Low-energy parity-violation and new physics

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The new physics sensitivity of a variety of low-energy parity-violating (PV) observables is analyzed. A comparison is made between atomic PV for a single isotope, atomic PV using isotope ratios, and PV electron-hadron and electron-electron scattering. The complementarity among these observables, as well as with high-energy processes, is emphasized. Theoretical uncertainties entering the interpretation of low-energy measurements are discussed. [S0556-2813(99)03207-0]

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I. INTRODUCTION

Low-energy parity-violating (PV) observables have played an important role in uncovering the structure of the electroweak sector of the standard model. Now that the predictions of the standard model have been tested and confirmed at the one-loop level over a wide range of processes and energies [1], attention has turned to the search for physics beyond the standard model. In this regard, low-energy parity violation continues to provide important information. As has been noted by several authors [2–4], the recent precise determination of the cesium weak charge Q_W in an atomic parity-violation (APV) experiment performed by the Boulder group [5] places stringent constraints on a variety of new physics scenarios. The importance of this benchmark measurement is reflected, in part, by the efforts of other experimental groups to determine Q_W for cesium as well as other atoms [6–8]. Future improvements in the APV sensitivity to new physics poses a challenge to both atomic experimentalists and theorists. Indeed, given the experimental precision reported by the Boulder group, atomic theory error now constitutes the dominant uncertainty associated with the intepretation of atomic PV (APV) observables. Whether this atomic theory error can be reduced to the level of the experimental uncertainty remains to be seen. An experimental strategy for circumventing the atomic theory uncertainty is to measure PV observables for different atoms along an isotope chain. Standard model predictions for *ratios* of such observables are largely atomic theory independent. Consequently, several groups have undertaken APV isotope ratio measurements in the hopes of minimizing the impact of atomic theory uncertainties on the extraction of new physics constraints [7–9].

Historically, the use of polarized electrons produced in accelerator experiments has, along with APV, played a part in testing the standard model [10–12]. In the past decade, however, PV electron scattering (PVES) has received less attention than APV in this respect since (a) the experimental precision achievable with APV has improved markedly and (b) interest in PVES has focused on its use in probing the nucleon's $s\bar{s}$ sea. The interest in nucleon strangeness has spawned a program of experiments at MIT-Bates, Mainz, and the Jefferson Laboratory to measure the left-right asymmetry A_{LR} on a variety of targets [13–15]. Recently, the attention of the PVES community has returned to the use of these experiments to probe new physics [16]. In the purely leptonic sector, work on a high-precision PV Möller scattering experiment has begun at SLAC [17]. In addition, a program of ‘‘second generation’’ PVES experiments—designed to look for physics beyond the standard model—is under consideration for the Jefferson Lab. The feasibility of such PVES new physics searches stems, in part, from the high luminosity and remarkably stable and clean electron beam achieved by the CEBAF accelerator [18].

Although numerous discussions of Q_W(cesium) have appeared in the literature recently, relatively little attention has been paid to the other low-energy PV observables mentioned above. In this paper, we therefore consider the new physics sensitivities of APV isotope ratios and PVES asymmetries, making a comparison with the sensitivities of Q_W and high-energy observables. In doing so, we focus on ‘‘direct’’ new physics, that is, extensions of the standard model which manifest themselves at low-energies as new four-fermion contact interactions. The sensitivity of APV and PVES to ‘‘oblique’’ new physics has been discussed elsewhere [2,19–21,24].[1] After quantifying the generic new physics sensitivities of PV observables, we specify to a variety of models in order to illustrate the complementarity of prospective measurements. In particular, we show that to leading order, the elastic ep asymmetry $A_{LR}(^1\mathrm{H})$ and APV isotope ratios $\mathcal{R}$ are sensitive to the same combination of possible new interactions. These two observables, while subject to different systematic and theoretical corrections, provide the same window on direct new physics. We also find that a 2–3% determination of $A_{LR}(^1\mathrm{H})$ would improve the new physics reach of low-energy PV by nearly a factor of 2 over the present cesium APV sensitivity. A similar improvement would obtain if the present cesium atomic theory error were

[1] Previously, isotope ratios were shown to display a significantly different sensitivity to oblique new physics than does Q_W for a single isotope. We show that a similar situation holds in the case of direct new physics.

improved by a factor of 4. Apart from the Möller asymmetry, the remaining asymmetries display a smaller new physics reach than Q_W, $A_{LR}(^1\mathrm{H})$, or $\mathcal{R}$. In illustrating model variations on the general pattern of new physics sensitivity, we consider additional neutral gauge bosons, leptoquarks, and fermion compositeness. We also discuss the sensitivity of PV observables to R-parity-violating supersymmetric interactions and compare this sensitivity with up-dated bounds from superallowed β decay.

The use of low-energy measurements to probe new physics requires that conventional, many-body physics associated with atoms and nuclei be sufficiently well understood. From this standpoint, we show that, in principle, PVES provides the theoretically ''cleanest'' new physics probe. This feature is most apparent for PV Möller scattering, as it is a purely leptonic process. In the case of semileptonic PV observables, the reason for minimal theoretical uncertainty is twofold: (a) A_{LR} depends on a ratio of electroweak amplitudes, from which the largest hadronic effects cancel, leaving essentially a dependence on Q_W of the target nucleus and (b) the largest remaining hadronic corrections to this cancellation can be separated from Q_W and measured by exploiting their kinematic dependence. Consequently, the dominant uncertainty in the interpretation of PVES new physics studies is likely to be experimental. We illustrate these features in the case of $A_{LR}(^1\mathrm{H})$ and discuss the kinematics to make such a clean separation of Q_W feasible.

The situation in the case of APV differs from that of PVES. The atomic theory uncertainty associated with extracting Q_W from cesium APV is about four times larger than the experimental error. This situation has prompted the consideration of the isotope ratios $\mathcal{R}$, from which the dominant atomic theory uncertainties cancel. Unfortunately, $\mathcal{R}$ carries a problematic sensitivity to changes in the neutron distribution $\rho_n(r)$ from one nucleus to the next along an isotope change. Following on the earlier work of Refs. [25–27], we analyze the impact which the uncertainties in the neutron distribution $\rho_n(r)$ have on the extraction of new physics limits from $\mathcal{R}$. We consider several atoms presently under experimental consideration and quantify the level of ρ_n uncertainty acceptable in order for relevant new physics limits to be obtained from $\mathcal{R}$.

Our discussion of these issues is organized as follows. In Sec. II, we outline our conventions and definitions and in Sec. III we discuss general new physics sensitivities of low-energy PV observables. In Sec. IV we illustrate these sensitivities for different new physics scenarios. Section V contains an analysis of theoretical uncertainties. A discussion of kinematic considerations for a prospective PV elastic ep experiment is also included. In Sec. VI we summarize our conclusions.

II. NEW PHYSICS AND THE WEAK CHARGE

For each PV observable, the quantity of interest here is the weak charge Q_W of the nucleus (electron), which characterizes the strength of the electron axial vector × nucleus (electron) vector weak neutral current interaction:

$$Q_W = Q_W^0 + \Delta Q_W . \tag{1}$$

Here, Q_W^0 gives the contribution in the standard model while ΔQ_W indicates possible contributions from new interactions. We consider Q_W to be generated by the low-energy effective Lagrangian

$$\mathcal{L} = \mathcal{L}_{\mathrm{SM}}^{\mathrm{PV}} + \mathcal{L}_{\mathrm{new}}^{\mathrm{PV}}, \tag{2}$$

where

$$\mathcal{L}_{\mathrm{SM}}^{\mathrm{PV}} = \frac{G_F}{2\sqrt{2}} g_A^e \bar{e}\gamma_\mu\gamma_5 e \sum_f g_V^f \bar{f}\gamma^\mu f, \tag{3}$$

$$\mathcal{L}_{\mathrm{new}}^{\mathrm{PV}} = \frac{4\pi\kappa^2}{\Lambda^2} \bar{e}\gamma_\mu\gamma_5 e \sum_f h_V^f \bar{f}\gamma^\mu f. \tag{4}$$

Here $g_V^f = 2T_3^f - 4Q_f \sin^2\theta_W$ and $g_A^f = -2T_3^f$ are the tree level standard model fermion-Z^0 couplings, h_V^f characterizes the interaction of the electron axial vector current with the vector current of fermion f for a given extension of the standard model, Λ is the mass scale associated with the new physics, and κ sets the coupling strength. Generally speaking, strongly interacting theories take $\kappa^2 \sim 1$ while for weakly interacting extensions of the standard model one has $\kappa^2 \sim \alpha$. For scenarios in which the interaction of Eq. (4) is generated by the exchange of a new heavy particle between the electron and fermion, the constant $h_V^f = \tilde{g}_A^e \tilde{g}_V^f$, where $\tilde{g}_A^e(\tilde{g}_V^f)$ are the heavy particle axial vector (vector) coupling to the electron (fermion).

