70}Br	^{70}Se	**Q_{EC}(sa)**	9970 ± 170 [Da80]			**9970 ± 170**	
^{74}Rb	^{74}Kr	ME(p)	−51905 ± 18 [He02]	−51915.2 ± 4.0 [Ke07]		−51914.7 ± 3.9	1.0
		ME(d)	−62332.0 ± 2.1 [Ro06]			−62332.0 ± 2.1	
		Q_{EC}(sa)				**10417.3 ± 4.4**	

[a]Abbreviations: gs, transition between ground states; sa, superallowed transition; p, parent; d, daughter; ME, mass excess; $E_x(0^+)$, excitation energy of the 0^+ (analog) state. Thus, for example, Q_{EC}(sa) signifies the Q_{EC} value for the superallowed transition, ME(d), the mass excess of the daughter nucleus; and $ME(d0^+)$, the mass excess of the daughter's 0^+ state.
[b]Result based on Refs. [Ba88] and [Ba89].
[c]Average result includes the results of Q_{EC} pairs; see Table II.
[d]Result based on Refs. [Bi03], [Se05], and [Je07].
[e]Result obtained from data elsewhere in this table.
[f]Result based on Refs. [Is80], [Al82], [Hu82], [Be85], [Pr90], [Ki91], and [Wa92].
[g]Result based on Refs. [Wa83], [Ra83], and [Li94].
[h]Result based on Refs. [Zi87] and [Ki89].

excluded from our data base, the reason for its exclusion is listed in Table VII.

One particularly significant, longstanding reference had to be excluded for the first time from this survey. Our decision to do so deserves a more detailed explanation. In 1977, Vonach *et al.* published in a single paper [162] the Q_{EC} values for seven superallowed emitters (^{14}O, ^{26}Al^m, ^{34}Cl, ^{42}Sc, ^{46}V, ^{50}Mn, and ^{54}Co), which they had determined from the Q values for (^{3}He,t) reactions on their stable daughters. They had used a "precision time-of-flight measuring system" with the Q3D spectrograph of the Munich Tandem Laboratory to produce uncertainties of 0.4–0.6 keV. For the time, these were very precise results and consequently they had a major impact on the superallowed data base for the following three decades.

The first indication that the Vonach *et al.* results might have a problem came with the first Penning-trap measurement of a superallowed Q_{EC} value [149]. The new measurement for ^{46}V quoted 0.4-keV uncertainty and differed from the old result by 2.4 keV, four of Vonach's claimed standard deviations. Within a year, a second Penning-trap measurement [61] had confirmed the new ^{46}V result and had also found that the ^{42}Sc Q_{EC} value differed from the Vonach result by six times the latter's quoted uncertainty. Two years later, another Penning-trap measurement [62] indicated that the ^{50}Mn and ^{54}Co Q_{EC} values also differed from Vonach's results by a similar amount. These most-recent Penning-trap results quoted 0.1-keV uncertainties. A current overview of the situation for all seven of the superallowed transitions measured by Vonach appears in Fig. 1, where each Vonach result is compared with the equivalent result(s) from a Penning trap, and both are compared to the average of all results for the same transition.

TABLE II. Q_{EC} value differences for superallowed β-decay branches. These data are also used as input to determine some of the average Q_{EC} values listed in Table I. (See Table VIII for the correlation between the alphanumeric reference code used in this table and the actual reference numbers.)

Parent nucleus 1	Parent nucleus 2	$Q_{EC2} - Q_{EC1}$ (keV)	
		Measurement	Average[a]
^{14}O	^{26}Al^m	1401.68 ± 0.13 [Ko87]	1401.43 ± 0.26
^{26}Al^m	^{42}Sc	2193.5 ± 0.2 [Ko87]	2193.62 ± 0.32
^{42}Sc	^{50}Mn	1207.6 ± 2.3 [Ha74d]	1208.17 ± 0.30
^{42}Sc	^{54}Co	1817.2 ± 0.2 [Ko87]	1818.09 ± 0.41
^{50}Mn	^{54}Co	610.09 ± 0.17 [Ko87] [Ko97b]	609.92 ± 0.29

[a]Average values include the results of direct Q_{EC} value measurements; see Table I.

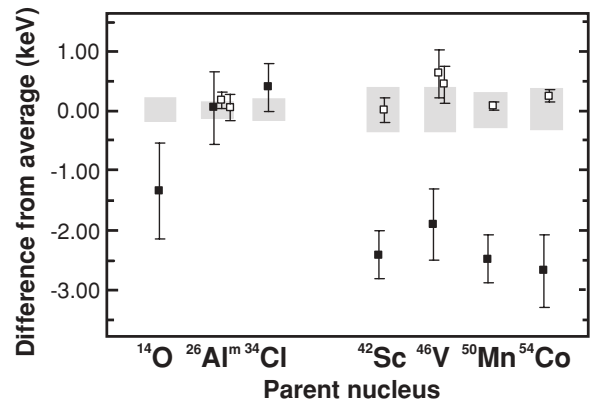

FIG. 1. Differences between individual measurements and the averages of all measurements for the seven parent nuclei studied by Vonach *et al.* [162]. The filled squares are the results of the (^{3}He, t) measurements of Vonach *et al.*; the open squares are recent Penning-trap results [61,62,73,149]. For each parent nucleus, the gray band about the zero line represents the uncertainty of the average for that case. Note that all the averages include the results of Vonach *et al.*, the Penning-trap results, and any other relevant measurements appearing in Table I.

With only two transitions, those from ^{26}Al^m and ^{34}Cl, showing agreement and the four cases already mentioned displaying serious disagreement, we believed that the best approach was to eliminate all the results published in the original Vonach *et al.* reference [162]. This conclusion has been further supported by a very recent (^{3}He,t) measurement of the ^{46}V Q_{EC} value [63] made with much of the same experimental equipment originally used by Vonach *et al.* 30 years ago. The new result disagrees with the old measurement and confirms the new Penning-trap values.

Before Penning traps could be applied to these measurements, all superallowed Q_{EC} values were determined via nuclear reactions. In addition to those employing (^{3}He,t) reactions, two other types of experiment led to rather precise results: the measurement of (p,n) thresholds and the combined measurements of (p,γ) and (n,γ) Q values on the same target, one reaction leading to the superallowed parent and the other to the daughter. We are now in a position to compare the different types of measurement to examine whether there are any systematic differences among them. A careful study of this issue [175], restricted to the region around $A = 26$, was undertaken several years ago and revealed no evidence of any systematic differences. We can now confirm this conclusion over a wider mass range with the help of Fig. 2. In that figure we consider nine superallowed transitions, which we will refer to as the "traditional nine" cases. They are the only superallowed transitions that populate a stable daughter nucleus and, for obvious reasons, were the only ones whose Q values could be measured to high precision in the pre-trap era. There are no systematic deviations apparent in the figure, leading us to conclude that, whatever problems plagued the measurements of Vonach *et al.* [162], they were associated with that particular experiment and were not endemic to a whole class of experiments. Of course, this conclusion could be strengthened by new Penning-trap data for ^{10}C, ^{14}O, ^{34}Cl, and ^{38}K^m.

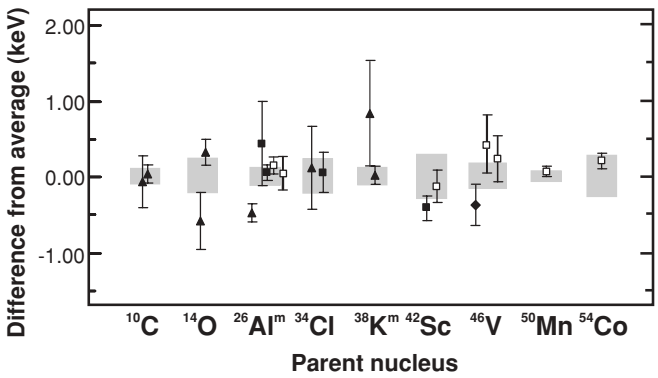

FIG. 2. Differences between individual measurements and average values for the "traditional nine" transitions; the results of Vonach *et al.* [162] have been removed from the averages. The open squares are the results of Penning-trap measurements [61,62,73,149]; the filled squares are from combined (p,γ) and (n,γ) measurements (see Ref. [55] and the references listed in Footnotes f, g, and h of Table I); the filled triangles are from (p,n) threshold measurements [26,27,31,45,87,106,160,166]; and the filled diamond is the new (^{3}He,t) measurement for ^{46}V [63]. For each parent nucleus, the gray band about the zero line represents the uncertainty of the average for that case.

B. Data tables

The Q_{EC} value data appear in Tables I and II. For the "traditional nine" superallowed decays—those of ^{10}C, ^{14}O, ^{26}Al^m, ^{34}Cl, ^{38}K^m, ^{42}Sc, ^{46}V, ^{50}Mn, and ^{54}Co—with stable daughter nuclei, their Q_{EC} values were all determined in the past by direct reaction measurements of that property. More recently, a growing number of Penning-trap measurements, also extending to nuclei outside of the traditional nine, determine the parent and daughter masses in a single experiment, thus effectively measuring the Q_{EC} value directly. Measurements of both types are identified in column 3 of Table I by "Q_{EC}(sa)" and each individual result is itemized with its appropriate reference in the next three columns. The weighted average of all measurements for a particular decay appears in column 7, with the corresponding scale factor (see Sec. II A) in column 8. A few of these cases, such as ^{34}Cl and all the cases from ^{42}Sc to ^{62}Ga, have no further complications. There are other cases, however, in which Q_{EC} value differences have been measured in addition to the individual Q_{EC} values. These measurements are presented in Table II. They have been dealt with in combination with the direct Q_{EC} value measurements, as described in Ref. [5], with the final average Q_{EC} value appearing in column 7 of Table I and the average difference in column 4 of Table II. Both are flagged with footnotes to indicate the interconnection.

There are two cases, ^{26}Al^m and ^{38}K^m, in which the superallowed decay originates from an isomeric state. For both, there are Q_{EC} value measurements that correspond to the ground state as well as to the isomer. Obviously, these two sets of measurements are simply related to one another by the excitation energy of the isomeric state in the parent. In Table I the set of measurements for the ground-state Q_{EC} value and for the excitation energy of the isomeric state appear in separate rows, each with its identifying property given in column 3 and its weighted average appearing in column 7. In the row below, the average value given in column 7 for the superallowed transition is the weighted average not only of the direct superallowed Q_{EC} value measurements in that row but also of the result derived from the two preceding rows. Note that in all cases the Q_{EC} value for the superallowed transition appears in bold-face type.

For some decays that lead to radioactive daughter nuclei, there is no direct measurement of the Q_{EC} value for the superallowed transition or the one that exists is rather imprecise. In these cases the Q_{EC} value must depend on the measured mass excesses of the parent and daughter nuclei, together with the excitation energy of the analog 0^+ state in the daughter. Each of these properties is identified in column 3 of Table I, with the individual measurements of that property, their weighted average, and a scale factor appearing in columns to the right. The average Q_{EC} value listed for the corresponding superallowed transition is obtained from these separate averages. If a direct measurement of the superallowed Q_{EC} value exists, then it is also included in the final average.

As in our previous survey [5], we have not used the 2003 mass tables [19] to derive the Q_{EC} values of interest. Our

TABLE III. Half-lives, $t_{1/2}$, of superallowed β emitters. (See Table VIII for the correlation between the alphanumeric reference code used in this table and the actual reference numbers.)

Parent nucleus	Measured half-lives, $t_{1/2}$ (ms)				Average value	
	1	2	3	4	$t_{1/2}$ (ms)	Scale
$T_z = -1$						
^{10}C	19280 ± 20 [Az74]	19295 ± 15 [Ba90]	19310 ± 4 [Ia08]		19308.0 ± 3.8	1.0
^{14}O	70480 ± 150 [Al72]	70588 ± 28 [Cl73]	70430 ± 180 [Az74]	70684 ± 77 [Be78]		
	70613 ± 25 [Wi78]	70560 ± 49 [Ga01]	70641 ± 20 [Ba04]	70696 ± 52 [Bu06]	70620 ± 15	1.2
^{18}Ne	1669 ± 4 [Al75]	1687 ± 9 [Ha75]	1665.6 ± 1.9 [Gr07]		1667.0 ± 1.7	1.0
^{22}Mg	3857 ± 9 [Ha75]	3875.5 ± 1.2 [Ha03]			3875.2 ± 2.4	2.0
^{26}Si	2210 ± 21 [Ha75]	2240 ± 10 [Wi80]	2228.3 ± 2.7 [Ma08]		2228.8 ± 2.9	1.1
^{30}S	1180 ± 40 [Ba67]	1220 ± 30 [Mo71]	1178.3 ± 4.8 [Wi80]		1179.4 ± 4.7	1.0
^{34}Ar	844.5 ± 3.4 [Ha74a]	843.8 ± 0.4 [Ia06]			843.8 ± 0.4	1.0
^{38}Ca	470 ± 20 [Ka68]	439 ± 12 [Ga69]	450 ± 70 [Zi72]	430 ± 12 [Wi80]	440.0 ± 7.8	1.2
^{42}Ti	200 ± 20 [Ni69]	202 ± 5 [Ga69]	173 ± 14 [Al69]		198.8 ± 6.3	1.4
$T_z = 0$						
^{26}Al^m	6346 ± 5 [Fr69a]	6346 ± 5 [Az75]	6339.5 ± 4.5 [Al77]	6346.2 ± 2.6 [Ko83]	6345.0 ± 1.9	1.0
	6345 ± 14 [Sc05]					
^{34}Cl	1526 ± 2 [Ry73a]	1525.2 ± 1.1 [Wi76]	1527.7 ± 2.2 [Ko83]	1526.8 ± 0.5 [Ia06]	1526.55 ± 0.44	1.0
^{38}K^m	925.6 ± 0.7 [Sq75]	922.3 ± 1.1 [Wi76]	921.71 ± 0.65 [Wi78]	924.15 ± 0.31 [Ko83]		
	924.4 ± 0.6 [Ba00]	924.46 ± 0.14 [Ba08]			924.33 ± 0.27	2.3
^{42}Sc	680.98 ± 0.62 [Wi76]	680.67 ± 0.28 [Ko97a]			680.72 ± 0.26	1.0
^{46}V	422.47 ± 0.39 [Al77]	422.28 ± 0.23 [Ba77a]	422.57 ± 0.13 [Ko97a]		422.50 ± 0.11	1.0
^{50}Mn	284.0 ± 0.4 [Ha74b]	282.8 ± 0.3 [Fr75]	282.72 ± 0.26 [Wi76]	283.29 ± 0.08 [Ko97a]		
	283.10 ± 0.14 [Ba06]				283.21 ± 0.11	1.7
^{54}Co	193.4 ± 0.4 [Ha74b]	193.0 ± 0.3 [Ho74]	193.28 ± 0.18 [Al77]	193.28 ± 0.07 [Ko97a]	193.271 ± 0.063	1.0
^{62}Ga	115.84 ± 0.25 [Hy03]	116.19 ± 0.04 [Bl04a]	116.09 ± 0.17 [Ca05]	116.01 ± 0.19 [Hy05]		
	116.100 ± 0.025 [Gr08]				116.121 ± 0.040	1.9
^{66}As	95.78 ± 0.39 [Al78]	95.77 ± 0.28 [Bu88]	97 ± 2 [Ji02]		95.79 ± 0.23	1.0
^{70}Br	80.2 ± 0.8 [Al78]	78.54 ± 0.59 [Bu88]			79.12 ± 0.79	1.7
^{74}Rb	64.90 ± 0.09 [Oi01]	64.761 ± 0.031 [Ba01]			64.776 ± 0.043	1.5

approach is to include all pertinent measurements for each property; typically, only a subset of the available data is included as input to the mass tables. Furthermore, we have examined each reference in detail and either accepted the result, updated it to modern calibration standards, or rejected it for cause. The updating procedures are outlined, reference by reference, in Table VI and the rejected results are similarly documented in Table VII. With a comparatively small data set, we could afford to pay the kind of individual attention that is impossible when one is considering all nuclear masses.

The half-life data appear in Table III in similar format to Table I. For obvious reasons, half-life measurements do not lend themselves to being updated. Consequently, a number of mostly pre-1970 measurements have been rejected because they were not analyzed with the "maximum-likelihood" method. The importance of using this technique for precision measurements was not recognized until that time [67] and, without access to the primary data, there is no way a new analysis can be applied retroactively. All rejected half-life measurements are also documented in Table VII.

Finally, the branching-ratio measurements are presented in Table IV. The decays of the $T_z = 0$ parents are the most straightforward since, in all these cases, the superallowed branch accounts for >99.5% of the total decay strength. Thus, even imprecise measurements of the weak nonsuperallowed branches can be subtracted from 100% to yield the superallowed branching ratio with good precision. For the higher Z parents of this type, particularly ^{62}Ga and heavier, it has been shown theoretically [88] and experimentally [64,133] that numerous very-weak Gamow-Teller transitions occur, which, in total, can carry significant decay strength. Where such unobserved transitions are expected to exist and have not already been accounted for in the quoted references, we have used a combination of experiment and theory to account for the unobserved strength, with uncertainties being adjusted accordingly.

The branching ratios for decays from $T_z = -1$ parents are much more challenging to determine, since the superallowed branch is usually one of several strong branches—with the notable exception of ^{14}O—and, in two of the measured cases, it actually has a branching ratio of less than 10%. The decays of ^{18}Ne, ^{26}Si, ^{30}S, ^{34}Ar, and ^{42}Ti thus required special treatment. In each case, the absolute branching ratio for a single β transition has been measured. The branching ratios for other β transitions then had to be determined from the relative intensities of β-delayed γ rays in the daughter. The relevant γ-ray intensity measurements appear in Table V, with their averages then being used to determine the superallowed branching-ratio averages shown in bold type in Table IV. These cases are also flagged with a table footnote.

TABLE IV. Measured results from which the branching ratios, R, have been derived for superallowed β transitions. The lines giving the average superallowed branching ratios themselves are in bold print. (See Table VIII for the correlation between the alphanumeric reference code used in this table and the actual reference numbers.)

Parent/daughter nuclei		Daughter state E_x (MeV)	Measured branching ratio, R (%)		Average value	
			1	2	R (%)	Scale
$T_z = -1$						
^{10}C	^{10}B	2.16	$0^{+0.0008}_{-0}$ [Go72]		$0^{+0.0008}_{-0}$	
		1.74	1.468 ± 0.014 [Ro72]	1.473 ± 0.007 [Na91]		
			1.465 ± 0.009 [Kr91]	1.4625 ± 0.0025 [Sa95]		
			1.4665 ± 0.0038 [Fu99]		**1.4646 ± 0.0019**	**1.0**
^{14}O	^{14}N	gs	0.68 ± 0.10 [Sh55, To05]	0.74 ± 0.05 [Fr63, To05]]		
			0.54 ± 0.02 [Si66, To05]		0.571 ± 0.068	3.7
		3.95	0.062 ± 0.007 [Ka69]	0.058 ± 0.004 [Wi80]		
			0.053 ± 0.002 [He81]		0.0545 ± 0.0019	1.1
		2.31			**99.374 ± 0.068**	
^{18}Ne	^{18}F	**1.04**	9 ± 3 [Fr63]	7.70 ± 0.21[a] [Ha75]	**7.70 ± 0.21**	**1.0**
^{22}Mg	^{22}Na	**0.66**	54.0 ± 1.1 [Ha75]	53.15 ± 0.12 [Ha03]	**53.16 ± 0.12**	**1.0**
^{26}Si	^{26}Al	1.06	21.8 ± 0.8 [Ha75]	21.21 ± 0.64 [Ma08]	21.4 ± 0.5	1.0
		0.23			**75.49 ± 0.57**[a]	
^{30}S	^{30}P	gs	20 ± 1 [Fr63]		20 ± 1	
		0.68			**77.4 ± 1.0**[a]	
^{34}Ar	^{34}Cl	0.67	2.49 ± 0.10 [Ha74a]		2.49 ± 0.10	
		gs			**94.45 ± 0.25**[a]	
^{42}Ti	^{42}Sc	0.61	56 ± 14 [Al69]		56 ± 14	
		gs			**43 ± 14**[a]	
$T_z = 0$						
^{26}Al^m	^{26}Mg	**gs**	> 99.997 [Ki91]		$\mathbf{100.000^{+0}_{-0.003}}$	
^{34}Cl	^{34}S	**gs**	> 99.988 [Dr75]		$\mathbf{100.000^{+0}_{-0.012}}$	
^{38}K^m	^{38}Ar	3.38	< 0.0019 [Ha94]	< 0.0008 [Le08]	$0^{+0.0008}_{-0}$	
		gs(^{38}K)	0.0330 ± 0.0043 [Le08]		0.0330 ± 0.0043	
		gs			$\mathbf{99.9670^{+0.0043}_{-0.0044}}$	
^{42}Sc	^{42}Ca	1.84	0.0063 ± 0.0026 [In77]	0.0022 ± 0.0017 [De78]		
			0.0103 ± 0.0031 [Sa80]	0.0070 ± 0.0012 [Da85]	0.0059 ± 0.0014	1.6
		gs			**99.9941 ± 0.0014**	
^{46}V	^{46}Ti	2.61	0.0039 ± 0.0004 [Ha94]		0.0039 ± 0.0004	
		4.32	0.0113 ± 0.0012 [Ha94]		0.0113 ± 0.0012	
		ΣGT[b]	< 0.004		$0^{+0.004}_{-0}$	
		gs			$\mathbf{99.9848^{+0.0013}_{-0.0042}}$	
^{50}Mn	^{50}Cr	3.63	0.057 ± 0.003 [Ha94]		0.057 ± 0.003	
		3.85	< 0.0003 [Ha94]		$0^{+0.0003}_{-0}$	
		5.00	0.0007 ± 0.0001 [Ha94]		0.0007 ± 0.0001	
		gs			**99.9423 ± 0.0030**	
^{54}Co	^{54}Fe	2.56	0.0045 ± 0.0006 [Ha94]		0.0045 ± 0.0006	
		ΣGT[b]	< 0.03		$0^{+0.03}_{-0}$	
		gs			$\mathbf{99.9955^{+0.0006}_{-0.0300}}$	
^{62}Ga	^{62}Zn	ΣGT[b]	0.142 ± 0.008 [Fi08]	0.107 ± 0.024 [Be08]	0.139 ± 0.011	1.4
		gs			**99.862 ± 0.011**	
^{74}Rb	^{74}Kr	ΣGT[b]	0.50 ± 0.10 [Pi03]		0.50 ± 0.10	
		gs			**99.50 ± 0.10**	

[a]Result also incorporates data from Table V.

[b]Designates total Gamow-Teller transitions to levels not explicitly listed; values were derived with the help of calculations in Ref. [88].

III. THE $\mathcal{F}t$ VALUES

With the input data now settled, we can proceed to derive the ft values for the 20 superallowed transitions included in the tables. In our last survey [5], we described and used a new computer code for calculating the statistical rate function f, which surpassed the precision then being obtained with measurements of Q_{EC}. Since then, with the advent of Penning-trap mass measurements, experimental uncertainties

TABLE V. Relative intensities of β-delayed γ rays in the superallowed β-decay daughters. These data are used to determine some of the branching ratios presented in Table IV. (See Table VIII for the correlation between the alphanumeric reference code used in this table and the actual reference numbers.)

Parent/daughter nuclei		Daughter ratios[a]	Measured γ-ray ratio		Average value	
			1	2	Ratio	Scale
^{18}Ne	^{18}F	$\gamma_{660}/\gamma_{1042}$	0.021 ± 0.003 [Ha75]	0.0169 ± 0.0004 [He82]		
			0.0172 ± 0.0005 [Ad83]		0.0171 ± 0.0003	1.0
^{26}Si	^{26}Al	$\gamma_{1622}/\gamma_{829}$	0.149 ± 0.016 [Mo71]	0.134 ± 0.005 [Ha75]		
			0.1245 ± 0.0023 [Wi80]	0.1301 ± 0.0062 [Ma08]	0.1269 ± 0.0026	1.3
		$\gamma_{1655}/\gamma_{829}$	0.00145 ± 0.00032 [Wi80]		0.0015 ± 0.0003	
		$\gamma_{1843}/\gamma_{829}$	0.013 ± 0.003 [Mo71]	0.016 ± 0.003 [Ha75]		
			0.01179 ± 0.00027 [Wi80]		0.0118 ± 0.0003	1.0
		$\gamma_{2512}/\gamma_{829}$	0.00282 ± 0.00010 [Wi80]		0.0028 ± 0.0001	
		$\gamma_{\text{total}}/\gamma_{829}$			0.1430 ± 0.0026	
^{30}S	^{30}P	$\gamma_{709}/\gamma_{677}$	0.006 ± 0.003 [Mo71]	0.0037 ± 0.0009 [Wi80]	0.0039 ± 0.0009	1.0
		$\gamma_{2341}/\gamma_{677}$	0.033 ± 0.002 [Mo71]	0.0290 ± 0.0006 [Wi80]	0.0293 ± 0.0011	1.9
		$\gamma_{3019}/\gamma_{677}$	0.00013 ± 0.00006 [Wi80]		0.0001 ± 0.0001	
		$\gamma_{\text{total}}/\gamma_{677}$			0.0334 ± 0.0014	
^{34}Ar	^{34}S	$\gamma_{461}/\gamma_{666}$	0.28 ± 0.16 [Mo71]	0.365 ± 0.036 [Ha74a]	0.361 ± 0.035	1.0
		$\gamma_{2580}/\gamma_{666}$	0.38 ± 0.09 [Mo71]	0.345 ± 0.01 [Ha74a]	0.345 ± 0.010	1.0
		$\gamma_{3129}/\gamma_{666}$	0.67 ± 0.08 [Mo71]	0.521 ± 0.012 [Ha74a]	0.524 ± 0.022	1.8
		$\gamma_{\text{total}}/\gamma_{666}$			1.231 ± 0.043	
^{42}Ti	^{42}Sc	$\gamma_{2223}/\gamma_{611}$	0.012 ± 0.004 [Ga69]		0.012 ± 0.004	
		$\gamma_{\text{total}}/\gamma_{611}$	0.023 ± 0.012 [Ga69, En90]		0.023 ± 0.012	

[a] γ-ray intensities are denoted by γ_E, where E is the γ-ray energy in keV.

TABLE VI. References for which the original decay-energy results have been updated to incorporate the most recent calibration standards. (See Table VIII for the correlation between the alphanumeric reference code used in this table and the actual reference numbers.)

References (parent nucleus)[a]	Update procedure
Bo64 (^{18}Ne), Ba84 (^{10}C), Br94 (^{26}Al^m) Ba98 (^{10}C), Ha98 (^{38}K^m), To03 (^{14}O)	We have converted all original (p,n) threshold measurements to Q values using the most recent mass excesses [Au03].
Ry73a (^{34}Cl), Sq76 (^{46}V), Ba77c (^{34}Cl) Wh77 (^{14}O)	These (p,n) threshold measurements have been adjusted to reflect recent calibration α energies [Ry91] before being converted to Q values.
Pr67 (^{18}Ne)	Before conversion to a Q value, this (p,n) threshold was adjusted to reflect a new value for the ^{7}Li(p,n) threshold [Wh85], which was used as calibration.
Ja78 (^{38}K^m)	This (p,n) threshold was measured relative to those for ^{10}C and ^{14}O; we have adjusted it based on average Q values obtained for those decays in this work.
Bu79 (^{38}K^m)	Before conversion to a Q value, this (p,n) threshold was adjusted to reflect the modern value for the ^{35}Cl(p,n) threshold [Au03], which was used as calibration.
Bu61 (^{14}O), Ba62 (^{14}O)	These ^{12}C(^{3}He, n) threshold measurements have been adjusted for updated calibration reactions based on current mass excesses [Au03].
Ha74d (^{34}Cl)	These (^{3}He,t) reaction Q values were calibrated by the ^{27}Al(^{3}He,t) reaction to excited states in ^{27}Si; they have been revised according to modern mass excesses [Au03] and excited-state energies [En98].
Ba88 and Ba89 (^{10}C)	These measurements of excitation energies in ^{10}B have been updated to modern γ-ray standards [He00].
Ki89 (^{42}Sc)	This ^{41}Ca(p,γ) reaction Q value was measured relative to that for ^{40}Ca(p,γ); we have slightly revised the result based on modern mass excesses [Au03].
Ha74c (^{22}Mg, ^{26}Si, ^{30}S, ^{34}Ar)	These (p,t) reaction Q values have been adjusted to reflect the current Q value for the ^{16}O(p,t) reaction [Au03], against which they were calibrated.

[a] These references all appear in Table I under the appropriate parent nucleus.

TABLE VII. References from which some or all results have been rejected even though their quoted uncertainties qualified them for inclusion. (See Table VIII for the correlation between the alphanumeric reference code used in this table and the actual reference numbers.)

References (parent nucleus)	Reason for rejection
1. Decay energies	
Pa72 (^{30}S)	No calibration is given for the measured (p,t) reaction Q values; update is clearly required but none is possible.
No74 (^{22}Mg)	Calibration reaction Q values have changed but calibration process is too complex to update.
Ro74 (^{10}C)	P. H. Barker (co-author) later considered that inadequate attention had been paid to target surface purity [Ba84].
Ba77b (^{10}C)	P. H. Barker (co-author) later stated [Ba84] that the (p,t) reaction Q value could not be updated to incorporate modern calibration standards.
Vo77 (^{14}O, ^{26}Al^m, ^{34}Cl, ^{42}Sc, ^{46}V, ^{50}Mn, ^{54}Co)	Most of the results in this reference disagree significantly with more recent and accurate measurements. Our justification for rejection is presented in more detail in the text.
Wh81 and Ba98 (^{14}O)	The result in [Wh81] was updated in [Ba98] but then eventually withdrawn by P. H. Barker (co-author) in [To03].
2. Half-lives	
Ja60 (^{26}Al^m), He61 (^{14}O), Ba62 (^{14}O), Fr63 (^{14}O), Fr65b (^{42}Sc, ^{46}V, ^{50}Mn) Si72 (^{14}O)	Quoted uncertainties are too small, and results likely biased, in light of statistical difficulties more recently understood (see [Fr69a]). In particular, "maximum-likelihood" analysis was not used.
Ha72a (^{26}Al^m, ^{34}Cl, ^{38}K^m, ^{42}Sc)	All four quoted half-lives are systematically higher than more recent and accurate measurements.
Ro74 (^{10}C)	P. H. Barker (co-author) later considered that pile-up had been inadequately accounted for [Ba90].
Ch84 (^{38}K^m)	"Maximum-likelihood" analysis was not used.
3. Branching ratios	
Fr63 (^{26}Si)	Numerous impurities are present; result is obviously wrong.

have shrunk even further. The level of precision possible has currently reached ~0.001%, at least for the Q_{EC} values of ^{50}Mn and ^{54}Co, so it is now necessary to include in the f calculation a provision for atomic excitation of the daughter nucleus if the calculation is to continue to match the precision of the input data. Our method for including this effect is described in Appendix A; and we also present there, in Table X, a comparison of f values both with and without this small correction. It can be seen that the effect of the correction is comparable to a shift of 0.001%–0.004% in the Q_{EC} value, an amount significant enough to warrant its inclusion in future. Our final f values are recorded in the second column of Table IX. They were evaluated by using our updated code and the Q_{EC} values with their uncertainties from column 7 of Table I.