For simplicity, we do not consider contributions to ΔQ_W arising from new scalar-pseudoscalar or tensor-pseudotensor interactions. We also do not consider $V(e)\times A(f)$ interactions, as they do not contribute to Q_W. Although the standard model $V(e)\times A(f)$ interaction is suppressed due to the small value of $g_A^e = -1 + 4\sin^2\theta_W$, resulting in an enhanced sensitivity to new physics of this type, one is at present not able to extract the $V(e)\times A(f)$ amplitudes from PV observables with the level of experimental precision attainable for Q_W. Moreover, the hadronic axial vector current is not protected by current conservation from hadronic effects which may cloud the interpretation of the hadronic axial vector amplitude in terms of new physics [29].

It is straightforward to write down the corrections to the weak charge of a given system arising from $\mathcal{L}_{\mathrm{new}}^{\mathrm{PV}}$. Specifically, we consider the nucleon and electron

$$\Delta Q_W^P = \zeta(2h_V^u + h_V^d), \tag{5}$$

$$\Delta Q_W^N = \zeta(h_V^u + 2h_V^d), \tag{6}$$

$$\Delta Q_W^e = \zeta h_V^e, \tag{7}$$

where

$$\zeta = \frac{8\sqrt{2}\pi\kappa^2}{\Lambda^2 G_F}. \tag{8}$$

To the extent that the couplings g_V^f and h_V^f entering Q_W and ΔQ_W are of the same order of magnitude, the fractional correction induced by new physics is

$$\frac{\Delta Q_W}{Q_W^0}=\frac{8\sqrt{2}\pi\kappa^2}{\Lambda^2 G_F}. \tag{9}$$

A one percent determination of Q_W then affords a lower bound on the mass scale associated with new physics of

$$\Lambda \geqslant \left[\frac{8\sqrt{2}\pi\kappa^2}{0.01 G_F}\right]^{1/2} \approx 20\kappa \ \text{TeV}. \tag{10}$$

In short, determinations of Q_W at the one percent or better level probe new physics at the TeV scale for weakly interacting theories and the ten TeV scale for new strong interactions.

III. OBSERVABLES

In this section, we discuss some of the general features of the low-energy PV observables used to determine Q_W. In particular, we consider a general atomic PV observable for a single isotope $A_{\rm PV}(N)$, ratios involving $A_{\rm PV}$ for different isotopes $\mathcal{R}$, and the left-right asymmetry for scattering polarized electrons from a given target A_{LR}. Of these, the simplest is the atomic PV observable for a single isotope $A_{\rm PV}(N)$. The nuclear spin-independent (NSID) part of this observable is given by

$$A_{\rm PV}^{\rm NSID}(N)=\xi Q_W=\xi[Q_W^0+Z\Delta Q_W^P+N\Delta Q_W^N], \tag{11}$$

where ξ is an atomic structure-dependent coefficient and where

$$Q_W^0=Z(1-4\sin^2\theta_W)-N \tag{12}$$

at tree level and

$$Z\Delta Q_W^P+N\Delta Q_W^N=\zeta[(2Z+N)h_V^u+(2N+Z)h_V^d]. \tag{13}$$

A determination ξ generally requires theoretical knowledge of the relevant atomic wavefunction and, therefore, introduces theoretical uncertainty into the extraction of Q_W. The relative sensitivity of $A_{\rm PV}^{\rm NSID}(N)$ to new physics can be seen by rewriting Q_W as

$$Q_W=Q_W^0[1+\delta_N], \tag{14}$$

where

$$\begin{aligned}\delta_N&=(Z\Delta Q_W^P+N\Delta Q_W^N)/Q_W^0\\ &\approx-\zeta\left[\left(\frac{2Z+N}{N}\right)h_V^u+\left(\frac{2N+Z}{N}\right)h_V^d\right]\\ &=-\zeta[(Z/N)(2h_V^u+h_V^d)+(2h_V^d+h_V^u)],\end{aligned} \tag{15}$$

where the approximation $Q_W^0\approx -N$ has been made in light of the small value for $1-4\sin^2\theta_W\approx 0.1$. From Eq. (15) we observe that for atoms having $Z\approx N$, the weak charge is roughly equally sensitive to the new up- and down-quark vector current interactions.

The use of ‘‘isotope ratios’’ involving $A_{\rm PV}^{\rm NSID}(N)$ and $A_{\rm PV}^{\rm NSID}(N')$ largely eliminates the dependence on the atomic structure-dependent constant ξ and the associated atomic theory uncertainty. We consider two such ratios:

$$\mathcal{R}_1=\frac{A_{\rm PV}^{\rm NSID}(N')-A_{\rm PV}^{\rm NSID}(N)}{A_{\rm PV}^{\rm NSID}(N')+A_{\rm PV}^{\rm NSID}(N)} \tag{16}$$

and

$$\mathcal{R}_2=\frac{A_{\rm PV}^{\rm NSID}(N')}{A_{\rm PV}^{\rm NSID}(N)}. \tag{17}$$

To the extent that ξ does not vary appreciably along the isotope chain, one has

$$\mathcal{R}_1=\frac{Q_W(N')-Q_W(N)}{Q_W(N')+Q_W(N)}, \tag{18}$$

$$\mathcal{R}_2=\frac{Q_W(N')}{Q_W(N)}. \tag{19}$$

It is straightforward to work out the sensitivity of these ratios to new physics. To this end, we write

$$\mathcal{R}_i=\mathcal{R}_i^0(1+\delta_i), \tag{20}$$

where

$$\mathcal{R}_1^0=\frac{Q_W^0(N')-Q_W^0(N)}{Q_W^0(N')+Q_W^0(N)}, \tag{21}$$

$$\mathcal{R}_2^0=\frac{Q_W^0(N')}{Q_W^0(N)}, \tag{22}$$

give the ratios in the standard model and δ_i's give corrections arising from new physics. Letting $N'=N+\Delta N$ and dropping small contributions containing $1-4\sin^2\theta_W$ one has

$$\mathcal{R}_1^0\approx\frac{\Delta N}{2N}, \tag{23}$$

$$\mathcal{R}_2^0\approx 1+\frac{\Delta N}{N} \tag{24}$$

and

$$\delta_1\approx\zeta\left(\frac{2Z}{N+N'}\right)(2h_V^u+h_V^d), \tag{25}$$

$$\delta_2\approx\zeta\left(\frac{Z}{N}\right)\left(\frac{\Delta N}{N'}\right)(2h_V^u+h_V^d). \tag{26}$$

At first glance, the dependence of δ_i, $i=1,2$, on $\Delta Q_W^P=\zeta(2h_V^u+h_V^d)$ only and not on $\Delta Q_W^N=\zeta(2h_V^d+h_V^u)$ may seem puzzling. To first order in ζ, however, the shifts ΔQ_W^N appearing in the numerator and denominator of each $\mathcal{R}_i$ cancel. In the case of $\mathcal{R}_1$, for example, one has

$$Q_W(N')-Q_W(N)\approx -N'+N+(N'-N)\Delta Q_W^N \tag{27}$$

$$=(N-N')[1-\Delta Q_W^N] \tag{28}$$

and

$$Q_W(N')+Q_W(N)\approx -(N+N')+(N+N')\Delta Q_W^N+2Z\Delta Q_W^P \tag{29}$$

$$=-(N+N')\left[1-\Delta Q_W^N - \left(\frac{2Z}{N+N'}\right)\Delta Q_W^P\right] \tag{30}$$

so that in the ratio, the dependence on ΔQ_W^N cancels to first order. Hence, the $\mathcal{R}_i$'s are twice as sensitive to new physics involving u-quarks than to new physics which couples to d quarks. The weak charge of a single isotope, on the other hand, has essentially the same sensitivity to u- and d-quark new physics.