The third column of Table IX lists (as percentages) the electron-capture fraction, P_{EC}, calculated for each of the 20 superallowed transitions. The method of calculation was described in our last survey [5], to which the reader is referred for more details. The partial half-life, t, for each transition is then obtained from its total half-life, $t_{1/2}$, branching ratio, R, and electron-capture fraction according to the following formula:

$$t = \frac{t_{1/2}}{R}(1 + P_{EC}). \tag{2}$$

The resultant values for the partial half-lives and the corresponding ft values appear in columns 4 and 5 of the table.

To obtain the $\mathcal{F}t$ value from each ft value, we use Eq. (1) to apply the small transition-dependent correction terms, δ'_R, δ_{NS}, and δ_C. We take the values of these terms from our recent re-evaluation of the corrections to superallowed β decay [7]. The first term, δ'_R, which is listed in column 6 of Table IX, is taken from Table V in Ref. [7]. The two nuclear-structure-dependent terms, combined in the form $(\delta_C - \delta_{NS})$, are listed in column 7. In Ref. [7], δ_C is expressed as the sum of δ_{C1} and δ_{C2}, the former being listed in Table III of that reference and the latter in Table II; δ_{NS} is taken from Table VI of the same reference. Finally, the resulting $\mathcal{F}t$ values are listed in column 8 of our Table IX.

Both the uncorrected ft values and the fully corrected $\mathcal{F}t$ values are plotted in Fig. 3 for the 13 most precisely measured transitions. The differences between the former, in the top panel, and the latter, in the bottom panel, illustrate the effects of our including the correction terms. It is also worth remarking that the values of δ'_R are very nearly the same for 11 of the 13 cases plotted: Only ^{10}C and ^{14}O have slightly higher values. Thus, most of the differences between the two panels of the figure are due to the effects of the nuclear-structure-dependent terms, δ_{NS} and δ_C.

A. CVC test

There are now 13 superallowed transitions whose $\mathcal{F}t$ values have uncertainties less than ±0.4%, with the best case, ^{26}Al^m, being known an order of magnitude better than that. These

TABLE VIII. Reference key, relating alphanumeric reference codes used in Tables I–VII to the actual reference numbers.

Table code	Reference number	Table code	Reference number	Table code	Reference number	Table code	Reference number	Table code	Reference number	Table code	Reference number
Ad83	[9]	Aj88	[10]	Aj91	[11]	Al69	[12]	Al72	[13]	Al75	[14]
Al77	[15]	Al78	[16]	Al82	[17]	An70	[18]	Au03	[19]	Az74	[20]
Az75	[21]	Ba62	[22]	Ba67	[23]	Ba77a	[24]	Ba77b	[25]	Ba77c	[26]
Ba84	[27]	Ba88	[28]	Ba89	[29]	Ba90	[30]	Ba98	[31]	Ba00	[32]
Ba01	[33]	Ba04	[34]	Ba06	[35]	Ba08	[36]	Be68	[37]	Be78	[38]
Be85	[39]	Be08	[40]	Bi03	[41]	Bl04a	[42]	Bl04b	[43]	Bo64	[44]
Br94	[45]	Bu61	[46]	Bu79	[47]	Bu88	[48]	Bu06	[49]	Ca05	[50]
Ch84	[51]	Cl73	[52]	Da80	[53]	Da85	[54]	De69	[55]	De78	[56]
Dr75	[57]	En90	[58]	En98	[59]	Er06a	[60]	Er06b	[61]	Er08	[62]
Fa09	[63]	Fi08	[64]	Fr63	[65]	Fr65b	[66]	Fr69a	[67]	Fr75	[68]
Fu99	[69]	Ga69	[70]	Ga01	[71]	Ge07	[72]	Ge08	[73]	Gi72	[74]
Go72	[75]	Gr07	[76]	Gr08	[77]	Ha67	[78]	Ha68	[79]	Ha72a	[80]
Ha74a	[81]	Ha74b	[82]	Ha74c	[83]	Ha74d	[84]	Ha75	[85]	Ha94	[86]
Ha98	[87]	Ha02	[88]	Ha03	[89]	He61	[90]	He81	[91]	He82	[92]
He00	[93]	He01	[94]	He02	[95]	Ho64	[96]	Ho74	[97]	Hu82	[98]
Hy03	[99]	Hy05	[100]	Ia06	[101]	Ia08	[102]	In77	[103]	Is80	[104]
Ja60	[105]	Ja78	[106]	Je07	[107]	Ji02	[108]	Ka68	[109]	Ka69	[110]
Ke07	[111]	Ki89	[112]	Ki91	[113]	Ko83	[114]	Ko87	[115]	Ko97a	[116]
Ko97b	[117]	Kr91	[118]	Le08	[119]	Li94	[120]	Ma94	[121]	Ma08	[122]
Mc67	[123]	Mi67	[124]	Mo71	[125]	Mu04	[126]	Na91	[127]	Ni69	[128]
No74	[129]	Oi01	[130]	Pa72	[131]	Pa05	[132]	Pi03	[133]	Pr67	[134]
Pr90	[135]	Ra83	[136]	Re85	[137]	Ri07	[138]	Ro70	[139]	Ro72	[140]
Ro74	[141]	Ro75	[142]	Ro06	[143]	Ry73a	[144]	Ry91	[145]	Sa80	[146]
Sa95	[147]	Sa04	[148]	Sa05	[149]	Sc05	[150]	Sc07	[151]	Se73	[152]
Se05	[153]	Sh55	[154]	Si66	[155]	Si72	[156]	Sq75	[157]	Sq76	[158]
Ti95	[159]	To03	[160]	To05	[161]	Vo77	[162]	Wa83	[163]	Wa92	[164]
We68	[165]	Wh77	[166]	Wh81	[167]	Wh85	[168]	Wi76	[169]	Wi78	[170]
Wi80	[171]	Zi72	[172]	Zi87	[173]						

data are sufficient to provide a very demanding test of the CVC assertion that the $\mathcal{F}t$ values should be constant for all nuclear superallowed transitions of this type. The data in column 8 of Table IX clearly satisfy the test, the weighted average of the 13 most precise results being

$$\overline{\mathcal{F}t} = 3072.08 \pm 0.79 \text{ s}, \tag{3}$$

with a corresponding chi-square per degree of freedom of $\chi^2/\nu = 0.29$. That these 13 $\mathcal{F}t$ values form a consistent set is also clearly evident from the bottom panel of Fig. 3. Since $\mathcal{F}t$ is inversely proportional to the *square* of the vector coupling constant, G_V, then Eq. (3) can be said to confirm the constancy of G_V—and to verify this key component of the CVC hypothesis—at the level of 1.3×10^{-4}.

Compared with the results of our last survey [5], the value of $\overline{\mathcal{F}t}$ in Eq. (3) is somewhat lower but carries a similar uncertainty. However, the new analysis is more demanding since, for the first time, it includes the ^{62}Ga transition, which has been very significantly improved in the past four years. This effectively increases the span of masses over which the CVC test is being applied; yet even with this addition, the χ^2/ν is actually lower than it was previously. The small reduction in the central value of $\overline{\mathcal{F}t}$ is within the uncertainty of the previous value; it has arisen from a combination of factors: changes in the experimental data base, which have acted to increase $\overline{\mathcal{F}t}$, and recent improvements in the nuclear-structure-dependent correction terms [7], which have acted to lower it.

B. $\mathcal{F}t$ value error budgets

We show the contributing factors to the individual $\mathcal{F}t$ value uncertainties in Fig. 4 for the "traditional nine" cases and in Fig. 5 for the remaining 11. For most of the cases that contribute to the CVC test—^{26}Al^m to ^{54}Co in Fig. 4 as well as ^{62}Ga and ^{74}Rb in Fig. 5—the theoretical uncertainties are greater than, or comparable to, the experimental ones. In these cases, the nuclear-structure-dependent correction, $\delta_C - \delta_{NS}$, contributes an uncertainty of 3–7 parts in 10^4 to all $\mathcal{F}t$ values between ^{26}Al^m and ^{54}Co but jumps up to 20–30 parts in 10^4 for ^{62}Ga and ^{74}Rb because of nuclear-model ambiguities. For its part, the nucleus-dependent radiative correction, δ'_R, has an uncertainty that starts very small but grows smoothly with Z^2. This is because the contribution to δ'_R from order $Z^2\alpha^3$ has only been estimated from its leading logarithm [176] and the magnitude of this estimate has been taken as the uncertainty in δ'_R. As a result, though, for ^{50}Mn and ^{54}Co it becomes the leading uncertainty, indicating that a closer look at the order $Z^2\alpha^3$ contribution to δ'_R would certainly now be worthwhile.

TABLE IX. Derived results for superallowed Fermi β decays.

Parent nucleus	f	P_{EC} (%)	Partial half-life t (ms)	ft (s)	δ_R' (%)	$\delta_C - \delta_{NS}$ (%)	$\mathcal{F}t$ (s)
$T_z = -1$							
^{10}C	2.3004 ± 0.0012	0.297	1322300 ± 1800	3041.7 ± 4.3	1.679 ± 0.004	0.520 ± 0.039	3076.7 ± 4.6
^{14}O	42.772 ± 0.023	0.088	71127 ± 51	3042.3 ± 2.7	1.543 ± 0.008	0.575 ± 0.056	3071.5 ± 3.3
^{18}Ne	134.47 ± 0.15	0.081	21660 ± 590	2912 ± 79	1.506 ± 0.012	0.855 ± 0.052	2931 ± 80
^{22}Mg	418.39 ± 0.17	0.069	7295 ± 17	3052.0 ± 7.2	1.466 ± 0.017	0.605 ± 0.030	3078.0 ± 7.4
^{26}Si	1029.4 ± 2.2	0.064	2954 ± 23	3041 ± 24	1.438 ± 0.023	0.650 ± 0.034	3065 ± 25
^{30}S	1966.9 ± 8.0	0.066	1524 ± 21	2998 ± 44	1.423 ± 0.029	1.040 ± 0.032	3009 ± 44
^{34}Ar	3414.5 ± 1.5	0.069	894.0 ± 2.4	3052.7 ± 8.2	1.412 ± 0.035	0.845 ± 0.058	3069.6 ± 8.5
^{38}Ca	5327.2 ± 1.8	0.075			1.414 ± 0.042	0.940 ± 0.072	
^{42}Ti	7040 ± 30	0.088	470 ± 160	3300 ± 1100	1.428 ± 0.050	1.170 ± 0.080	3300 ± 1100
$T_z = 0$							
^{26}Al^m	478.237 ± 0.080	0.082	6350.2 ± 1.9	3036.9 ± 0.9	1.478 ± 0.020	0.305 ± 0.027	3072.4 ± 1.4
^{34}Cl	1995.96 ± 0.47	0.080	$1527.77^{+0.44}_{-0.47}$	$3049.4^{+1.1}_{-1.2}$	1.443 ± 0.032	0.735 ± 0.048	3070.6 ± 2.1
^{38}K^m	3297.88 ± 0.34	0.085	925.42 ± 0.28	3051.9 ± 1.0	1.440 ± 0.039	0.755 ± 0.060	3072.5 ± 2.4
^{42}Sc	4472.24 ± 1.15	0.099	681.44 ± 0.26	3047.6 ± 1.4	1.453 ± 0.047	0.630 ± 0.059	3072.4 ± 2.7
^{46}V	7209.47 ± 0.90	0.101	422.99 ± 0.11	3049.5 ± 0.9	1.445 ± 0.054	0.655 ± 0.063	3073.3 ± 2.7
^{50}Mn	10745.97 ± 0.51	0.107	283.68 ± 0.11	3048.4 ± 1.2	1.444 ± 0.062	0.695 ± 0.055	3070.9 ± 2.8
^{54}Co	15766.6 ± 2.9	0.111	$193.495^{+0.063}_{-0.086}$	$3050.8^{+1.1}_{-1.5}$	1.443 ± 0.071	0.805 ± 0.068	$3069.9^{+3.2}_{-3.3}$
^{62}Ga	26400.2 ± 8.3	0.137	116.441 ± 0.042	3074.1 ± 1.5	1.459 ± 0.087	1.52 ± 0.21	3071.5 ± 7.2
^{66}As	32125 ± 470	0.155			1.468 ± 0.095	1.62 ± 0.40	
^{70}Br	38600 ± 3600	0.175			1.49 ± 0.11	1.69 ± 0.25	
^{74}Rb	47300 ± 110	0.194	65.227 ± 0.078	3084.9 ± 7.8	1.50 ± 0.12	1.71 ± 0.31	3078 ± 13
					Average (best 13), $\overline{\mathcal{F}t}$		3072.08 ± 0.79
					χ^2/ν		0.29

For all the transitions from $T_z = 0$ parent nuclei, the experimental branching ratios are >99% and have very small associated uncertainties with the exception of ^{54}Co, which has a 3×10^{-4} fractional uncertainty attributed to its branching ratio, and ^{74}Rb, which has 10×10^{-4}. In both cases, this is because they are predicted to have weak Gamow-Teller branches that have not yet been observed. We have used an estimate of the strength of the missing branches, taken from a shell-model calculation [88], to assign an uncertainty to the superallowed branching ratio. Numerous weak Gamow-Teller branches become an increasingly significant issue for the heavier-mass nuclei with $A \geqslant 62$, where they present a major experimental challenge if they are to be fully characterized. Only in the case of ^{62}Ga has this been accomplished so far.

For the decays of ^{10}C and ^{14}O, and for all the decays depicted in Fig. 5 except for ^{62}Ga and ^{74}Rb, the predominant uncertainties are experimental in origin. Many of the experimental branching ratios and some of the Q values and half-lives have yet to be measured precisely for the cases in Fig. 5, but recent advances in experimental techniques have been improving this situation and are likely to improve it even more within the next few years.

C. Accounting for systematic uncertainties

So far, we have dealt with the internuclear behavior of $\mathcal{F}t$ values, examining their constancy as a test of CVC. With that test passed at high precision, we are now in a position to use the average $\mathcal{F}t$ value obtained from these concordant nuclear data to go beyond nuclei, obtaining first the vector coupling constant [see Eq. (1)] and then the V_{ud} matrix element. Before doing so, however, we must address one more possible source of uncertainty. Though the average $\mathcal{F}t$ value given in Eq. (3) includes a full assessment of the uncertainties attributable to experiment and to the particular calculations used to obtain the correction terms, it does not incorporate any provision for a common systematic error that could arise from the type of calculation chosen to model the nuclear-structure effects. To discuss this, we divide the problem into two parts: the accuracy of the model as an approximation to the formally complete treatment and the possible existence of systematic uncertainties within the model.