From a comparison of δ_N with the δ_i, we also observe that, for a given experimental precision, the isotope ratios are generally less sensitive to direct new physics than is the weak charge for a single isotope. This feature is particularly evident in the case of $\mathcal{R}_2$, since δ_2 contains the explicit factor $\Delta N/N'$. Taking $Z\approx N$ for the case of $\mathcal{R}_1$, we find that a single isotope is three times more sensitive to new physics which couples to d quarks and 1.5 times more sensitive to the u-quark coupling. For new physics scenarios which favor new e-d interactions over e-u interactions (e.g., E_6 models, discussed below), the weak charge for a single isotope consititutes a more sensitive probe.

An alternative method for obtaining Q_W is to scatter longitudinally polarized electrons from fixed targets. Flipping the incident electron helicity and comparing the helicity difference cross section with the total cross section filters out the PV part of the weak neutral current interaction. The resulting left-right asymmetry for elastic scattering has the general form [21–23]

$$A_{LR}=\frac{N_+-N_-}{N_++N_-}\approx\frac{2M_{\rm NC}^{\rm PV}}{M_{\rm EM}}=\frac{G_F|q^2|}{4\sqrt{2}\pi\alpha}\left[\frac{Q_W}{Q_{\rm EM}}+F(q)\right]. \tag{31}$$

Here, N_+ (N_-) are the number of detected electrons for a positive (negative) helicity incident beam, $M_{\rm EM}$ and $M_{\rm NC}^{\rm PV}$ are, respectively, the electromagnetic and parity-violating neutral current electron-nucleus scattering amplitudes, $Q_{\rm EM}$ is the nuclear EM charge, and $F(q)$ is a correction involving hadronic and nuclear form factors. In general, the latter term can be separated from the term containing the charges by varying electron energy and angle. For elastic scattering, the weak charge term can be isolated by going to forward angles and low energies. In the case of PV Möller scattering, one has $F(q)\equiv 0$. The present PV electron scattering program at MIT-Bates, Mainz-MAMI, and the Jefferson Laboratory seeks to determine the $F(q)$ for a variety of targets, with a special emphasis on contributions from strange quarks.

In order to compare the sensitivities of different scattering experiments to new physics, we specify the terms in Eq. (31) for the following processes: elastic scattering from the proton $A_{LR}(^1H)$, elastic scattering from $(J^\pi,T)=(0^+,0)$ nuclei $A_{LR}(0^+,0)$, excitation of the $\Delta(1232)$ resonance $A_{LR}(N\to\Delta)$, and Möller scattering $A_{LR}(e)$. The corresponding charge terms are (neglecting standard model radiative corrections)

$$Q_W(^1H)/Q_{\rm EM}(^1H)=(1-4\sin^2\theta_W)[1+\delta_P], \tag{32}$$

$$Q_W(0^+,0)/Q_{\rm EM}(0^+,0)=-4\sin^2\theta_W[1+\delta_{00}], \tag{33}$$

$$Q_W(e)/Q_{\rm EM}(e)=(-1+4\sin^2\theta_W)[1+\delta_e], \tag{34}$$

while for the $N\to\Delta$ transition one replaces the ratio of charges by the ratio of isovector weak neutral current and EM couplings

$$Q_W(N\to\Delta)/Q_{\rm EM}(N\to\Delta)\to 2(1-2\sin^2\theta_W)[1+\delta_\Delta]. \tag{35}$$

The new physics corrections δ are given by

$$\delta_P=\zeta(2h_V^u+h_V^d)/(1-4\sin^2\theta_W), \tag{36}$$

$$\delta_{00}=-3\zeta(h_V^u+h_V^d)/(4\sin^2\theta_W), \tag{37}$$

$$\delta_e=-\zeta h_V^e/(1-4\sin^2\theta_W), \tag{38}$$

$$\delta_\Delta=\zeta(h_V^u-h_V^d)/[2(1-2\sin^2\theta_W)]. \tag{39}$$

For completeness, we also write down the corresponding expressions for PV deep inelastic scattering (DIS). We consider only the case of deuterium, which was the target in the first PV scattering experiment and was proposed in the early 1990's as the target for a new SLAC experiment [30,31]. An analysis of new physics contributions to the PV DIS asymmetry requires that we consider the more general four fermion Lagrangian

$$\mathcal{L}^{\rm new}=\frac{4\pi\kappa^2}{\Lambda^2}\sum_{q,i,j} h_{ij}^q\bar{e}_i\gamma_\mu e_i\bar{q}_j\gamma^\mu q_j, \tag{40}$$

where i and j denote the handedness of the given fermion. h_V^q's of Eq. (4) represent one linear combination of h_{ij}^q's:

$$h_V^q=(h_{RR}^q-h_{LL}^q+h_{RL}^q-h_{LR}^q)/4. \tag{41}$$

The PV DIS asymmetry for a deuterium target is [21]

$$A_{LR}^{\rm DIS}(^2{\rm H})=\frac{G_F|q^2|}{4\sqrt{2}\pi\alpha}\frac{9}{5}\left\{\tilde{a}_1+\tilde{a}_2\left[\frac{1-(1-y)^2}{1+(1-y)^2}\right]\right\}, \tag{42}$$

where

$$\tilde{a}_1=\left(1-\frac{20}{9}\sin^2\theta_W\right)[1+\tilde{\delta}_1], \tag{43}$$

TABLE I. Relative sensitivities of PV observables to new physics, assuming a $h_V^u = h_V^d$, tree-level values for the corresponding weak charges (except for the Möller asymmetry, as noted in the text), and $\sin^2\hat{\theta}_W = 0.2314$. The scale factor $f_i = \sqrt{\delta_i/\delta_N}$ can be used to scale mass bounds from the cesium APV bounds to the bounds for observable i assuming the same precision for both δ_N and δ_i. Note that we have assumed $h_V^u = h_V^d$ so that $\delta_\Delta = 0$.

Correction δ	Scale factor f_i
$\delta_N \approx 5.1\zeta$	1
$\delta_1 \approx 1.9\zeta$	0.6
$\delta_2 \approx 0.4\zeta$	0.3
$\delta_P \approx 40\zeta$	2.8
$\delta_{00} \approx 6.5\zeta$	1.1
$\delta_e \approx 22\zeta$	2.1
$\delta_\Delta = 0$	0
$\tilde{\delta}_1 \approx \zeta/12$	0.13

$$\tilde{a}_2 = (1 - 4\sin^2\theta_W)[1 + \tilde{\delta}_2]. \tag{44}$$

The $\tilde{\delta}_i' s$ contain standard model radiative corrections, corrections involving the quark distribution functions [21], and contributions from new physics. Writing only the latter, we obtain

$$\tilde{\delta}_1 = \frac{\zeta}{24(1 - 20\sin^2\theta_W/9)}(2h_V^u - h_V^d) \approx \zeta(2h_V^u - h_V^d)/12, \tag{45}$$

$$\tilde{\delta}_2 = \frac{\zeta}{2(1 - 4\sin^2\theta_W)} \sum_q Q_q[h_{RR}^q - h_{LL}^q + h_{LR}^q - h_{RL}^q], \tag{46}$$

$$\approx 6.5\zeta \sum_q Q_q[h_{RR}^q - h_{LL}^q + h_{LR}^q - h_{RL}^q], \tag{47}$$

where Q_q is the quark EM charge.

The expressions for the various δ_i's allow us to make a few observations regarding the relative sensitivities the corresponding observables to new physics. For this purpose, we take $h_V^u = h_V^d = 1$ and specify δ_N for the case of ^{133}Cs. We also use cesium for the isotope ratios and take a reasonable range of neutron numbers $N = 75$, $N' = 95$ [32]. In Table I we show the δ_i in units of ζ. The third column gives a scale factor f defined as

$$f_i = \sqrt{\delta_i/\delta_N}. \tag{48}$$

The factor f_i can be used to scale the cesium APV sensitivity to the new physics mass scale Λ to those obtainable from any other observable when measured with the same precision as $Q_W(\mathrm{Cs})$: $\Lambda(i) = f_i \Lambda(\mathrm{Cs})$. Alternatively, the sensitivity of any other observable will be the same as that of cesium when the precision is f_i^2 times the cesium uncertainty. The numbers shown in the table are obtained using the $\overline{\mathrm{MS}}$ value $\sin^2\theta_W = 0.2314$ [33] in tree-level expressions for the weak charges. The entries for the Möller asymmetry have been modified to account for one-loop electroweak radiative corrections, according to the calculation of Ref. [34]. In the latter case, these corrections reduce the asymmetry by ≈40% from its tree-level value. Radiative corrections do not appreciably alter the relative new physics sensitivities of the other observables listed in Table I.

As Table I illustrates, $A_{LR}(^1\mathrm{H})$ and $A_{LR}(e)$ display the greatest sensitivities to new physics for a given level of error in the observables. The reason is the suppression of Q_W^0 for the proton and electron, which goes as $(1 - 4\sin^2\theta_W)$ at tree level, as well as the additional suppression of Q_W^e due to radiative corrections. This suppression, however, renders the attainment of high precision more difficult than for some of the other cases, since the statistical uncertainty in A_{LR} goes as $1/A_{LR}$ [35,21]. To set the scale, we note that a 10% $A_{LR}(^1\mathrm{H})$ measurement would be as sensitive as the present cesium APV determination to the mass scale Λ. Given the performance of the beam and detectors at the Jefferson Lab, it appears that a future measurement of $A_{LR}(^1\mathrm{H})$ with 5% or better precision could be feasible [18]. Such a determination would yield new physics limits comparable to those from cesium APV should the atomic theory error be reduced to the level of the present experimental error. A 2.5% ep measurement would strengthen the present APV bounds by a factor of two. Sensitivity at this level would be competitive with those expected from high energy colliders by the end of the next decade [36,37]. The physics reach of a 6% determination of the Möller asymmetry would be similar to that of the present cesium measurement, though PV ee scattering is in general sensitive to a different set of new interactions than arise in the eq sector.

In the case of isotope ratios, which depend like $A_{LR}(^1\mathrm{H})$ on ΔQ_W^P, a 0.5% determination $\mathcal{R}_1$ would give new physics limits comparable to the present cesium results. The prospects for achieving this precision are promising. The Berkeley group, for example, expects to perform a 0.1% determination of $\mathcal{R}_1$ using the isotopes of Yb $N = 100 \rightarrow N = 106$ [8].[2] A measurement of such precision would double the present cesium sensitivity, neglecting nuclear structure corrections. Similarly, the Seattle group plans to conduct studies on the isotopes of Ba^+ ions [38]. For both Yb and Ba, the scale factors f_1 are similar to that for cesium isotopes, whereas f_2 depends strongly on the range ΔN.

As the discussion of the following section illustrates, variations from this general pattern of relative sensitivities occur when specific new physics scenarios are considered. For example, our assumption of purely isoscalar new interactions ($h_V^u = h_V^d$) in arriving at Table I renders the PV $N \rightarrow \Delta$ correction zero. In the case of purely isovector interactions, the scale factor for PV ep scattering becomes 6.6 while that for the PV $N \rightarrow \Delta$ asymmetry is 2.5. In short, the weak charge for a single heavy isotope is relatively insensitive to new isovector interactions. As a second example, the

[2] Nuclear structure uncertainties may cloud the interpretation of such a measurement, however (see Sec. V).

Möller asymmetry is at least an order of magnitude less sensitive to leptoquarks than are the other observables, even though it generally displays a relatively strong sensitivity to new heavy physics (see discussion in Sec. IV). Similarly, in E_6 models which give rise to leptoquarks, one has $h_V^u=0$ while $h_V^d\neq 0$. In this case, systems having a relatively large *d*- to *u*-quark ratio are advantageous. The scale factor *f* for PV *ep* scattering, for example, is reduced to 2.2 when considering such E_6 models. A similar reduction occurs in the scale factors for the isotope ratios $\mathcal{R}_i$, since these ratios, such as $A_{LR}(^1\mathrm{H})$, are sensitive primarily to ΔQ_W^P. We also note in passing that limits from high energy colliders are sometimes quoted assuming that the new physics couplings to *u* and *d* quarks are the same as in the standard model. While there is no *a priori* reason to invoke this assumption, it would imply that the new physics shifts δ_P and δ_i ($i=1,2$) are suppressed by the same $1-4\sin^2\theta_W$ factor which enters Q_W^P at tree level.

Finally, we make a few observations regarding the new physics corrections to the DIS asymmetry. The correction $\tilde{\delta}_1$ depends on the same combination of the h_{ij}^q that arises in the other PV observables, but with a different *u*- and *d*-quark weighting than appears anywhere else. As reflected in Table I, however, the sensitivity of $\tilde{\delta}_1$ to new physics is much weaker than for most of the other observables. The correction $\tilde{\delta}_2$, on the other hand, is significantly more sensitive to new four fermion interactions than is $\tilde{\delta}_1$. Moreover, its dependence on the h_{ij}^q differs from that of all the other PV observables discussed here. In fact, certain scenarios proposed for evading the atomic PV limits on the h_{ij}^q, such as SU(12) symmetry [4], would not apply to bounds of comparable strength obtained from $\tilde{\delta}_2$. Unfortunately, a precise determination of the $\tilde{a}_2$ term in the DIS asymmetry appears to be difficult.

IV. MODEL ILLUSTRATIONS

The interaction of Eq. (4) may be specified for different new physics scenarios. In what follows, we consider three examples which illustrate the relative sensitivities of PV observables to different models: (a) additional neutral gauge bosons, (b) leptoquarks and *R*-parity-violating supersymmetric models, and (c) fermion compositeness.

A. Additional neutral gauge bosons

The existence of additional, neutral gauge bosons is natural in the context of superstring-inspired E_6 theories, in which the spontaneous breakdown of E_6 symmetry results in the existence of one or more U(1) gauge symmetries beyond the $U(1)_Y$ of the standard model [39–42]. Additional neutral gauge bosons may also arise in left-right symmetric models [40,42]. It is conceivable that at least one of the neutral gauge bosons is sufficiently light to be of interest to low-energy neutral current processes. We let Z' and Z denote the ''new'' and standard model neutral gauge bosons, respectively. The exisentence of a light Z' which mixes with the Z is ruled out by Z-pole observables. In the event that the Z-Z' mixing angle is ≈ 0, however, LEP and SLC measurements provide rather weak constraints [41]. Consequently, we consider the case of zero mixing.

For the sake of illustration, we follow the E_6 analysis of Ref. [39], in which the different symmetry breaking scenarious can be parametrized by writing the Z' as

$$Z'=\cos\phi Z_\psi+\sin\phi Z_\chi. \tag{49}$$

Z_ψ and Z_χ arise, for example, from the breakdown $E_6\to SO(10)\times U(1)_\psi$ and $SO(10)\to SU(5)\times U(1)_\chi$. Since the multiplets of SO(10) contain both f and $\overline{f}$ for the leptons and quarks of the standard model, *C* invariance implies that the Z_ψ can have only axial vector couplings to these fermions. As a result, it cannot contribute at tree-level to low-energy PV observables. In the case of SU(5), however, the left-handed *d* quark and e^+ live in a different multiplet from the left-handed $\overline{d}$ and e^-, whereas the *u* and $\overline{u}$ live in the same multiplet. The Z_χ correspondingly has both vector and axial vector couplings to the electron and *d* quarks, and only axial vector *u*-quark couplings. In short, E_6 Z' bosons yield $h_V^u=0$ and h_V^d, $h_V^e\propto\sin\phi$.