1. The model approximation

Very recently Miller and Schwenk [177] have explored the formally complete approach to isospin-symmetry breaking. Their starting point is to define the Fermi matrix element as

$$M_F = \langle f|\tau_+|i\rangle = \sum_\alpha \langle f|a_\alpha^\dagger b_\alpha|i\rangle = \sum_{\alpha,\pi} \langle f|a_\alpha^\dagger|\pi\rangle\langle\pi|b_\alpha|i\rangle, \tag{4}$$

where $a_\alpha^\dagger$ creates a neutron and b_α annihilates a proton in state α. Here $|i\rangle$ and $|f\rangle$ are the *exact* state vectors for the full Hamiltonian. If this Hamiltonian commutes with the isospin

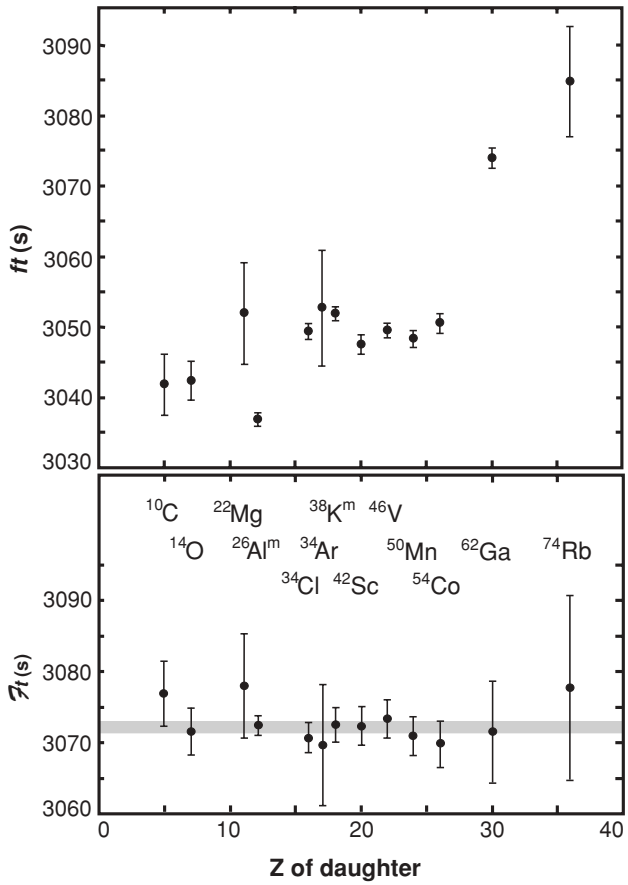

FIG. 3. The uncorrected experimental ft values as a function of the charge on the daughter nucleus (top panel) and the corresponding $\mathcal{F}t$ values (bottom panel); the latter differ from the ft values by the inclusion of the correction terms δ'_R, δ_{NS}, and δ_C. The horizontal gray band in the bottom panel gives one standard deviation around the average $\overline{\mathcal{F}t}$ value.

operators, then $|i\rangle$ and $|f\rangle$ are exact isospin analogs of each other, $\langle\pi|b_\alpha|i\rangle = \langle f|a^\dagger_\alpha|\pi\rangle^*$, and the symmetry-limit matrix

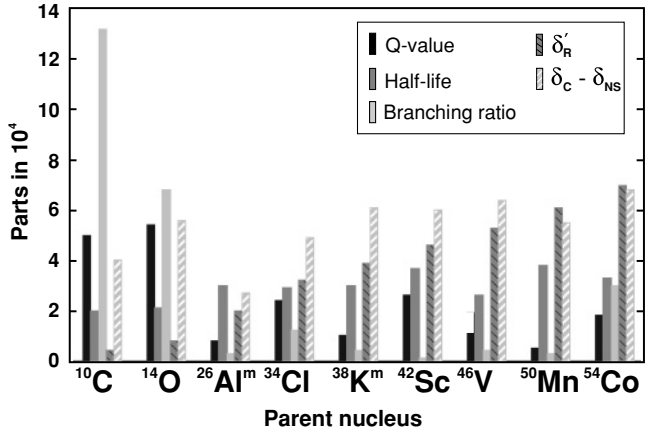

FIG. 4. Summary histogram of the fractional uncertainties attributable to each experimental and theoretical input factor that contributes to the final $\mathcal{F}t$ values for the "traditional nine" superallowed transitions.

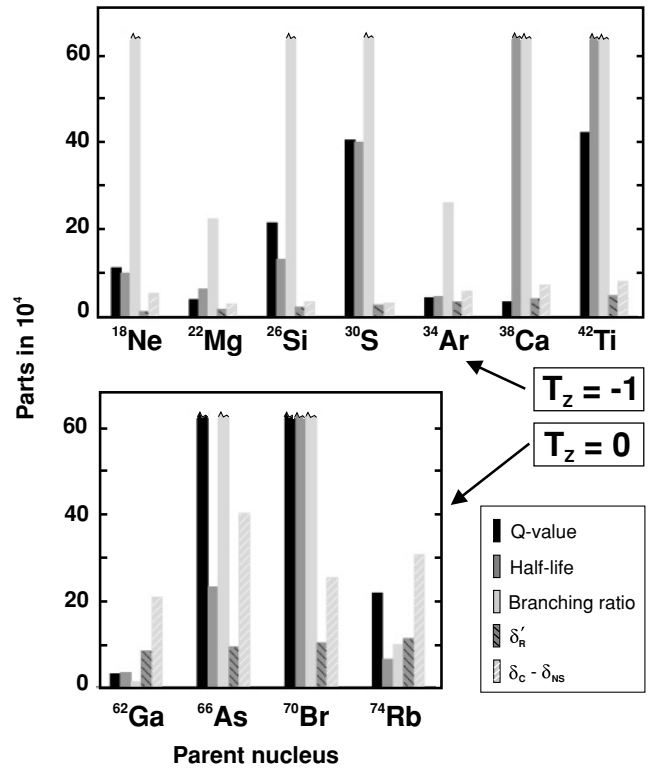

FIG. 5. Summary histogram of the fractional uncertainties attributable to each experimental and theoretical input factor that contributes to the final $\mathcal{F}t$ values for the 11 other superallowed transitions. Where the error is shown as exceeding 60 parts in 10^4, no useful experimental measurement has been made.

element is

$$M_0 = \sum_{\alpha,\pi} |\langle f|a^\dagger_\alpha|\pi\rangle|^2. \tag{5}$$

If isospin is not an exact symmetry, then $|i\rangle$ and $|f\rangle$ are not isospin analogs and a correction to M_0 needs to be evaluated. This is the isospin-symmetry-breaking correction, δ_C, we seek to determine. It is defined by

$$M_F^2 = M_0^2(1-\delta_C). \tag{6}$$

Ideally, to obtain δ_C one would compute Eq. (4) using the shell model and introduce Coulomb and other charge-dependent terms into the shell-model Hamiltonian. However, because the Coulomb force is long range, the shell-model space would have to be huge to include all the potential states with which the Coulomb interaction might connect. Currently, this is not a practical proposition.

To proceed with a manageable calculation, we have developed a model approach [7,178,179] in which δ_C is divided into two parts:

$$\delta_C = \delta_{C1} + \delta_{C2}. \tag{7}$$

For δ_{C1}, we compute

$$\sum_{\alpha,\pi} \langle \bar{f}|a^\dagger_\alpha|\pi\rangle\langle\pi|b_\alpha|\bar{i}\rangle = M_0(1-\delta_{C1})^{1/2}, \tag{8}$$

where $|\bar{i}\rangle$ and $|\bar{f}\rangle$ are not the exact eigenstates that appear in Eq. (4) but are the shell-model eigenstates of an effective Hamiltonian (including charge-dependent terms) evaluated in a modest-sized shell-model space. Since this space does not allow for nodal mixing, we correct for that limitation by computing the second component, δ_{C2}, obtained from

$$\sum_{\alpha,\pi} |\langle \bar{f}|a^\dagger_\alpha|\pi\rangle|^2 r^\pi_\alpha = M_0(1-\delta_{C2})^{1/2}, \tag{9}$$

where each r^π_α is a radial overlap integral of proton and neutron radial functions. We justify the efficacy of this second term by the following arguments: If the radial functions were identical, then δ_{C2} would vanish as it should. Otherwise the proton radial functions, $u^p(r)$, could be expanded in terms of a complete set of neutron functions, $u^n_N(r)$, including all possible radial nodes, N:

$$u^p(r) = \sum_N a_N u^n_N(r). \tag{10}$$

The isospin-symmetry-breaking correction, δ_{C2}, could then be expressed in terms of a_N, which from perturbation theory could itself be written in terms of matrix elements of the Coulomb interaction. This would be equivalent to the nodal mixing included in Eq. (4) but left out of the calculation of δ_{C1} in Eq. (8). The idea is that δ_{C1} is the result of a tractable shell-model calculation that does not include any nodal mixing, while δ_{C2} then corrects for the nodal mixing that would be present if the shell-model space were larger.

Clearly our charge-dependent correction terms [7,178,179] are based on a model and required approximations to make the computation possible. Since no one has yet made a complete calculation without approximations, it is impossible to be definitive about any systematic errors that might be introduced by our methods. Only for the lightest superallowed emitter, ^{10}C, has it been possible so far even to come close to an exact treatment. Caurier *et al.* [180] have reported a large no-core shell-model calculation for that system but, even though they were able to extend their basis states up to $8\hbar\omega$, their calculated δ_C still had not converged to a stable value. However, they used their results together with perturbation theory to estimate that the full value of δ_C should be about 0.19%. This result, which in effect used Eq. (4) and did not split δ_C into two parts, agrees completely with our calculated value for $\delta_C = \delta_{C1} + \delta_{C2}$ of 0.18(2)% (see Table VII in Ref. [7]). This agreement certainly supports the validity of our model.

Furthermore, it must be noted that our model approach has allowed us to use well-established shell-model and related parameters, which were determined from experimental data that are completely independent of the superallowed ft values. As is clearly evident from Fig. 3, these calculated corrections do a remarkable job in converting widely scattered ft values into a consistent set of $\mathcal{F}t$ values. Not only that but, as shown in Ref. [7], they also closely reproduce the measured results for isospin-forbidden $0^+ \rightarrow 0^+$ β transitions in all nuclei for which the shell-model calculation is well specified. (This is not the case for ^{62}Ga.) Of course, although these successes demonstrate that our calculated δ_C values correctly reproduce the nucleus-to-nucleus variations observed by experiment, they cannot rule out a constant shift in the corrections for all nuclei. Even so, it seems highly unlikely that a faulty approximation could lead to relative results that are correct in every detail, while being consistently wrong—and by the same constant amount—in the absolute values for each and every case.

Under the circumstances, we see no justification at this time to assign any additional systematic error to account for possible inadequacies of the model we use to calculate the charge-dependent correction terms.

2. *Systematic uncertainty within the model*

As introduced in Eq. (7) the isospin-symmetry-breaking correction, δ_C, is separated into two pieces: δ_{C1} comes from configuration mixing in a modest-sized shell-model calculation with charge-dependent interactions and δ_{C2} involves radial overlap integrals. The calculation of δ_{C1}, the smaller of the two terms, requires a reliable shell-model description of the nuclei involved but it can also be further constrained by independent experimental data—for example, the measured coefficients of the corresponding isobaric multiplet mass equation (IMME). We take our corrections from Ref. [7] where the uncertainties attached to the calculated values of δ_{C1} for the 20 cases of interest here already include ample provision for differences between competing shell-model parametrizations as well as for experimental uncertainties on the IMME coefficients used as constraints.

The values of the radial-overlap term, δ_{C2}, which we use as input to Table IX, were also taken from our recently published calculations [7]. Those calculations used radial wave functions derived from a Saxon-Woods potential with either the well depth or one of the surface terms in the potential adjusted so that the binding energy of each computed eigenfunction matched the corresponding measured separation energy. The quoted uncertainties included provision for any variations in the results depending on which parameter was used in the adjustment. However, no provision was included for possible differences that might occur if another method entirely were used to derive the radial wave functions. In the past [5], we have accounted for this uncertainty by comparing our results for δ_{C2} with those of Ormand and Brown [181–183], who used Hartree-Fock eigenfunctions and obtained consistently smaller corrections than those we found with a Saxon-Woods potential. We treated this as a valid source of systematic uncertainty and incorporated it by deriving two average $\overline{\mathcal{F}t}$ values, one for each set of δ_{C2} calculations, then taking the average of the two and assigning a systematic uncertainty equal to half the spread between them.

This specific comparison is no longer tenable. The Ormand and Brown calculations are in some cases more than two decades old: They use smaller shell-model spaces than are now known to be necessary [7] and they are not available at all for some of the transitions that we now need to include. To remedy these deficiencies we have undertaken our own Hartree-Fock calculations. They are described in detail in Appendix B, where Table XI lists the values of δ_{C2} we compute from Hartree-Fock-derived wave functions and compares them with our earlier results from the Saxon-Woods potential [7],

the same results that we used to evaluate $\overline{\mathcal{F}t}$ in Table IX. Both methods used exactly the same shell-model calculations to determine the full parentage of the states involved.

With these new Hartree-Fock calculations we can now follow a similar procedure to the one we employed with the old calculations in our previous survey [5]. We begin by substituting the Hartree-Fock δ_{C2} values for the Saxon-Woods ones in deriving the δ_C values used in Table IX. When we do this the $\overline{\mathcal{F}t}$ value result becomes 3071.55 ± 0.89 s with $\chi^2/\nu = 0.93$. This normalized chi-square is three times what we obtained in Table IX with the Saxon-Woods corrections, which arguably could justify our rejecting the Hartree-Fock results outright. However, to be safe, we proceed as before and take the average of the Hartree-Fock and Saxon-Woods results, adding a systematic uncertainty equal to half the spread between the two results. Thus, we obtain

$$\begin{aligned}\overline{\mathcal{F}t} &= 3071.81 \pm 0.79_{\text{stat}} \pm 0.27_{\text{syst}}\ \text{s} \\ &= 3071.81 \pm 0.83\ \text{s}, \end{aligned} \tag{11}$$

where on the second line the two uncertainties have been added in quadrature. Our new systematic adjustment amounts to only 0.27 s, much smaller and of opposite sign to the 0.90-s correction applied previously [5].

It is the value for $\overline{\mathcal{F}t}$ in Eq. (11) that we carry forward to subsequent sections where we obtain V_{ud} and test the unitarity of the CKM matrix.

IV. THE IMPACT ON WEAK-INTERACTION PHYSICS

A. The value of V_{ud}

With a mutually consistent set of $\mathcal{F}t$ values, we can now use the adjusted average value in Eq. (11) to determine the vector coupling constant, G_V, from Eq. (1). The value of G_V itself is of little interest but, together with the weak interaction constant for the purely leptonic muon decay, G_F, it yields the much more interesting up-down element of the Cabibbo-Kobayashi-Maskawa (CKM) quark-mixing matrix. The basic relationship is $V_{ud} = G_V/G_F$, which in terms of $\overline{\mathcal{F}t}$ becomes

$$V_{ud}^2 = \frac{K}{2G_F^2\left(1+\Delta_R^V\right)\overline{\mathcal{F}t}}, \tag{12}$$

where Δ_R^V is the nucleus-independent radiative correction. This correction has recently been carefully re-examined by Marciano and Sirlin [6], who very substantially reduced its uncertainty. Expressing their new result in a way that is consistent with the definition of our other correction terms, we obtain (see Eq. (41) in Ref. [7])

$$\Delta_R^V = (2.361 \pm 0.038)\%. \tag{13}$$

Using the Particle Data Group (PDG) [174] value for the weak interaction coupling constant from muon decay of $G_F/(\hbar c)^3 = (1.16637 \pm 0.00001) \times 10^{-5}\ \text{GeV}^{-2}$, we obtain from Eq. (12) the result

$$|V_{ud}|^2 = 0.94916 \pm 0.00044. \tag{14}$$

Note that the total uncertainty here is almost entirely due to the uncertainties contributed by the theoretical corrections. By far the largest contribution, 0.00035, arises from the uncertainty in Δ_R^V; 0.00020 comes from the nuclear-structure-dependent corrections $\delta_C - \delta_{\text{NS}}$ and 0.00004 is attributable to δ_R'. Only 0.00016 can be considered to be experimental in origin.

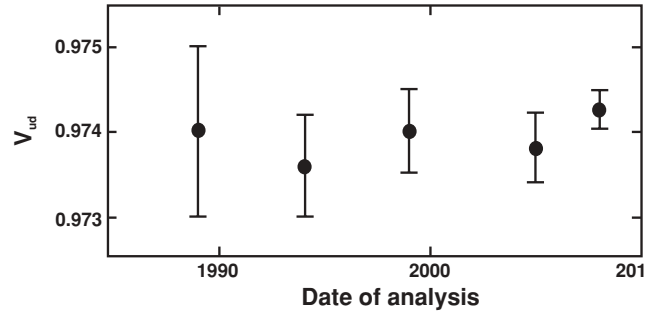

FIG. 6. Values of V_{ud} as determined from superallowed $0^+ \to 0^+$ β decays plotted as a function of analysis date, spanning the past two decades. In order, from the earliest date to the most recent, the values are taken from Refs. [4,5,184,185] and this work.

The corresponding value of V_{ud} is

$$|V_{ud}| = 0.97425 \pm 0.00022, \tag{15}$$

a result that is consistent with, but more precise than, values we have obtained in previous analyses of superallowed β decay. To emphasize the consistency and steady improvement that has characterized the value of V_{ud} as derived from nuclear β decay, in Fig. 6 we plot our new result together with V_{ud} values published at various times over the past two decades [4,5,184,185]. However, this steady improvement evident from survey to survey does not reveal the dramatic events that have occurred between the last survey, published in 2005, and the present one. Although the four-year net change in V_{ud} has been within the previously quoted uncertainty, if snapshots of V_{ud} had been taken at regular intervals during the past four years, would they have revealed larger upheavals?

As explained in Sec. II B, the upheavals began with the first Penning-trap Q_{EC} value results. These results ultimately demonstrated that an important set of reaction-based Q_{EC} value measurements, which were published in 1977 and have played an important role in data surveys ever since, were seriously flawed. The removal of this data set affected seven Q_{EC} values in all and was perhaps the most serious blow that could have been dealt to the data base in a single stroke. To isolate the impact of these and other experimental changes from the impact of the theoretical changes that have also taken place during the same time period, we have analyzed the experimental data in Tables I–IV using exactly the same theoretical corrections as were used in the 2005 survey. The result for $\overline{\mathcal{F}t}$ is 3074.1(7) s, with $\chi^2/\nu = 1.1$, compared with the 2005 value of 3072.7(8) s, with $\chi^2/\nu = 0.42$. The shift in the central value is nearly twice the standard deviation of the 2005 result. Even so, if we proceed in the same spirit to calculate V_{ud} using the new experimental data combined with the 2005 theoretical corrections, we obtain the result 0.9736(4), which is to be compared with the 2005 value quoted for V_{ud} of 0.9738(4). In contrast to the $\overline{\mathcal{F}t}$ values, these numbers agree well within their uncertainties, although it has to be noted that the error bar quoted for V_{ud} was—and

still is—dominated by provisions for theoretical uncertainties. It is this large window for possible theoretical shortcomings that absorbs the unexpectedly large shift in the experimental input instead.

Clearly the large number of measurements that contribute to the nuclear determination of V_{ud} provides a robust data base, one that has proven capable of absorbing even a severe blow without causing a radical shift in the critical results. There is a further feature of these nuclear decays, whose importance is also illustrated by the events of the past four years: the fact that there are many transitions—currently 13—that all should give consistent results. As we have already indicated, the new experimental data base analyzed with the 2005 theoretical corrections yields a χ^2/ν value for $\overline{\mathcal{F}t}$ that is nearly three times larger than the result in 2005. It is also obvious by a visual inspection of a plot of these $\mathcal{F}t$ values that something is wrong, and it was this observation that led to improvements in the theoretical corrections, which have been applied in this survey.

B. Unitarity of the CKM matrix

The CKM matrix transforms one set of quark basis states into another: It transforms the quark-mass eigenstates into the weak-interaction eigenstates. If both sets are complete and orthonormal, then the transformation matrix itself must be unitary. The standard model does not prescribe the individual elements of the CKM matrix—they must be determined experimentally—but absolutely fundamental to the model is the requirement that the matrix be unitary. Whether unitarity is satified in practice can be tested experimentally, the severity of the test depending on the precision with which the CKM matrix elements can be determined.

To date, the most demanding test of CKM unitarity comes from the sum of squares of the top-row elements, $|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2$, which should equal one. Since $|V_{ud}|^2$ constitutes 95% of this sum, the precision on V_{ud} is of paramount importance. The value of $|V_{ud}| = 0.97425(22)$ derived in Sec. IV A has a precision of 0.02%, which is the most precise result so far obtained for this matrix element and is, by more than an order of magnitude, the most precisely determined value for any element in the CKM matrix. Alternative methods of obtaining V_{ud} from neutron β decay [$V_{ud} = 0.9746(18)$] and from pion β decay [$V_{ud} = 0.9749(26)$]—with both values taken from the PDG's 2008 compilation [174]—are much less precise and have been hampered by experimental difficulties. In the case of the neutron, not only is physical containment a problem but the axial-vector contribution to its β decay must be separated from the vector contribution by a β-asymmetry measurement; for pion β decay a very small branching fraction, $\mathcal{O}(10^{-8})$, must be measured.

At the time of our last survey [5] the value of V_{us} was in a state of flux. The 2004 PDG value was $|V_{us}| = 0.2200(26)$, based mostly on measurements that were at least two decades old, but new results then emerging were suggesting a value some two standard deviations higher. In the past four years, these new results on the semileptonic decays, $K_{\ell 3}$, of both charged and neutral kaons—from BNL-E865 [186], KTeV [187], NA48 [188], KLOE [189], and ISTRA+ [190]—have all combined to clarify the situation. Now, current averages by the 2008 PDG [174] and FlaviaNet [8] for kaon semileptonic branching fractions are based only on recent, high-statistics experiments, which are also consistent with one another. The best current value, presented at the CKM2008 Workshop [191] by the FlaviaNet group, is

$$f_+(0)|V_{us}| = 0.21673 \pm 0.00046. \tag{16}$$

Here $f_+(0)$ is the semileptonic decay form factor at zero-momentum transfer. Its value is close to unity. In fact, the CVC hypothesis in the exact SU(3) symmetry limit establishes its value to be exactly one, but SU(3) symmetry is broken to some extent and a theoretical calculation is required to estimate the departure of $f_+(0)$ from unity. Currently, there are two classes of evaluation: analytic or semianalytic approaches based on chiral perturbation theory [192–196] and those based on lattice QCD [197–199]. We will follow the FlaviaNet group and adopt the lattice value of $f_+(0) = 0.9644 \pm 0.0049$ from the RBC-UKQCD Collaboration [197], which yields

$$|V_{us}| = 0.2247 \pm 0.0012. \tag{17}$$

An independent determination of V_{us} can be obtained from the purely leptonic decay of the kaon, the most important mode being $K^+ \to \mu^+\nu$. If it is considered as a ratio with the leptonic decay of the pion, $\pi^+ \to \mu^+\nu$, the hadronic uncertainties can be minimized and the result yields the ratio of the CKM matrix elements $|V_{us}|/|V_{ud}|$. In the analysis of the FlaviaNet group [8] the current result is

$$\frac{|V_{us}|}{|V_{ud}|} \times \frac{f_K}{f_\pi} = 0.2760 \pm 0.0006, \tag{18}$$

where f_K and f_π are the kaon and pion decay constants, respectively. This ratio of pseudoscalar decay constants has to be obtained from theory, for which lattice QCD seems to be the only reliable source. Again, following the FlaviaNet group we adopt the lattice result from the MILC-HPQCD Collaboration [200], $f_K/f_\pi = 1.189 \pm 0.007$, and obtain

$$\frac{|V_{us}|}{|V_{ud}|} = 0.2321 \pm 0.0015. \tag{19}$$

Thus, we now have three pieces of data—$|V_{ud}|$ from nuclear decays, Eq. (15), $|V_{us}|$ from $K_{\ell 3}$ decays, Eq. (17), and the ratio $|V_{us}|/|V_{ud}|$ from $K_{\ell 2}$ decays, Eq. (19)—from which to determine two parameters, $|V_{ud}|$ and $|V_{us}|$. We perform a nonlinear least squares fit to obtain the result

$$|V_{ud}| = 0.97424(22), \quad |V_{us}| = 0.22534(93). \tag{20}$$

Note that the value of $|V_{ud}|$ obtained from this fitting procedure has only changed by one unit in the last figure compared to Eq. (15); and the change in $|V_{us}|$ compared to Eq. (17), though somewhat larger, is still well within the quoted uncertainties.

The third element of the top row of the CKM matrix, V_{ub}, is very small and hardly impacts on the unitarity test at all. Its value from the 2008 PDG compilation [174] is $|V_{ub}| = (3.93 \pm 0.35) \times 10^{-3}$. Combining this number with the ones in Eq. (20) we find the sum of the squares of the top-row elements of the CKM matrix to be

$$|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 0.99995 \pm 0.00061, \tag{21}$$

a result that shows unitarity to be fully satisfied at the 0.06% level. Only V_{us} and V_{ud} contribute perceptibly to the uncertainty and their contributions are almost equal to one another. This may seem surprising since V_{ud} is known to much higher precision than V_{us}, but it follows from the fact that $|V_{ud}|^2$ contributes 95% to the unitarity sum.

C. Limit on scalar interactions

1. Fundamental scalar current

In our previous survey [5] we explained in detail how a scalar current, if it existed, would affect the $\mathcal{F}t$ value data. We demonstrated that its effect on $\mathcal{F}t$ would be approximately proportional to $\langle 1/W \rangle$, the average inverse decay energy of each β^+ transition, so its presence would be manifest by $\mathcal{F}t$ values that are not constant as a function of Z. Since $\langle 1/W \rangle$ increases monotonically as Z decreases, the largest deviation of $\mathcal{F}t$ from constancy would occur for the superallowed transitions from nuclei with the lowest Z, ^{10}C, and ^{14}O.

We have now repeated the same analysis on our new survey results. We evaluated the statistical rate functions, f, with a shape-correction factor that included the presence of a scalar current via the Fierz interference term, b_F, which we treated as an adjustable parameter. We then sought the value of b_F that minimized χ^2 in a least-squares fit to the expression $\mathcal{F}t =$ constant. The result we obtained is

$$b_F = -0.0022 \pm 0.0026, \tag{22}$$

which is consistent with zero, as it was in 2005 [5]. In Fig. 7 we illustrate the sensitivity of this analysis by plotting the measured $\mathcal{F}t$ values together with the loci of $\mathcal{F}t$ values that would be expected if $b_F = \pm 0.004$. Obviously, the measured $\mathcal{F}t$ values do not exhibit any statistically significant curvature.

The result in Eq. (22) can also be expressed in terms of the coupling constants that Jackson, Treiman, and Wyld [201] used in writing a general form for the weak-interaction Hamiltonian. Since we are dealing only with Fermi superallowed transitions, we can restrict ourselves to scalar and vector couplings, for which that Hamiltonian becomes

$$\begin{aligned} H_{S+V} = &(\overline{\psi}_p \psi_n)(C_S \overline{\phi}_e \phi_{\overline{\nu}_e} + C'_S \overline{\phi}_e \gamma_5 \phi_{\overline{\nu}_e}) \\ &+ (\overline{\psi}_p \gamma_\mu \psi_n)[C_V \overline{\phi}_e \gamma_\mu (1+\gamma_5) \phi_{\overline{\nu}_e}], \end{aligned} \tag{23}$$

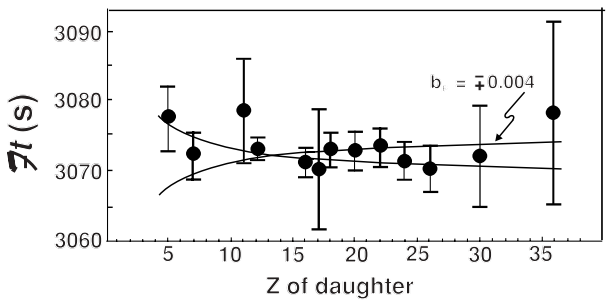

FIG. 7. Corrected $\mathcal{F}t$ values from Table IX plotted as a function of the charge on the daughter nucleus, Z. The curved lines represent the approximate loci the $\mathcal{F}t$ values would follow if a scalar current existed with $b_F = \pm 0.004$.

where we have taken the vector current to be maximally parity violating, as indicated by experiment [202]. The complexity of the relationship between b_F and the couplings C_S, C'_S, and C_V depends on what assumptions are made about the properties of the scalar current. If we take the most restrictive conditions, that the scalar and vector currents are time-reversal invariant (i.e., C_S and C_V are real) and that the scalar current, like the vector current, is maximally parity violating (i.e., $C_S = C'_S$), then we can write

$$\frac{C_S}{C_V} = -\frac{b_F}{2} = +0.0011 \pm 0.0013. \tag{24}$$

This limit from superallowed β decay is, by far, the tightest limit on the presence of a scalar current under the assumptions stated.

If we remove the condition that the scalar current be maximally parity violating, then the expression contains two unknowns,

$$b_F = \frac{-2C_V(C_S + C'_S)}{2|C_V|^2 + |C_S|^2 + |C'_S|^2} \simeq -\left(\frac{C_S}{C_V} + \frac{C'_S}{C_V}\right), \tag{25}$$

and cannot be solved individually for C_S/C_V and C'_S/C_V. However, the β-ν angular-correlation coefficient, a, for a superallowed $0^+ \to 0^+$ β transition provides another independent measure of C_S and C_V. In that case,

$$\begin{aligned} a &= \frac{2|C_V|^2 - |C_S|^2 - |C'_S|^2)}{2|C_V|^2 + |C_S|^2 + |C'_S|^2} \\ &\simeq 1 - \frac{1}{2}\left(\left|\frac{C_S}{C_V}\right|^2 + \left|\frac{C'_S}{C_V}\right|^2\right), \end{aligned} \tag{26}$$

which, together with Eq. (25), can be used set limits on both C_S/C_V and C'_S/C_V. Currently, the most precise measurement of such a β-ν angular correlation is for the superallowed decay of ^{38}K [203]. In this case, what was actually measured is $\tilde{a} = a/(1 + \gamma b_F m_e/\langle W \rangle)$, where $\gamma = \sqrt{1-(\alpha Z)^2}$ and m_e is the mass of the electron.[1] The results in terms of C_S/C_V and C'_S/C_V are plotted in Fig. 8. The value of $\tilde{a}$ taken from Ref. [203] leads to the gray annulus plotted in the figure, whereas our result for b_F from Eq. (22) is responsible for the narrow diagonal band. The intersection of these two regions, which is in black, defines the 68% confidence limit (one standard deviation) for C_S/C_V and C'_S/C_V. It corresponds to the limit

$$\left|\frac{C_S}{C_V}\right| \leqslant 0.065 \tag{27}$$

and exactly the same limit for C'_S/C_V.

The Jackson-Treiman-Wyld Hamiltonian [201] is a parametrization of one possible extension to the standard model. The coupling constants C_S and C'_S are not prescribed by the authors but are simply parameters that must be determined from experiment. One model that actually introduces scalar interactions in a natural way is the minimal supersymmetric standard model (MSSM). A valuable review of low-energy

[1] Our b_F is defined differently from the b used in Ref. [203]. The two are related by $b_F = b/\gamma$.

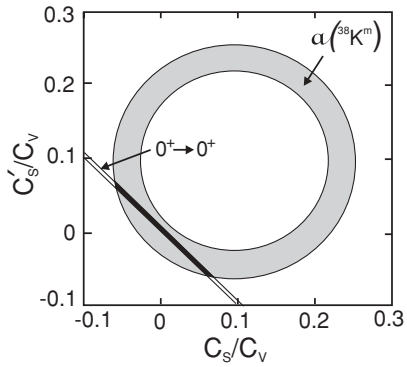

FIG. 8. Allowed range of values for C_S/C_V and C_S'/C_V as determined from the ^{38}K β-ν angular-correlation measurement [203] (gray annulus) and from superallowed $0^+ \to 0^+$ β decays (narrow diagonal band). The overlap region is shaded in black. We plot the 68% confidence limits.

tests of this model has recently been published by Ramsey-Musolf and Su [204]. In the MSSM a radiative correction to the β-decay amplitude involves box graphs with exchanged supersymmetric fermions. These produce an energy dependence in the β-decay amplitude that shows up in the Fierz interference term. Unfortunately, estimates by Profumo, Ramsey-Musolf, and Tulin [205] indicate that the resulting value of b_F would be less than 10^{-3}, which is an order of magnitude smaller than our current experimental upper limit on that quantity. Thus it will likely be some time before superallowed β decay can provide useful constraints for this class of supersymmetric models.