According to the notation of Eq. (4), we have for E_6 models

$$\kappa^2=\alpha', \tag{50}$$

$$\Lambda^2=M_{Z'}^2, \tag{51}$$

$$h_V^u=0, \tag{52}$$

$$h_V^d=-h_V^e=[\sin^2\phi-\sqrt{15}\sin\phi\cos\phi/3]/20, \tag{53}$$

where α' is the fine structure constant associated with the new gauge coupling. Generally, one has [40]

$$\alpha'\lesssim\frac{5}{3}\frac{\alpha}{\cos^2\theta_W}\approx 2.2\alpha. \tag{54}$$

Different models for the Z' correspond to different choices for ϕ. Examples include the Z_η ($\tan\phi=-\sqrt{3/5}$) and the $Z_I(\tan\phi=-\sqrt{5/3})$, where the latter is associated with an additional ''inert'' SU(2) gauge group not contributing to the electromagnetic charge. From the standpoint of phenomenology, it is worth noting the dependence of h_V^d and h_V^e on the value of ϕ. For $\phi=\phi_c=\tan^{-1}(\sqrt{5/3})\approx 52°$, $h_V^d=0=h_V^e$. For $\phi>\phi_c$, $h_V^d>0$. From Eq. (15), we observe that δ_N is negative for $h_V^u=0$ and $h_V^d>0$. The most recent value of δ_N for cesium implies that $h_V^d>0$ at the one σ level, and therefore could not be explained models giving $\phi<\phi_c$. The model that gives nearly the largest possible contribution to the weak charge is the Z_χ, which corresponds to $\phi=90°$.

An interesting variation on the idea of extended gauge group symmetry is that of left-right symmetric theories. In such theories, the low-energy gauge group becomes $SU(2)_L\times SU(2)_R\times U(1)_{B-L}$, where $B-L=1/3$ for baryons and -1 for leptons. In the case of ''manifest'' left-

TABLE II. Present and prosepctive limits on two species of additional neutral gauge bosons. The third column gives the ratio of Fermi constants as defined in the text. The fourth and fifth columns give lower bounds on masses for the Z_χ and Z_{LR}, respectively, assuming the precision given in column two.

Observable	Precision	G'_χ/G_F	M_{Z_χ} (GeV)	$M_{Z_{LR}}$ (GeV)
$Q_W(\mathrm{Cs})$	1.3%	0.006	730	790
	0.35%	0.0016	1410	1520
$\mathcal{R}_1$	0.3%	0.006	740	360
	0.1%	0.002	1300	630
$Q_W(^1\mathrm{H})/Q_{\mathrm{EM}}(^1\mathrm{H})$	10%	0.010	580	285
	3%	0.003	1100	520
$Q_W(0^+,0)/Q_{\mathrm{EM}}(0^+,0)$	1%	0.004	910	920
$Q_W(e)/Q_{\mathrm{EM}}(e)$	7%	0.004	910	460
$A_{LR}(N\to\Delta)$	1%	0.013	490	920
$\tilde{a}_1$	1%	0.15	145	320

right symmetry the $\mathrm{SU}(2)_L$ and $\mathrm{SU}(2)_R$ couplings are identical. For this case, a second low-mass neutral gauge boson Z_R couples to fermions with the strengths [42]

$$h_V^u=-\frac{3}{5}\frac{\alpha}{4}\left(\frac{\alpha}{4}-\frac{1}{6\alpha}\right), \tag{55}$$

$$h_V^d=\frac{3}{5}\frac{\alpha}{4}\left(\frac{\alpha}{4}+\frac{1}{6\alpha}\right), \tag{56}$$

$$h_V^e=\frac{3}{5}\frac{\alpha}{4}\left(\frac{\alpha}{4}-\frac{1}{2\alpha}\right), \tag{57}$$

where

$$\alpha=\left(\frac{1-2\sin^2\theta_W}{\sin^2\theta_W}\right)^{1/2}\approx 1.53. \tag{58}$$

With this set of couplings, the combination appearing in the correction ΔQ_W^P is $2h_V^u+h_V^d\approx 0.012 \ll h_V^u, h_V^d$. Consequently, the sensitivies of the R_i and $A_{LR}(^1\mathrm{H})$ are suppressed relative to their generic scale. The corresponding mass limits on M_{Z_R} are weaker than those obtainable from cesium APV or $A_{LR}(0^+,0)$.

In Table II, we give the present and prospective sensitivities for two species of additional neutral gauge bosons, the Z_χ and Z_R. In particular, we show lower bounds on the Fermi constant associated with the new gauge boson Z', defined as

$$\frac{G'_a}{\sqrt{2}}\equiv\frac{g'^2}{8M^2_{Z'_a}}, \tag{59}$$

where g' is the coupling associated with the additional $\mathrm{U}(1)_a$ gauge group. Low-energy PV observables constrain the ratio $(g'/M_{Z'_a})$ and do not provide separate limits on the mass and coupling. Consequently, the ratio of G'_χ/G_F characterizes the strength of a new $\mathrm{U}(1)_\chi$ gauge interaction relative to the strength of the standard model. In general, mass bounds for the Z' can be obtained from the limits on G' under specific assumptions for g'. A comparison of such mass bounds is often instructive, so we quote such bounds in the final two columns of Table II. Lower bounds on M_χ are quoted assuming the maximal value for g' as given by Eq. (54). In the case of LR symmetry models with manifest LR symmetry, one has $g'=g$. The corresponding mass limits for the Z_{LR} are given in the final column of Table II. Since we only discuss the case of manifest LR symmetry above, we do not include bounds on G'_{LR}/G_F.

The limits in Table II lead to several observations. Primary among these is that low-energy PV already constrains the strength of new, low-energy gauge interactions to be at most a few parts in a thousand relative to the strength of the $\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y$ sector. When reasonable assumptions are made about new gauge couplings strengths, low-energy mass bounds now approach one TeV. The significance of these bounds becomes more apparent when a comparison is made with the results of collider experiments. The present 110 pb^{-1} $p\bar{p}$ data set analyzed by the CDF Collaboration yields a lower bound on $M_{Z_{LR}}$ of 620 GeV, assuming manifest LR symmetry [43]. The lower bound for M_{Z_χ} is 585 GeV, assuming no Z_χ decays to supersymmetric particles [43]. The sensitivity of cesium APV already exceeds these Tevatron bounds. In fact, collider experiments and low-energy PV provide complementary probes of extended gauge group structure. PV observables are sensitive to the vector couplings of the Z' to fermions. For a model for which this coupling is small or vanishing [e.g., the Z_ψ having $\phi=0°$ in Eq. (49)], PV observables cannot yield significant information. Collider experiments, on the other hand, retain a sensitivity to such Z' interactions. For models in which the ffZ' coupling is not suppressed, low-energy PV presently displays the greatest sensitivity.

A look to the future suggests that PV could continue to play such a complementary role. Assuming the collection of 10 fp^{-1} of data at TeV33, for example, the current Tevatron bounds on $M_{Z'}$ would increase by roughly a factor of 2 [36]. The prospective sensitivity of cesium APV, assuming a reduction in atomic theory error to the level of the present experimental uncertainty, would exceed the collider reach by $\sim$50%. Precise determinations of the isotope ratio $\mathcal{R}_1$ or various PV electron scattering asymmetries could also yield sensitivities which match or exceed the prospective TeV33 bounds. Only with the advent of the LHC or $\gtrsim$60 TeV hadron collider will high-energy machines probe masses significantly beyond those accessible with low-energy PV [36].

Finally, Table II illustrates the model-sensitivity of different PV observables. For the models considered here, the mass bounds do not scale with the f_i of Table I since $h_V^u \neq h_V^d$. Both the Z' in E_6 and the Z_{LR} couple more strongly neutrons than protons. Consequently, both R_1 and $A_{LR}(^1\mathrm{H})$ display weaker sensitivity to new gauge interactions than their generic sensitivities to new physics indicated in Table I.

B. Leptoquarks and supersymmetry

In early 1997, the H1 [44] and ZEUS [45] Collaborations reported the presence of anomalous events in high-$|q^2|$ e^+p collisions at HERA. These events have been widely interpreted as arising from s-channel lepton-quark resonances with mass $M_{\rm LQ}\approx 200$ GeV [46,47]. Given the stringent limits on the existence of vector leptoquarks (LQ's) obtained at Fermilab [46,48,49], scalar leptoquarks are the favored interpretation of the HERA events. Although the results remain controversial, they are nonetheless provocative and suggest a consideration of LQ effects in low-energy PV processes. To that end, we consider general LQ interactions of the form

$$\mathcal{L}_{\rm LQ}^{S}=\lambda_S(\phi\bar{e}_L q_R+\text{H.c.}), \tag{60}$$

$$\mathcal{L}_{\rm LQ}^{V}=\lambda_V(\bar{e}_L\gamma_\mu q_L\phi^\mu+\text{H.c.}), \tag{61}$$

where ϕ and ϕ^μ denote scalar and vector LQ fields, respectively. For simplicity, we do not explicitly consider the corresponding interactions obtained from Eq. (61) with $L\leftrightarrow R$. The corresponding analysis is similar to what follows. Assuming $M_{\rm LQ}^2\gg|q^2|$, the process $eq\to LQ\to eq$ gives rise to the following PV interactions:

$$\mathcal{L}_{\rm PV}^{S}=(\lambda_S/2M_{\rm LQ})^2[\bar{e}q\bar{q}\gamma_5 e-\bar{e}\gamma_5 q\bar{q}e], \tag{62}$$

$$\mathcal{L}_{\rm PV}^{V}=(\lambda_V/2M_{\rm LQ})^2[\bar{e}\gamma_\mu q\bar{q}\gamma^\mu\gamma_5 e+\bar{e}\gamma_\mu\gamma_5 q\bar{q}\gamma^\mu e]. \tag{63}$$

After a Fierz transformation, these become

$$\mathcal{L}_{\rm PV}^{S}=(\lambda_S/2\sqrt{2}M_{\rm LQ})^2\Big[\bar{e}\gamma_\mu\gamma_5 e\bar{q}\gamma^\mu q-\bar{e}\gamma_\mu e\bar{q}\gamma^\mu\gamma_5 q+\frac{1}{4}\bar{e}\sigma_{\mu\nu}e\bar{q}\sigma^{\mu\nu}\gamma_5 q-\frac{1}{4}\bar{e}\sigma_{\mu\nu}\gamma_5 e\bar{q}\sigma_{\mu\nu}q\Big], \tag{64}$$

$$\mathcal{L}_{\rm PV}^{V}=-(\lambda_V/2M_{\rm LQ})^2[\bar{e}\gamma_\mu e\bar{q}\gamma^\mu\gamma_5 q+\bar{e}\gamma_\mu\gamma_5 e\bar{q}\gamma^\mu q]. \tag{65}$$

In terms of the interaction in Eq. (4), we may identify

$$\Lambda^2=M_{\rm LQ}^2, \tag{66}$$

$$\kappa^2=\lambda^2/16\pi, \tag{67}$$

and $h_V^q=1/2(h_V^q=-1)$ for scalar (vector) LQ interactions.

Assuming for simplicity that either a u- or d-type LQ (but not both) contributes to low-energy PV processes, the results from cesium APV, together with Eqs. (4) and (15), yield the following 1σ limits on LQ couplings and masses:

$$\lambda_S\leqslant\begin{cases}0.042(M_{\rm LQ}/100\ \text{GeV}), & u\ \text{type},\\ 0.04(M_{\rm LQ}/100\ \text{GeV}), & d\ \text{type}\end{cases} \tag{68}$$

and

$$\lambda_V\leqslant\begin{cases}0.030(M_{\rm LQ}/100\ \text{GeV}), & u\ \text{type},\\ 0.028(M_{\rm LQ}/100\ \text{GeV}), & d\ \text{type}.\end{cases} \tag{69}$$

TABLE III. Present and prosepctive limits on leptoquark interactions. The third and fourth columns give γ_q for a q-type leptoquark, as defined in Eq. (70). The leptoquark sensitivity of Möller asymmetry does not behave according to Eq. (70), so that no limits on γ_q are attainable.

Observable	Precision	γ_u	γ_d
$Q_W(\text{Cs})$	1.3%	0.04	0.042
	0.35%	0.021	0.022
$\mathcal{R}_1$	0.3%	0.04	0.028
	0.1%	0.023	0.016
$Q_W(^1\text{H})/Q_{\rm EM}(^1\text{H})$	10%	0.05	0.036
	3%	0.028	0.02
$Q_W(0^+,0)/Q_{\rm EM}(0^+,0)$	1%	0.033	0.033
$Q_W(e)/Q_{\rm EM}(e)$	7%		
$A_{LR}(N\to\Delta)$	1%	0.06	0.06
$\tilde{a}_1$	1%	0.14	0.20

Substituting the HERA value of $M_{\rm LQ}\approx 200$ GeV into Eq. (68) yields an upper bound of $\lambda_S\leqslant 0.08$. On general grounds, one might have expected $\kappa^2\sim\alpha$ or $\lambda_S\sim 0.6$. The cesium APV results require the coupling for a 200 GeV scalar LQ to be about an order of magnitude smaller than this expectation. Alternatively, if one does not interpret the HERA results as a 200 GeV LQ and assumes $\kappa^2\sim\alpha$, the APV bounds on the scalar LQ mass are $M_{\rm LQ}>1.5$ TeV. These LQ constraints are consistent with those obtained from high-energy collider experiments, though low- and high-energy processes generally provide complementary information. The constraints from the Tevatron [50], for example, are essentially λ_S independent, while providing bounds on $M_{\rm LQ}$ and LQ decay branching fraction [46,51].

Table III gives comparable bounds on the LQ coupling-to-mass ratio for the other PV observables discussed in Sec. III. The bounds are characterized by the quantity γ_q, defined as

$$\lambda_S\leqslant\gamma_q(M_{\rm LQ}/100\ \text{GeV}), \tag{70}$$

where q denotes the quark flavor.

Note that no bounds are given for the Möller asymmetry, as LQ's do not contribute at tree level. The leading contributions arise from the loop graphs of Fig. 1. We have evalu-

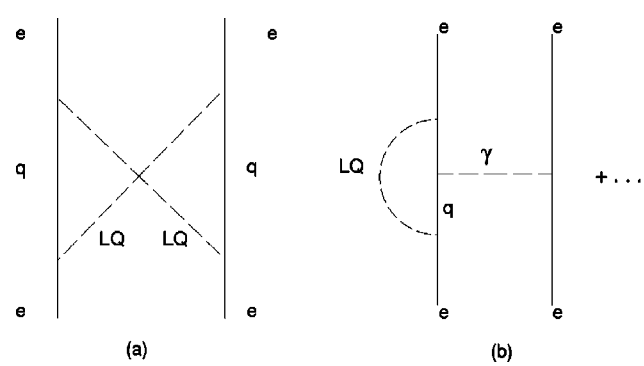

FIG. 1. Leptoquark (LQ) one-loop contributions to PV Möller scattering.

ated the amplitudes for these diagrams and obtain the following contributions to the PV effective ee interaction (to leading order in fermion masses and momenta):

$$\mathcal{L}^{\mathrm{PV}}_{(\mathrm{a})}=\left(\frac{\lambda_S^2}{16\pi M_{\mathrm{LQ}}}\right)^2\bar{e}\gamma_\mu e\bar{e}\gamma^\mu\gamma_5 e, \tag{71}$$

$$\mathcal{L}^{\mathrm{PV}}_{(\mathrm{b})}=\frac{\alpha Q_q}{12\pi}\left(\frac{\lambda_S}{M_{\mathrm{LQ}}}\right)^2\ln\frac{m_q}{M_{\mathrm{LQ}}}\bar{e}\gamma_\mu e\bar{e}\gamma^\mu\gamma_5 e, \tag{72}$$

where m_q and Q_q are the intermediate state quark mass and EM charge. For $M_{\mathrm{LQ}}=100$ GeV, a 7% determination of the Möller asymmetry would yield

$$\lambda_S\leqslant 1.06 \tag{73}$$

from graph (a) and

$$\lambda_S\leqslant\begin{cases}0.88, & d \text{ type},\\ 0.6, & u \text{ type}\end{cases} \tag{74}$$

from graph (b). The limits for a vector LQ are comparable. The prospective Möller bounds are more than an order of magnitude weaker than those attainable with semileptonic PV. Any deviation of the Möller asymmetry from the standard model prediction is unlikely to be due to LQ's.