2. Induced scalar current

If we consider only the vector part of the weak interaction, for composite spin-1/2 nucleons the most general form of that interaction is written [206] as

$$H_V = \overline{\psi}_p(g_V\gamma_\mu - f_M\sigma_{\mu\nu}q_\nu + if_S q_\mu)\psi_n\,\overline{\phi}_e\gamma_\mu(1+\gamma_5)\phi_{\overline{\nu}_e}, \tag{28}$$

with q_μ being the four-momentum transfer, $q_\mu = (p_p - p_n)_\mu$. The values of the coupling constants g_V (vector), f_M (weak magnetic), and f_S (induced scalar) are prescribed so long as the CVC hypothesis—that the weak vector current is just an isospin rotation of the electromagnetic vector current—is correct. In particular, since CVC implies that the vector current is divergenceless, it follows that f_S should equal zero. An independent argument [207], that there be no second-class currents in the hadronic weak interaction, also requires f_S to vanish. We proved in our previous survey [5] that the presence of a nonzero f_S would manifest itself in exactly the same way as a nonzero C_S: by a $\langle 1/W\rangle$ dependence in the $\mathcal{F}t$ value data.

In the same manner that we obtained Eqs. (22) and (24), we determine from our present survey results that

$$m_e f_S/g_V = -(0.0011 \pm 0.0013). \tag{29}$$

This result is a vindication for the CVC hypothesis, which predicts $g_V = 1$ and $f_S = 0$. We confirm this prediction at the level of 24 parts in 10^4. Our result can also be interpreted as setting a limit on vector second-class currents in the semileptonic weak interaction.

D. Limits on extensions to the standard model

The unitarity sum established in Sec. IV B can be used to set limits on new physics beyond the standard model. A list of possible extensions includes, but is certainly not limited to, right-hand currents, extra Z bosons, scalars, supersymmetry, a fourth generation of quarks, and exotic muon decay. Marciano surveyed many of these possibilities at the CKM2008 Workshop [191]. His general conclusion was that, although the CKM unitarity test yields no sign of new physics, it does place important constraints on the possibilities. In the case of supersymmetric models these constraints have been explored by Ramsey-Musolf and co-workers [204,208,209]. In the minimal supersymmetric version (MSSM), corrections to low-energy observables arise only via loop effects, whereas in extensions that allow for R-parity violating (RPV) interactions new tree-level effects appear. In general, the presence of new physics may modify low-energy semileptonic electroweak observables in two ways: (i) directly, via a new semileptonic interaction (e.g. right-hand currents, MSSM with RPV interactions), and (ii) indirectly, via loop graphs contributing to the radiative correction (e.g. extra Z bosons, MSSM). In what follows we give one example of each type of modification: right-hand currents and extra Z bosons.

1. Right-hand currents

In the standard model, parity violation is considered to be maximal. What if this condition were to be relaxed? For semileptonic transitions, Herczeg [210,211] extends the general form of the weak interaction to read

$$\begin{aligned} H_{s\ell} = {} & a_{LL}(V-A)(V-A) + a_{LR}(V-A)(V+A) \\ & + a_{RL}(V+A)(V-A) + a_{RR}(V+A)(V+A), \end{aligned} \tag{30}$$

where, in each term, the first factor represents the lepton currents and the second the hadron currents. In particular, for the vector lepton current, V stands for either $\overline{\phi}_e\gamma_\mu\phi^L_{\overline{\nu}_e}$ or $\overline{\phi}_e\gamma_\mu\phi^R_{\overline{\nu}_e}$ depending on whether the chirality of the neutrino is left-handed, as it is for $V-A$ coupling, or right-handed, as it is for $V+A$ coupling. In the standard model, $a_{LL} = 1$, and $a_{LR} = a_{RL} = a_{RR} = 0$. For Fermi β decay, only the vector part of the weak hadron current contributes, so the decay rate is given by the following proportionality [5]:

$$\begin{aligned} \Gamma_\beta &\propto |a_{LL} + a_{LR}|^2 + |a_{RL} + a_{RR}|^2 \\ &\simeq |a_{LL}|^2(1 + 2\,\mathrm{Re}\,\overline{a}_{LR} + \cdots), \end{aligned} \tag{31}$$

where $\overline{a}_{LR} = a_{LR}/a_{LL}$. In the second line of the equation, we have only retained quantities that are first order in the (presumably) small quantities $\overline{a}_{LR}$, $\overline{a}_{RL}$, and $\overline{a}_{RR}$.

To determine the effect that right-hand currents would have on the value of V_{ud} obtained from experiment, we also need to consider the role of such currents on the purely leptonic muon decay. Herczeg [210] writes the effective Hamiltonian,

in analogy to Eq. (30), as

$$H_\ell = c_{LL}(V-A)(V-A) + c_{LR}(V-A)(V+A) + c_{RL}(V+A)(V-A) + c_{RR}(V+A)(V+A). \quad (32)$$

The coupling constants in Eqs. (32) and (30) are related by the CKM matrix elements through

$$a_{LL} = c_{LL}V_{ud}^L, \quad a_{LR} = c_{LR}e^{i\alpha}V_{ud}^R, \quad a_{RL} = c_{RL}V_{ud}^L, \quad a_{RR} = c_{RR}e^{i\alpha}V_{ud}^R. \quad (33)$$

Here V_{ud}^L is the element of the CKM matrix for left-handed chirality quarks, and V_{ud}^R is for right-handed chirality quarks. The phase α is a CP-violating phase in the right-handed CKM matrix. The decay rate for muon decay is constructed from an equal mix of vector and axial-vector interactions and is proportional to the following expression [5]:

$$\Gamma_\mu \propto |c_{LL}|^2 + |c_{LR}|^2 + |c_{RL}|^2 + |c_{RR}|^2 = |c_{LL}|^2(1 + |\bar{c}_{LR}|^2 + |\bar{c}_{RL}|^2 + |\bar{c}_{RR}|^2), \quad (34)$$

where $\bar{c}_{ij} = c_{ij}/c_{LL}$.

If we define $|V_{ud}|^2_{\rm expt}$ as being the quantity obtained from the ratio of measured β- and muon-decay rates, we can combine Eqs. (31) and (34) to relate this experimental result to the matrix element $|V_{ud}^L|^2$ by the relationship

$$|V_{ud}|^2_{\rm expt} \equiv \frac{\Gamma_\beta}{\Gamma_\mu} = |V_{ud}^L|^2 \frac{|1+\bar{a}_{LR}|^2 + |\bar{a}_{RL}+\bar{a}_{RR}|^2}{1+|\bar{c}_{LR}|^2+|\bar{c}_{RL}|^2+|\bar{c}_{RR}|^2} \simeq |V_{ud}^L|^2(1 + 2\,\mathrm{Re}\,\bar{a}_{LR}), \quad (35)$$

where, in the second line, only corrections to first order in small quantities are retained. If the situation is identical for the second (kaon decay) and third (B-meson decay) generations of quarks, with the interaction coupling constants a_{ij} and c_{ij} in $H_{s\ell}$ and H_ℓ being generation independent, then

$$\sum_i |V_{ui}|^2_{\rm expt} = \sum_i |V_{ui}^L|^2(1 + 2\,\mathrm{Re}\,\bar{a}_{LR}) = 1 + 2\,\mathrm{Re}\,\bar{a}_{LR}. \quad (36)$$

In writing the second line we have assumed that the CKM matrix for left-hand chirality quarks is strictly unitary. Since the left-hand side of Eq. (36) is the experimentally determined unitarity sum, given in Eq. (21) of Sec. IV B, this expression can clearly be used to set a limit on the coupling constant $\bar{a}_{LR}$. The result is

$$0.99995 \pm 0.00061 = 1 + 2\,\mathrm{Re}\,\bar{a}_{LR}, \quad \mathrm{Re}\,\bar{a}_{LR} = -0.00003 \pm 0.00030, \quad (37)$$

which is consistent with no right-hand currents—at least not in the LR sector.

2. Extra Z bosons

The existence of neutral gauge bosons, beyond the usual photon and Z boson of the standard $\mathrm{SU}(2)_L \times \mathrm{U}(1)$ model, would impact on the CKM unitarity test. To illustrate this we consider just one of the many models that appear in grand unified theories, namely the SO(10) model, whose group breakdown is

$$\mathrm{SO}(10) \rightarrow \mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1) \times \mathrm{U}(1)_\chi. \quad (38)$$

Here an extra U(1) group is introduced, $\mathrm{U}(1)_\chi$, and its corresponding neutral gauge boson is labeled Z_χ. The existence of such an extra Z boson would impact on the calculation of the electroweak radiative correction. One of the important diagrams in the hadron-independent radiative correction, Δ_R^V, is a WZ-box graph. This graph would have to be augmented by an additional WZ_χ-box graph, whose contribution is of order $\ln x_\chi$, where $x_\chi = m_{Z_\chi}^2/m_W^2$, the ratio of squared masses of the heavy bosons in the box diagram. If we assume that this correction is common to all quark flavours, then the same correction that occurs in the determination of $|V_{ud}|^2$ would also occur for $|V_{us}|^2$ and $|V_{ub}|^2$. If so, its impact can be incorporated into the unitarity test.

Following Marciano and Sirlin [212], we write

$$|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = (0.99995 \pm 0.00061) + \Delta, \quad (39)$$

where the numerical value for the experimental unitarity sum is from Eq. (21), and Δ is a calculated correction due to the extra Z boson. If we take the CKM matrix to be exactly unitary in three generations, then Eq. (39) can be used to set the following one-standard-deviation limits on Δ:

$$-0.00056 \leqslant \Delta \leqslant +0.00066. \quad (40)$$

Marciano and Sirlin [212] have computed the contribution of a putative Z_χ boson to the radiative correction and obtained

$$\Delta = -\frac{27\alpha}{40\pi \sin^2\theta_W} \times \frac{4}{3}|C_\chi|^2 \frac{\ln x_\chi}{x_\chi - 1}, \quad (41)$$

where α is the fine-structure constant, θ_W is the Weinberg angle ($\sin^2\theta_W \simeq 0.23$), and C_χ is a coupling constant linking the Z_χ boson to fermions. The normalization has been selected so that C_χ is unity at the SO(10) unification mass scale. Its value at lower energies has to be estimated, and Marciano and Sirlin use $|C_\chi|^2 = \frac{1}{2}$. Noting that the correction Δ is negative, we obtain from the lower limit in Eq. (40)

$$\frac{\ln x_\chi}{x_\chi - 1} \leqslant 0.12. \quad (42)$$

Taking $m_W = 81$ GeV for the W-boson mass, we arrive at the limit

$$m_{Z_\chi} > 430\ \mathrm{GeV}. \quad (43)$$

Impressive though this limit is, somewhat higher limits have been obtained in direct searches at proton and electron colliders. The CDF and D0 experiments at FermiLab in searches of $\bar{p}p \rightarrow e^+e^-$ have placed lower mass limits (at 95% C.L.) on m_{Z_χ} of 822 and 640 GeV, respectively, whereas at CERN the LEP2 experiment on $e^+e^- \rightarrow f\bar{f}$ (with f signifying a fermion) find a lower mass limit of 673 GeV. These limits are recorded in the survey of Erler and Langacker [213] in the 2008 PDG listings.

V. CONCLUSIONS

In our previous survey [5], only four years ago, we remarked on the excellent agreement among the derived $\mathcal{F}t$ values, lamented that the results of the unitarity test were still ambiguous, and predicted that the already well-measured ft values of the "traditional nine" superallowed decays were unlikely to be improved dramatically in the near future. Much has happened since then, not all of it expected. Today, we can say that the excellent $\mathcal{F}t$ value consistency remains—or, to be more accurate, it has been restored after Penning-trap $Q_{\rm EC}$ value measurements, nonexistant at the time of the last survey, did in fact make important improvements (and changes) in the known ft values, which in turn prompted improvements (and changes) in the calculated isospin-symmetry-breaking corrections. At the same time, the calculation of the nucleus-independent radiative correction, Δ_R^V, was improved, leading to a more precise result for V_{ud}, and the kaon-decay community mounted a concerted effort, which led to a new and reliable value for V_{us}. With these new results, and others, CKM unitarity has now been tested to unprecedented precision . . . and it has passed the test with flying colors.

Furthermore, we have demonstrated in Sec. IV how powerful these improved results can be in setting limits on new physics beyond the standard model, whether that new physics be a scalar interaction, right-hand currents, or extra Z bosons. We have seen that tiny uncertainties on the ft values are essential ingredients of a demanding test of CKM unitarity, which also leads to tight limits on new physics. The challenge now is whether those uncertainties can be reduced still further? The motivation is as strong as ever: to identify the need for new physics—or to limit the possible candidate theories even more definitively.

We have taken pains throughout this work to pay careful attention to all uncertainties, both theoretical and experimental. In Sec. IV A we detail the various contributions to the uncertainty in $|V_{ud}|^2$. Of these, by far the largest is still from Δ_R^V, even though its uncertainty has recently been improved significantly [6]. To improve it more must remain an important theoretical goal.

The next largest contributor to the error budget for $|V_{ud}|^2$ is the nuclear-structure-dependent corrections $(\delta_C - \delta_{\rm NS})$. Their uncertainties arise both from the input parameters used in their calculation—two-body matrix elements in the shell-model calculations, experimental uncertainties in charge-radii, etc. [7,179]—and from possible systematic differences between two different methods used for calculating radial wave functions (see Sec. III C). From a theoretical point of view, it would obviously be desirable to have a third completely different calculation, to reinforce the assessment of systematic uncertainties. However, in the absence of such a calculation, one must rely on experiment to test the accuracy of these calculated corrections. This has become, and should remain, a top priority for experiment.

The method, which is best described with reference to Fig. 9, is based on the validity of the CVC hypothesis that the corrected $\mathcal{F}t$ values for the superallowed $0^+ \rightarrow 0^+$ decays should be constant. In the figure we compare the uncorrected measured ft values (points and error bars) with the quantity $\overline{\mathcal{F}t}/[(1+\delta'_R)(1-\delta_C+\delta_{\rm NS})]$ shown as a band, the width of which represents the assigned theory error. The band corresponds to the calculated corrections normalized to the data via the measured average $\mathcal{F}t$ value, $\overline{\mathcal{F}t}$, taken from Table IX. Thus, although this comparison does not test the absolute values of the correction terms, it does test the collective ability of all three calculated correction terms to reproduce the significant variations in ft from one transition to another. In fact, since δ'_R is almost independent of Z when $Z > 10$, this test really probes directly the effectiveness of the calculated values of $(\delta_C - \delta_{\rm NS})$.

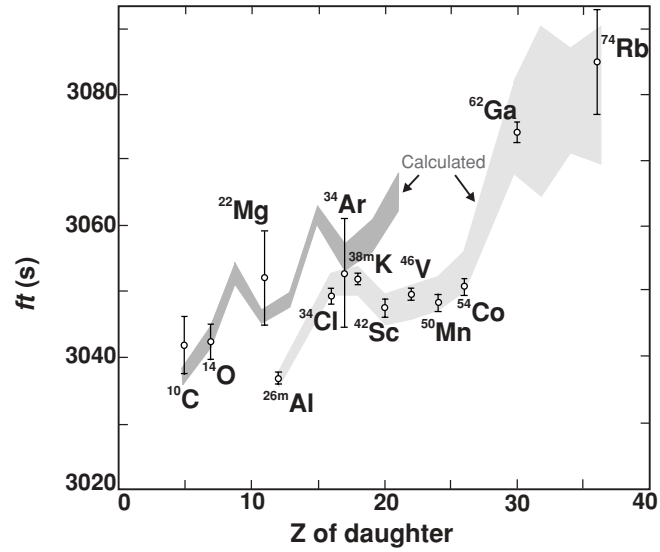

FIG. 9. Experimental ft values plotted as a function of the charge on the daughter nucleus, Z. Both bands represent the quantity $\overline{\mathcal{F}t}/[(1+\delta'_R)(1-\delta_C+\delta_{\rm NS})]$. The two separate bands distinguish those β emitters whose parent nuclei have isospin $T_z = -1$ (darker shading) from those with $T_z = 0$ (lighter shading).

It can be seen that there is remarkable agreement between theory and experiment. In assessing the significance of this agreement, it is important to recognize that the calculations of δ_C and $\delta_{\rm NS}$ for $Z \leqslant 26$ are based on well-established shell-model wave functions that were further tuned to reproduce measured binding energies, charge radii, and coefficients of the isobaric multiplet mass equation [7,179]. The origins of the calculated correction terms for all cases are completely independent of the superallowed decay data. Thus, the agreement in the figure between the measured superallowed data points and the theoretical band is already a powerful validation of the calculated corrections used in determining that band. The validation becomes even more convincing when we consider that it would require a pathological fault indeed in the theory to allow the observed nucleus-to-nucleus variations in δ_C and $\delta_{\rm NS}$ to be reproduced in such detail while failing to obtain the *absolute* values to comparable precision. As satisfactory as the agreement in Fig. 9 is, though, new experiments can still improve the test, making it even more demanding, and can ultimately serve to reduce the uncertainty in the nuclear-structure-dependent corrections even further.

These new experiments can follow different paths. In the past four years, the biggest impact has come from experiments that focused on the "traditional nine" superallowed transitions. New Penning-trap Q_{EC} value measurements have already been mentioned, but there have been new half-life and branching-ratio measurements as well (see Tables III and IV). More improvements are still possible, as a glance at Fig. 4 reveals. If we accept as a goal that experiment should be more than a factor of 2 more precise than theory, then we see that the Q_{EC} values for ^{10}C, ^{14}O, and ^{34}Cl, the half-lives of ^{26}Al^m,^{34}Cl, ^{42}Sc, and ^{50}Mn, and the branching ratios for ^{10}C and ^{14}O can all bear improvement. It is also particularly noteworthy that any improvements in the cases of ^{10}C and ^{14}O will lead directly to improvements in the limits on the possible existence of scalar currents. As is evident from Fig. 7 and the discussion in Sec. IV C1, it is on these two low-Z superallowed transitions that a scalar current would have the largest effect. Unfortunately, the branching ratios for both these transitions offer experimental obstacles that have proved very difficult to surmount.

A second experimental path is to expand the number of precisely measured superallowed emitters to include cases for which the calculated nuclear-structure-dependent corrections are larger, or show larger variations from nuclide to nuclide, than the values applied to the "traditional nine" cases. We argue that if the experimental ft values agree with the calculations where the nucleus-to-nucleus variations are large, then that must surely verify the calculations' reliability for the nine cases whose corrections are considerably smaller. Already four cases of this type have been carefully measured: ^{22}Mg, ^{34}Ar, ^{62}Ga, and ^{74}Rb. They appear to agree well with the calculations although, with the exception of ^{62}Ga, their uncertainties are still five times greater than those for the best known transitions. Undoubtedly, these uncertainties will be reduced and more cases added in the near future.

These new cases certainly present serious experimental challenges. The parent nuclei are more exotic than the traditional cases, which all have stable daughters, so they are more difficult to produce in pure and statistically significant quantities. They also exhibit more complex branching patterns: Each $T_Z = -1$ parent nucleus decays by Gamow-Teller transitions of comparable strength to the superallowed Fermi one, thus requiring the latter's branching ratio to be measured directly with high precision. For the $T_Z = 0$ parents with $A \geqslant 62$, each decay includes numerous weak Gamow-Teller transitions, which are very difficult to observe individually but which collectively constitute nonnegligible branching strength. In both regions, these problems are being, or have been, overcome, albeit with very specialized techniques. The recently published branching-ratio measurement [64] for ^{62}Ga is an example of how even meticulously detailed spectroscopic studies must be combined with theory [88] to ensure that missing transitions are properly accounted for in the decays of the heavy $T_Z = 0$ parents.

There is a further important issue that arises for the superallowed emitters with $A \geqslant 62$: The shell-model calculations of the structure-dependent corrections for these nuclei are not solidly based on spectroscopic measurements as they are for the lighter nuclei. Such measurements simply do not exist for most $N \simeq Z$ nuclei in this mass region. Furthermore, charge radii and coefficients for the isobaric multiplet mass equation are not known either and so cannot be used to constrain the radial wave functions or "tune" the charge dependence embedded in the two-body matrix elements. As a consequence, the uncertainties assigned to the calculated corrections are very large (see the broad band in this mass region in Fig. 9), considerably reducing the usefulness of these nuclei either in testing the corrections or in contributing to the determination of V_{ud}. It would be very valuable in this context for radioactive-beam facilities to direct some attention to determining a wide variety of spectroscopic information in this mass region with a view to obtaining a reasonably effective nuclear model, which, among other things, could lead to much improved calculations for the correction terms.

In conclusion, we can assert—as we did four years ago—that world data for superallowed $0^+ \rightarrow 0^+$ β decays strongly support the CVC expectation of an unrenormalized vector coupling constant and also set a tight limit, consistent with zero, on scalar currents. We can now add, though, that CKM unitarity is satisfied to within an uncertainty of 0.06%. This reconciliation with unitarity has come about as a result of significant changes in V_{us}; The value of V_{ud} determined from nuclear β decay has not varied outside of error bars in 20 years, during which time the size of those error bars has been reduced by a factor of 5. Finally, we have noted that the calculated nuclear-structure-dependent correction terms have recently been improved and continue to stand up favorably to experimental tests, an outcome that must further increase confidence in the nuclear results.

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APPENDIX A: ATOMIC OVERLAP CORRECTION TO THE STATISTICAL RATE FUNCTION

The statistical rate function, f, is an integral over phase space,

$$f = \int_1^{W_0} pW(W_0 - W)^2 F(Z, W) S(Z, W)\, dW, \quad \text{(A1)}$$

where W is the total energy of the electron in electron-rest-mass units, W_0 is the maximum value of W, $p = (W^2 - 1)^{1/2}$ is the momentum of the electron, Z is the atomic number of the daughter nucleus, $F(Z, W)$ is the Fermi function, and $S(Z, W)$

is the shape-correction function. The details of the calculation of $S(Z, W)$ were given in our previous survey [5] and will not be repeated here. What we address here is the inclusion for the first time of an additional factor in Eq. (A1) to account for the mismatch in the initial and final *atomic* states in the β decay. Since the nucleus changes charge by one unit in β decay, the final atomic state does not overlap perfectly with the initial atomic state, an effect that leads to a slight inhibition in the β-decay rate. In the past, this effect has justifiably been considered too small to be of practical concern but, with the advent of Penning-trap mass measurements, the experimental uncertainties in transition Q values have been reduced so much that they are now comparable to the effects of the imperfect atomic overlap.

We begin by writing

$$f = \int_1^{W_0} pW(W_0 - W)^2 F(Z, W) S(Z, W) r(Z, W)\, dW, \tag{A2}$$

where $r(Z, W)$ is the atomic overlap correction we are seeking. We then follow the method of Bahcall [214] by expressing f as a double integral with an energy-conserving delta function:

$$f = \iint pWq^2 F(Z, W) S(Z, W) \times \sum_{A'} |\langle A'|G\rangle|^2 \delta(E_f - E_i)\, dW\, dq, \tag{A3}$$

where q is the neutrino momentum. We have introduced into this equation an overlap of the initial and final atomic electron configurations: $|G\rangle$ is the state vector for the initial neutral atom with $(Z + 1)$ electrons, and $|A'\rangle$ is the state vector for the final *ionized* atom with $(Z + 1)$ electrons but only charge Z in the nucleus. There are many such possible final states, so a sum over A' is included.

Of the two energies within the delta function, the first, E_i, is the energy of the initial neutral atom in its atomic ground state:

$$E_i = \mathcal{M}_{Z+1}(G) = M_{Z+1} + (Z + 1)m_e - B(G), \tag{A4}$$

where $\mathcal{M}_{Z+1}(G)$ is the *atomic* mass, M_{Z+1} is the *nuclear* mass, m_e is the electron mass ($m_e = 1$ in electron-rest-mass units), and $B(G)$ is the total electron binding energy in the ground state of the atom. The sign of the latter is chosen so that $B(G) > 0$. The second energy, E_f, is that of the final state, which is composed of an *ionized* atom still with $(Z + 1)$ atomic electrons in an excited configuration plus an emitted β-decay positron and an emitted neutrino:

$$E_f = \mathcal{M}^{-1}_{Z+1}(A') + W + q = M_Z + (Z + 1)m_e - B(A') + W + q, \tag{A5}$$

where $\mathcal{M}^{-1}_{Z+1}(A')$ is the atomic mass of a negatively ionized atom (with superscript -1 denoting ionization) of $(Z + 1)$ electrons in configuration A', M_Z is the nuclear mass for the final nucleus in the β decay, and $B(A')$ is the total electron binding energy for the ionized atom. Thus the energy difference becomes

$$\begin{aligned} E_f - E_i &= (M_Z - M_{Z+1}) + W + q + B(G) - B(A') \\ &= (M_Z - M_{Z+1}) + W + q + [B(G) - B(G')] \\ &\quad - [B(A') - B(G')]. \end{aligned} \tag{A6}$$

In the second line of the equation we have introduced the total electron binding energy for the final *neutral* atom of charge Z in its atomic ground-state configuration, $B(G')$.

The $Q_{\rm EC}$ value is the difference in the atomic masses of neutral atoms in ground-state configurations:

$$\begin{aligned} Q_{\rm EC} &= \mathcal{M}_{Z+1}(G) - \mathcal{M}_Z(G') \\ &= [M_{Z+1} + (Z + 1)m_e - B(G)] - [M_Z + Zm_e - B(G')] \\ &= (M_{Z+1} - M_Z) + m_e + [B(G') - B(G)], \end{aligned} \tag{A7}$$

and the quantity W_0 in Eq. (A1) is related to $Q_{\rm EC}$ by the equation $W_0 = Q_{\rm EC} - m_e$. Thus, $E_f - E_i$ in Eq. (A6) can be written as

$$E_f - E_i = q + W - W_0 + [B(G') - B(A')]. \tag{A8}$$

For the energy-conserving delta function we now make a Taylor series expansion about the value $q + W - W_0$:

$$\delta(E_f - E_i) = \delta(q + W - W_0) + \delta'(q + W - W_0) \times [B(G') - B(A')] + \cdots. \tag{A9}$$

If the first term in this expansion is inserted into the double integral, Eq. (A3), then the expression for f reduces to the original form of Eq. (A1) since the atomic overlap factor is unity under the assumption that the sum over electronic configurations A' can be completed by closure; that is, $\sum_{A'} |\langle A'|G\rangle|^2 = \sum_{A'} \langle G|A'\rangle\langle A'|G\rangle = \langle G|G\rangle = 1$. The second term in Eq. (A9) involves a derivative of a delta function. This is handled by an integration by parts, in which the rest of the integrand is differentiated with respect to q. No boundary terms survive as the integrand vanishes at the boundaries. Thus the atomic overlap correction becomes

$$\begin{aligned} r(Z, W) &= 1 - \frac{2}{W_0 - W} \sum_{A'} |\langle A'|G\rangle|^2 [B(G') - B(A')] \\ &= 1 - \frac{2}{W_0 - W} \left(B(G') - \sum_{A'} |\langle A'|G\rangle|^2 B(A') \right). \end{aligned} \tag{A10}$$

Next, it is useful to recall the eigenvalue equations satisfied by the atomic states $|G\rangle$ and $|A'\rangle$:

$$\hat{H}_i|G\rangle = -B(G)|G\rangle \quad \text{with} \tag{A11}$$
$$\hat{H}_i = \sum_{i=1}^{Z+1} t_i - (Z + 1)e^2 \sum_{i=1}^{Z+1} \frac{1}{r_i} + e^2 \sum_{i<j=1}^{Z+1} \frac{1}{r_{ij}},$$

$$\hat{H}_f^{-1}|A'\rangle = -B(A')|A'\rangle \quad \text{with} \tag{A12}$$
$$\hat{H}_f^{-1} = \sum_{i=1}^{Z+1} t_i - Ze^2 \sum_{i=1}^{Z+1} \frac{1}{r_i} + e^2 \sum_{i<j=1}^{Z+1} \frac{1}{r_{ij}},$$

TABLE X. Comparison of statistical rate functions calculated without the atomic overlap correction, f_{without}, to those calculated with it included, f_{with}. The change in the Q_{EC} value that would lead to the same change in f is given in the last column.

Parent	f_{without}	f_{with}	df/f (%)	dQ/Q (%)	dQ (eV)
$T_z = -1$					
^{10}C	2.30089	2.30039	0.02178	0.00436	83
^{14}O	42.7779	42.7724	0.01277	0.00255	72
^{18}Ne	134.484	134.469	0.01093	0.00219	74
^{22}Mg	418.423	418.386	0.00877	0.00175	72
^{26}Si	1029.52	1029.44	0.00767	0.00153	74
^{30}S	1967.05	1966.91	0.00707	0.00141	77
^{34}Ar	3414.68	3414.46	0.00647	0.00129	78
^{38}Ca	5327.57	5327.24	0.00612	0.00122	81
^{42}Ti	7040.63	7040.21	0.00597	0.00119	84
$T_z = 0$					
^{26}Al^m	478.279	478.237	0.00880	0.00176	75
^{34}Cl	1996.10	1995.96	0.00711	0.00142	78
^{38}K^m	3298.10	3297.88	0.00663	0.00133	80
^{42}Sc	4472.52	4472.24	0.00643	0.00129	83
^{46}V	7209.90	7209.47	0.00598	0.00120	84
^{50}Mn	10746.6	10746.0	0.00565	0.00113	86
^{54}Co	15767.5	15766.6	0.00537	0.00107	89
^{62}Ga	26401.6	26400.2	0.00557	0.00111	102
^{66}As	32127.0	32125.3	0.00545	0.00109	104
^{70}Br	38602.2	38600.1	0.00539	0.00108	107
^{74}Rb	47296.9	47294.5	0.00523	0.00105	109

where t_i is the kinetic energy of electron i, r_i is its distance from the nucleus, and r_{ij} is the separation of electrons i and j. Note, in particular, that

$$\hat{H}_i - \hat{H}_f^{-1} = -e^2 \sum_{i=1}^{Z+1} \frac{1}{r_i} \simeq \frac{\partial}{\partial Z} \hat{H}_i . \tag{A13}$$

Inserting these Hamiltonian expressions into Eq. (A10) we obtain

$$\begin{aligned} r(Z, W) &= 1 - \frac{2}{W_0 - W} \left(B(G') + \sum_{A'} \langle G|\hat{H}_f^{-1}|A'\rangle\langle A'|G\rangle \right) \\ &= 1 - \frac{2}{W_0 - W} \left(B(G') + \langle G|\hat{H}_f^{-1}|G\rangle \right) \\ &= 1 - \frac{2}{W_0 - W} \big(B(G') - \langle G|\hat{H}_i - \hat{H}_f^{-1}|G\rangle \\ &\quad + \langle G|\hat{H}_i|G\rangle \big) \\ &= 1 - \frac{2}{W_0 - W} \left(B(G') + \frac{\partial}{\partial Z} B(G) - B(G) \right), \end{aligned} \tag{A14}$$

where in the last line the order of integration and differentiation has been reversed on the assumption that the binding energy as a function of Z behaves in a smooth way. Now, $B(G')$ is the electronic binding energy of a neutral atom with Z electrons, whereas $B(G)$ is the same quantity for an atom with $Z + 1$ electrons. Treating $B(G')$ as a function of Z, we can expand $B(G')$ in a Taylor series about $B(G)$:

$$B(G') = B(G) - \frac{\partial}{\partial Z} B(G) + \frac{1}{2} \frac{\partial^2}{\partial Z^2} B(G) - \cdots . \tag{A15}$$

Then, substituting this expression in Eq. (A14), we obtain our final expression for the atomic overlap correction:

$$r(Z, W) = 1 - \frac{1}{W_0 - W} \frac{\partial^2}{\partial Z^2} B(G). \tag{A16}$$

This expression was first obtained by Bahcall [214].