Scalar leptoquarks arise naturally in R-parity-violating supersymmetric theories from a term in the superpotential of the form [52]

$$\lambda'_{ijk}L^i_L Q^j_L \overline{D}^k_R, \tag{75}$$

where the chiral superfields L^i_L, Q^j_L, and $\overline{D}^k_R$ contain the left-handed lepton and quark doublets and right handed d-quark singlets, respectively, for generations i,j,k. This term includes a lepton-number-violating electron-quark-squark interaction [52–54]

$$\mathcal{L}=\lambda'_{1jk}[\bar{d}_{kR}e_L\phi^u_{jL}+\bar{e}_L d_{kR}\bar{\phi}^u_{jL}-\overline{e^c}u_{jL}\bar{\phi}^d_{kR}-\bar{u}_{jL}e^c\phi^d_{kR}], \tag{76}$$

where ϕ^u_{jL} is the squark of charge $+2/3$ associated with a left-handed $+2/3$ charged quark of generation j, etc. The first two terms in Eq. (76) contribute to the HERA processes for $k=1$ when a positron scatters from a valence d quark in the proton, while the last two terms contribute for scattering from a sea u quark. Low-energy PV receives a contribution from both terms. For illustrative purposes, we consider only the first two. Identifying the λ'_{1j1} with λ_S of Eq. (61), we obtain

$$\lambda'_{1j1}\leqslant 0.04(M_{\phi^u_{jL}}/100\ \ \mathrm{GeV}) \tag{77}$$

as the bound obtained from cesium APV. The prospective bounds attainable from other PV observables may be obtained from Table III.[3]

For completeness, we note that low-energy PV is sensitive to another R-breaking term

$$\lambda_{ijk}L^i_L L^j_L \overline{E}^k_R, \tag{78}$$

where $\overline{E}^k_R$ contains the right-handed charged-lepton singlet fields. This term generates a four-fermion contact interaction which contributes to μ decay [52]

$$\mathcal{L}=-(\lambda_{12k}/\sqrt{2}M_{\phi^e_{kR}})^2\bar{e}_L\gamma_\alpha\nu^e_L\bar{\nu}^\mu_L\gamma^\alpha\mu_L. \tag{79}$$

Because the strength of the weak neutral current amplitude $(g/M_W)^2$ is written in terms of the μ-decay Fermi constant G_μ, the interaction (79) induces a correction to low-energy PV interactions:

$$\frac{g^2}{8M_W^2}=\frac{G_\mu}{\sqrt{2}}-(\lambda_{12k}/2\sqrt{2}M_{\phi^e_{kR}})^2\equiv\frac{G_\mu}{\sqrt{2}}[1-\Delta_{12k}], \tag{80}$$

where

$$\Delta_{12k}=\frac{\lambda^2_{12k}}{4\sqrt{2}G_\mu M^2_{\phi^e_{kR}}}. \tag{81}$$

A 1% determination of *any* low-energy PV observable (including the Möller asymmetry) would yield the bounds

$$\lambda_{12k}\leqslant 0.08(M_{\phi^e_{kR}}/100\ \ \mathrm{GeV}). \tag{82}$$

It is instructive to compare this bound with that obtained from superallowed β decay. In the latter case, interaction (79) would cause the measured value of the CKM matrix element $|V_{ud}|$ to differ from a valued implied by CKM matrix unitarity [52]. Letting $|V_{ud}|_{\mathrm{EX}}$ denote the value extracted from experiment—assuming only the standard model—and $|V_{ud}|$ the value implied by unitarity, one has

$$|V_{ud}|^2_{\mathrm{EX}}=|V_{ud}|^2[1-2\Delta_{12k}]. \tag{83}$$

The experimental situation regarding superallowed β decay has generated some debate about the value of $|V_{ud}|_{\mathrm{EX}}$. Assuming the experimental values for $|V_{us}|$ and $|V_{ub}|$, one finds from a fit to nine precisely measured superallowed $\mathcal{F}t$ values [56,57]

$$|V_{ud}|^2_{\mathrm{EX}}-|V_{ud}|^2=-0.0023\pm 0.0013. \tag{84}$$

A recent measurement of the superallowed ^{10}C β decay, however, yields a value consistent with CKM unitarity at the 1σ level [58]

$$|V_{ud}|^2_{\mathrm{EX}}-|V_{ud}|^2=-0.001\pm 0.0027. \tag{85}$$

The ^{10}C result, together with Eq. (83), requires that

$$2\Delta_{12k}|V_{ud}|^2\leqslant 0.0027 \tag{86}$$

[3]The most stringent bounds on λ'_{111} are derived from neutrinoless double β decay [55].

or

$$\lambda_{12k} \leq 0.03(M_{\phi^e_{kR}}/100\ \ \text{GeV}). \tag{87}$$

If, on the other hand, one assumes that the 2σ deviation is due to some type of new physics, then it could be generated by the lepton number violating interaction (79), since Δ_{12k} enters Eq. (83) with the correct sign.[4] In this case, the inequalities in Eqs. (86) and (87) would be replaced by the appropriate equalities.

C. Compositeness

The standard model assumes the known bosons and fermions to be pointlike. The possibility that they possess internal structure, however, remains an intriguing one. Manifestations of such composite structure could include the presence of fermion form factors in elementary scattering processes [59] or the existence of new, low-energy contact interactions [60]. The latter could arise, for example, from the interchange of fermion constituents at very short distances [40]. A recent analysis of $p\bar{p}\rightarrow \ell^+\ell^-$ data by the CDF Collaboration limits the size of a lepton or quark to be $R<5.6\times 10^{-4}$ f when R is determined from the assumed presence of a form factor at the fermion-boson vertex [59]. More stringent limits on the distance scale associated with compositeness are obtained from the assumption of new contact interations governed by a coupling of strength $g^2=4\pi$. Collider experiments yield $R\sim 1/\Lambda < 6\times 10^{-5}$ f, where Λ is the mass scale associated with new dimension 6 lepton-quark operators [59].

It is conventional to write the lowest dimension contact interactions as

$$\mathcal{L}_{\text{comp}}=4\pi\sum_{ij}\frac{\eta_{ij}}{\Lambda_{ij}^2}\bar{e}_i\Gamma e_i\bar{q}_j\Gamma q_j\,, \tag{88}$$

where Γ is any one of the Dirac matrices and i,j denote the appropriate fermion chiralities (e.g., $\bar{e}_L e_R \bar{q}_L q_R$ or $\bar{e}_L\gamma_\mu e_L \bar{q}_R\gamma^\mu q_R$, etc.). For simplicity, we restrict our attention to $\Gamma=\gamma^\mu$. The quantities η_{ij} take on the values $\pm 1,0$ depending on one's model assumptions. In terms of the PV interaction of Eq. (4), the contribution from $\mathcal{L}_{\text{comp}}$ is

$$\pi\bar{e}\gamma_\mu\gamma_5 e\sum_q\left[\frac{\eta_{RR}}{\Lambda_{RR}^2}-\frac{\eta_{LL}}{\Lambda_{LL}^2}+\frac{\eta_{RL}}{\Lambda_{RL}^2}-\frac{\eta_{LR}}{\Lambda_{LR}^2}\right]\bar{q}\gamma^\mu q. \tag{89}$$

Writing this interaction in terms of a common mass scale Λ yields

$$\frac{\pi}{\Lambda^2}\bar{e}\gamma_\mu\gamma_5 e\sum_q\ [\tilde{\eta}_{RR}-\tilde{\eta}_{LL}+\tilde{\eta}_{RL}-\tilde{\eta}_{LR}], \tag{90}$$

[4] In general leptoquark models where both u- and d-type LQ's contribute to β decay, the sign of the corresponding correction to unitarity may also be consistent with Eq. (84) under certain assumptions [51].

TABLE IV. Present and prospective limits on compositeness scale for the LL scenario.