It remains for us to estimate the second derivative of the electronic binding energy of neutral atoms in their ground-state configuration. For this we use binding-energy values from the tables of Carlson *et al.* [215], which were obtained from self-consistent Hartree-Fock calculations and have been demonstrated to agree with experimental values to within 5%. We performed a fit to these tabulated values using a fitting function, aZ^b, in three ranges of Z values, with the following results:

$$B(G) = \begin{cases} 13.080 Z_i^{2.42}\ \text{eV}, & 6 \leqslant Z_i \leqslant 10, \\ 14.945 Z_i^{2.37}\ \text{eV}, & 11 \leqslant Z_i \leqslant 30, \\ 11.435 Z_i^{2.45}\ \text{eV}, & 31 \leqslant Z_i \leqslant 39, \end{cases} \tag{A17}$$

where Z_i is the charge of the parent atom in the β-decay process. It is conventional to use Z as the charge of the daughter nucleus in β decay; thus for positron decay $Z_i = Z + 1$. The second derivative is easily obtained from these expressions.

We have recomputed the statistical rate function f, with the results being listed in Table X. Those results obtained without

the atomic overlap correction, Eq. (A1), are given under the heading f_{without}, whereas those with the correction, Eqs. (A2) and (A16), are labeled f_{with}. The latter results also appear in column 2 of Table IX. The fractional difference between f_{with} and f_{without} in percent is given in column 4 and is of order 0.01%, decreasing with increasing mass value. This is a very small correction. Furthermore, the statistical rate function depends on the Q value to the fifth power, so the fractional change in Q that would lead to a change in f of the same size as that induced by the atomic overlap correction is even smaller: $1/5 \times df/f$. This percentage change is given in column 5 of Table X. As small as this effect is, it can be seen from the last column of the table that the equivalent change in Q value ranges from 70 to 110 eV, an amount that is similar to the experimental uncertainties on the most precisely measured Q values.

APPENDIX B: ACCOUNTING FOR SYSTEMATIC ERRORS

1. Isospin-symmetry-breaking correction, δ_{C2}

In the past (see, for example, Ref. [5]), we have added a systematic uncertainty to the average corrected $\overline{\mathcal{F}t}$ value to account for an apparent systematic difference between calculations of the isospin-symmetry-breaking correction, δ_{C2}, that used, on the one hand, Saxon-Woods eigenfunctions and, on the other hand, Hartree-Fock eigenfunctions. The former method, which was the one used by us [7,178,179], gave consistently larger corrections than the latter method, which was used by Ormand and Brown [181–183]. This was deemed to be a systematic effect caused by the different shapes of the Saxon-Woods and Hartree-Fock mean-field potentials. In this section, we readdress this issue for two reasons: (1) the Hartree-Fock calculations of Ormand and Brown are now 14–24 years old and are not available for all the nuclei under study here and (2) our most recent Saxon-Woods calculations [7] were performed in larger shell-model spaces not matched by the Hartree-Fock calculations. For a valid comparison both sets of calculations should be done in identical shell-model spaces.

The isospin-symmetry-breaking correction, δ_{C2}, for a superallowed β transition between 0^+, $T = 1$ analog states in nuclei with A nucleons is computed from the formula [7]

$$\delta_{C2} \simeq \sum_{\pi^<,\alpha} S_\alpha^< \Omega_\alpha^< - \frac{1}{2} \sum_{\pi^>,\alpha} S_\alpha^> \Omega_\alpha^>. \tag{B1}$$

Here S_α is the spectroscopic factor for the pickup of a nucleon in quantum state α from an A-particle state of spin 0^+ and isospin 1 to an $(A-1)$-particle state of spin α and isospin T_π. There is a complete-set sum over all the $(A-1)$-particle states (called parent states and denoted π) in Eq. (B1). The sum is divided into two parts: The first is over states with isospin $T_\pi = 1/2$ and is denoted by $\pi^<$; the second is over states with $T_\pi = 3/2$, denoted $\pi^>$. Further, the Ω_α^π are radial-mismatch factors, which depend on the difference between the radial wave function of a proton bound in the decaying nucleus, $u_{p,\alpha}^\pi(r)$, and that of a neutron bound in the daughter nucleus, $u_{n,\alpha}^\pi(r)$. Specifically, the radial-mismatch factors are given by

$$\begin{aligned}\Omega_\alpha^\pi &= \int_0^\infty u_{n,\alpha}^\pi(r)\left[u_{n,\alpha}^\pi(r) - u_{p,\alpha}^\pi(r)\right]dr \\ &= 1 - \int_0^\infty u_{n,\alpha}^\pi(r) u_{p,\alpha}^\pi(r)\, dr.\end{aligned} \tag{B2}$$

The radial functions are normalized to $\int |u(r)|^2\, dr = 1$. They are labeled by the parent state π because their asymptotic forms are matched to their separation energies, which in turn depend on the parent state. For example, if the parent state is the ground state of the $(A-1)$ system, then the proton separation energy would be S_p and the neutron separation energy S_n, two quantities given in terms of atomic masses and found in any atomic mass table. If, however, π represents an excited state of the $(A-1)$ system, then the proton and neutron separation energies would be $S_p + E_x$ and $S_n + E_x$, respectively, where E_x is the excitation energy of that parent state.

To compute δ_{C2} from Eq. (B1) one needs a set of spectroscopic factors, S_α, and a set of radial-mismatch integrals, Ω_α. The difference between Saxon-Woods and Hartree-Fock calculations lies in the method used to evaluate the radial-overlap integrals. In one case, the radial functions $u(r)$ are taken to be eigenfunctions of a Saxon-Woods potential; in the other case they are eigenfunctions of a Hartree-Fock mean-field potential. However, for a valid comparison between them the δ_{C2} must be calculated for both with the same spectroscopic factors, S_α. This was not done in the past, it being assumed that the model spaces were sufficiently comparable that this would not lead to any serious error. However, our most recent Saxon-Woods calculation [7] have significantly increased the model spaces over the ones used before, so now it is essential that a new Hartree-Fock calculation be undertaken. We report such a calculation here.

To be clear about the procedure, we illustrate it for the specific case of the decay of ^{34}Cl to ^{34}S. The decaying nucleus has $Z + 1 = 17$ protons; the daughter nucleus has $Z = 16$ protons. In the Saxon-Woods approach, the proton radial wave functions are taken to be eigenfunctions of a potential defined for a nucleus of mass A and charge $Z + 1$ as follows:

$$V(r) = -V_0 f(r) - V_s g(r) \mathbf{l}\cdot\boldsymbol{\sigma} + V_C(r), \tag{B3}$$

where

$$\begin{aligned} f(r) &= \{1 + \exp[(r-R)/a]\}^{-1}, \\ g(r) &= \left(\frac{\hbar}{m_\pi c}\right)^2 \frac{1}{a_s r} \exp\left(\frac{r-R_s}{a_s}\right) \\ &\quad \times \left\{1 + \exp\left(\frac{r-R_s}{a_s}\right)\right\}^{-2}, \\ V_C(r) &= Ze^2/r, \quad \text{for} \quad r \geqslant R_c \\ &= \frac{Ze^2}{2R_c}\left(3 - \frac{r^2}{R_c^2}\right), \quad \text{for} \quad r < R_c. \end{aligned} \tag{B4}$$

Here, $R = r_0(A-1)^{1/3}$ and $R_s = r_s(A-1)^{1/3}$. The three terms in Eq. (B3) are the central, spin-orbit, and Coulomb terms, respectively. In our calculations [7] most of the parameters were fixed at standard values, with the well depth V_0 being adjusted case by case so that the binding energy

of the eigenfunction being computed matched the separation energy to the corresponding parent state—in ^{33}S, for our example. Likewise the neutron radial functions were taken to be eigenfunctions of a similar potential but with the Coulomb term omitted.

Although the Hartree-Fock procedure is comparable, there is one issue unique to this approach that requires particular attention. For our illustrative example, a Hartree-Fock calculation might first be mounted for ^{34}Cl, which would yield a mean field with central, spin-orbit, and Coulomb terms. The required proton radial functions would then be taken as eigenfunctions of this mean field with the strength of the central term readjusted case by case so that the computed binding energy matched the appropriate separation energy. A second Hartree-Fock calculation might then be mounted for ^{34}S, from which the neutron radial functions would be similarly determined in the mean field, but without the Coulomb term. However, under these circumstances, if the Coulomb terms in the Hartree-Fock mean-field potential were to be compared with those in the Saxon-Woods potential, a very significant difference would emerge. In the Hartree-Fock case, the Coulomb term is

$$V_C(\mathbf{r}) = \int d^3\mathbf{r}' \frac{e^2}{|\mathbf{r}-\mathbf{r}'|}\rho_p(\mathbf{r}') - \frac{3e^2}{2}\left[\frac{3}{\pi}\rho_p(\mathbf{r})\right]^{1/3}, \quad \text{(B5)}$$

which depends on the proton density (of ^{34}Cl in our example) that is generated as part of the Hartree-Fock procedure. The two terms in Eq. (B5) are called the direct and exchange terms, respectively. If we take the asymptotic limit of the direct term for large r, we obtain

$$V_C^{\text{dir}}(\mathbf{r}) = \int d^3\mathbf{r}' \frac{e^2}{|\mathbf{r}-\mathbf{r}'|}\rho_p(\mathbf{r}') \xrightarrow{r\to\infty} \frac{e^2}{r}\int d^3\mathbf{r}'\rho_p(\mathbf{r}') = \frac{(Z+1)e^2}{r}. \quad \text{(B6)}$$

Since the Hartree-Fock proton density is normalized to $(Z+1)$ protons in ^{34}Cl, the asymptotic form of the Coulomb potential tends to $(Z+1)e^2/r$. However, this disagrees with the equivalent Saxon-Woods calculation, which has the form Ze^2/r [see Eq. (B4)].

This discrepancy is important and constitutes, in our opinion, a serious flaw in this Hartree-Fock calculation of the radial-mismatch factor. Since a proton removed from a nucleus of charge $Z+1$ leaves behind Z protons, its asymptotic interaction is with charge Z—as described by the Saxon-Woods potential—and not with charge $Z+1$. This deficiency in Hartree-Fock would be overcome in principle by the Coulomb exchange term. However, in Skyrme-Hartree-Fock calculations it is not possible to compute the exchange term exactly without sacrificing the simplicities that come with use of zero-range Skyrme interactions. The exchange term appearing in Eq. (B5) is a commonly used local approximation, which might well be appropriate for the nuclear interior and for the computation of bulk properties such as binding energies and radii, but it certainly does not do the job asymptotically, which is the region of greatest importance to our calculations.

To circumvent this difficulty, we have chosen to alter the Hartree-Fock protocol. Instead of mounting two Hartree-Fock calculations—for ^{34}Cl and ^{34}S—as just described, we mount a *single* calculation for the nucleus with $(A-1)$ nucleons and Z protons—^{33}S in our example. We then use the proton mean field from this calculation to generate the proton eigenfunctions $u^{\pi}_{p,\alpha}(r)$, and the neutron mean field from the same calculation to generate the neutron eigenfunctions $u^{\pi}_{n,\alpha}(r)$. In this procedure, the Coulomb interaction automatically has the correct asymptotic form. It is also fully consistent with the Saxon-Woods potential parametrization, Eq. (B4), which considers the nucleus of mass $(A-1)$ as the core to which the last particle is bound, since the radius of the potential is parametrized as $r_0(A-1)^{1/3}$ rather than $r_0A^{1/3}$. Calculations of δ_{C2} with this new Hartree-Fock protocol will be presented in the next section. It can be noted here, however, that these results are larger than those obtained with the conventional protocol by between 10% to 40% depending on the Skyrme interaction used and the nucleus under study. This change of protocol goes a long way in reducing the systematic error between Saxon-Woods and Hartree-Fock calculations.

2. New Hartree-Fock calculations for δ_{C2}

Here we present our new Hartree-Fock calculations for the isospin-symmetry-breaking correction, δ_{C2}, for the 20 cases of superallowed Fermi β decay considered in our survey: they range from ^{10}C to ^{74}Rb. Our procedure was that discussed at the end of the previous section and involved obtaining the mean field from a Hartree-Fock calculation in the $(A-1)$ system. The proton and neutron radial functions

TABLE XI. Adopted isospin-symmetry-breaking corrections, δ_{C2} in percent units, and their assigned uncertainties obtained from Hartree-Fock calculations. Also listed are earlier results obtained with Saxon-Woods (SW) eigenfunctions, as published in Ref. [7].

Nucleus	HF	SW
$T_z = -1$		
^{10}C	0.215(35)	0.165(15)
^{14}O	0.255(30)	0.275(15)
^{18}Ne	0.205(55)	0.410(25)
^{22}Mg	0.250(55)	0.370(20)
^{26}Si	0.335(55)	0.405(25)
^{30}S	0.540(55)	0.700(20)
^{34}Ar	0.510(60)	0.635(55)
^{38}Ca	0.600(60)	0.745(70)
^{42}Ti	0.535(60)	0.835(75)
$T_z = 0$		
^{26}Al	0.410(50)	0.280(15)
^{34}Cl	0.595(55)	0.550(45)
^{38}K	0.640(60)	0.550(55)
^{42}Sc	0.620(55)	0.645(55)
^{46}V	0.525(55)	0.545(55)
^{50}Mn	0.575(55)	0.610(50)
^{54}Co	0.635(55)	0.720(60)
^{62}Ga	0.93(16)	1.20(20)
^{66}As	1.10(35)	1.35(40)
^{70}Br	1.14(25)	1.25(25)
^{74}Rb	1.29(16)	1.50(30)

were obtained as eigenfunctions of this mean-field potential, whose overall strength was scaled on a case-by-case basis to ensure the eigenfunction's asymptotic solution matched the required separation energy. These eigenfunctions were used to compute the radial-mismatch factors, Ω_α^π, of Eq. (B2). For the spectroscopic factors needed in Eq. (B1), we ran several shell-model calculations with the model spaces and effective interactions used recently [7] in calculations with Saxon-Woods potentials. We also considered three choices of the Skyrme interaction: SGII [216], SkM* [217], and Ska [218]. The first two of these interactions were the ones used by Ormand and Brown [182,183] in their computations of δ_{C2}. The third interaction, Ska, is of similar quality and was one of the first to be fitted to the incompressibility of nuclear matter, a key constraint used in all later Skyrme interactions. More recent Skyrme interactions tend to be used in conjunction with pairing forces in Hartree-Fock-Bogliubov calculations [219–221]. Since we have not included pairing forces in the present work, we have not attempted to use any of these more recent interactions.

The results from our Hartree-Fock calculations for δ_{C2} are listed in column 2 of Table XI. For each transition, the central value is an average of the results obtained with the three choices of Skyrme interactions. To assign an uncertainty, we have examined the spread in results obtained with the different Skyrme interactions and with different shell-model effective interactions and model spaces. We also list in the third column of the table the values we adopted from our Saxon-Woods computations, which originally appeared in Ref. [7].

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