Observable	Precision	Λ_{LL}(TeV)
Q_W(Cs)	1.3%	17.3
	0.35%	33.3
$\mathcal{R}_1$	0.3%	21.6
	0.1%	37.4
$Q_W(^1\text{H})/Q_{\text{EM}}(^1\text{H})$	10%	17.5
	3%	31.8
$Q_W(0^+,0)/Q_{\text{EM}}(0^+,0)$	1%	21.6
$Q_W(e)/Q_{\text{EM}}(e)$	7%	17.1[a]
$\tilde{a}_1$	1%	2.6

[a] Möller limits refer to new ee compositeness interactions, while other enteries refer to eq interactions.

where

$$\tilde{\eta}_{ij}=\eta_{ij}\left(\frac{\Lambda}{\Lambda_{ij}}\right)^2. \tag{91}$$

The correspondence with $\mathcal{L}_{\text{new}}^{\text{PV}}$ is given by

$$\kappa^2=1/4, \tag{92}$$

$$h_V^q=\tilde{\eta}_{RR}-\tilde{\eta}_{LL}+\tilde{\eta}_{RL}-\tilde{\eta}_{LR}\,. \tag{93}$$

On the most general grounds, one has no strong argument for any of the h_V^q to vanish. Consequently, low-energy observables will generate lower bounds on Λ. To compare with the recent CDF limits, we consider the case of $\tilde{\eta}_{LL}=\pm 1$ and $\tilde{\eta}_{RR}=\tilde{\eta}_{RL}=\tilde{\eta}_{LR}=0$. In this case, the cesium APV results yield

$$\Lambda_{LL}\geq 17.3\ \ \text{TeV} \tag{94}$$

assuming $h_V^u=h_V^d=-\tilde{\eta}_{LL}$. Regarding other low-energy PV observables, we note that the general comparisons made in Sec. III apply here. Hence, a 10% measurement of δ_P with PV ep scattering would yield comparable bounds, while a measurement of the isotopte ratio $\mathcal{R}_1$ with 0.5% precision would be required to obtain comparable limits. Were the cesium APV theory error reduced to the level of the present experimental error, or were a 2–3% determination of δ_P achieved, the lower limit (94) would double. Specific sensitivities from present and prospective measurements are given in Table IV.

As with other new physics scenarios, the present and prospective low-energy limits on compositeness are competitive with those presently obtainable from collider experiments as well as those expected in the future. The CDF Collaboration has obtained lower bounds on $\Lambda_{LL}(eq)$ of 2.5 (3.7) TeV for $\tilde{\eta}_{LL}=+1(-1)$ [59]. One expects to improve these bounds to 6.5 (10) TeV with the completion of Run II and 14 (20) TeV with TeV33 [62]. It is conceivable that future improvements in determinations of Q_W with APV or scattering will yield stronger bounds that those expected from colliders. In the case of $\Lambda_{LL}(ee)$, Z-pole observables imply lower bounds of 2.4 (2.2) TeV for $\tilde{\eta}_{LL}=+1(-1)$ [61]. The prospective Möller PV lower bounds exceed the LEP limits considerably.

The strength of these low-energy PV bounds has inspired various proposals for evading them. These scenarios include requiring $\mathcal{L}_{\rm comp}$ to be parity invariant [63] ($\tilde{\eta}_{RR}=\tilde{\eta}_{LL}$, $\tilde{\eta}_{RL}=\tilde{\eta}_{LR}$) or to satisfy SU(12) symmetry [4] (in effect, $\tilde{\eta}_{iL}=-\tilde{\eta}_{iR}$, that is, the new quark currents are purely axial vector).

V. THEORETICAL UNCERTAINTIES

The PV Möller asymmetry is the theoretically cleanest low-energy PV new physics probe. The dominant theoretical uncertainties are associated with hadronic contributions to the Z-γ mixing tensor, and they do not appear to be problematic for the extraction of new physics limits [34]. The attainment of stringent limits on new physics scenarios from low-energy semileptonic PV observables, however, requires that conventional many-body physics of atoms and hadrons be sufficiently well understood. At present, the dominant uncertainty in $Q_W(\mathrm{Cs})$ is theoretical. A significant improvement in the precision with which this quantity is known requires considerable progress in atomic theory. The issues involved in reducing the atomic theory uncertainty are discussed elsewhere [5,64,65]. In this section, we discuss the many-body uncertainties associated with the other semileptonic observables discussed above.

A. Isotope ratios

It was pointed out in Refs. [25,26] that the isotope ratios $\mathcal{R}_i$ display an enhanced sensitivity to the neutron distribution $\rho_n(r)$ within atomic nuclei, and that uncertainties in $\rho_n(r)$ could hamper the extraction of new physics limits from the $\mathcal{R}_i$. In Ref. [26], only $\mathcal{R}_2$ was considered, and only the implications of $\rho_n(r)$ uncertainties for the determination of $\sin^2\theta_W$ were discussed. For completeness, we consider also $\mathcal{R}_1$—which displays a greater new physics sensitivity than $\mathcal{R}_2$—and quantify the implications of $\rho_n(r)$ uncertainties for the extraction of new physics limits.

In general, one may express the weak charge as

$$Q_W=ZQ_W^P q_p+NQ_W^N q_n, \tag{95}$$

where

$$q_p=(1/\mathcal{N})\int d^3x\langle P|\hat{\psi}_e^\dagger(\vec{x})\gamma_5\hat{\psi}_e(\vec{x})|S\rangle\rho_p(\vec{x}), \tag{96}$$

$$q_n=(1/\mathcal{N})\int d^3x\langle P|\hat{\psi}_e^\dagger(\vec{x})\gamma_5\hat{\psi}_e(\vec{x})|S\rangle\rho_n(\vec{x}), \tag{97}$$

where $\hat{\psi}_e(\vec{x})$ is the electron field operator, $|S\rangle$ and $|P\rangle$ are atomic $S_{1/2}$ and $P_{1/2}$ states, and $\mathcal{N}$ is the value of the electron matrix element at the origin. The latter matrix element may be written as

$$\langle P|\hat{\psi}_e^\dagger(\vec{x})\gamma_5\hat{\psi}(\vec{x})|S\rangle=\mathcal{N}f(x), \tag{98}$$

where $f(0)=1$. The effect of uncertainties in $\rho_p(\vec{x})$—which are smaller than those in $\rho_n(\vec{x})$—are suppressed in Q_W since q_p is multiplied by the small number Q_W^P. Consequently, we consider only q_n.

To obtain general features, we follow Refs. [25,26] and consider a simple model in which the nucleus is treated as a sphere of uniform proton and neutron number densities out to radii R_P and R_N, respectively. In this case, one obtains [26]

$$q_n=1-(Z\alpha)^2 f_2^N+\cdots, \tag{99}$$

where

$$f_2^N=\frac{3}{10}x_N^2-\frac{3}{70}x_N^4+\frac{1}{450}x_N^6, \tag{100}$$

$$x_N=R_N/R_P. \tag{101}$$

Letting δ_N^n and δ_i^n denote the $\rho_n(\vec{x})$ corrections to $Q_W(N)$ and $\mathcal{R}_i$, respectively, we obtain

$$\delta_N^n\approx-(Z\alpha)^2 f_2^N(x_N), \tag{102}$$

$$\delta_1^n\approx-(Z\alpha)^2(N'/\Delta N)f_2^{N'}(x_N)\Delta x_N, \tag{103}$$

$$\delta_2^n\approx-(Z\alpha)^2 f_2^{N'}(x_N)\Delta x_N, \tag{104}$$

where $\Delta x_N=(R_{N'}-R_N)/R_P$. Uncertainties in Q_W and $\mathcal{R}_i$ arise from *uncertainties* in these quantities:

$$\delta(\delta_N^n)\approx-(Z\alpha)^2 f_2^{N'}(x_N)\delta x_N, \tag{105}$$

$$\delta(\delta_1^n)\approx-(Z\alpha)^2(N'/\Delta N)[f_2^{N'}(x_N)\delta(\Delta x_N)+(\Delta x_N)f_2^{N''}(x_N)\delta x_N], \tag{106}$$

$$\delta(\delta_2^n)\approx-(Z\alpha)^2[f_2^{N'}(x_N)\delta(\Delta x_N)+(\Delta x_N)f_2^{N''}(x_N)\delta x_N], \tag{107}$$

where δx_N is the uncertainty in x_N, etc.

From the standpoint of extracting new physics limits, the impact of neutron distribution uncertainties is characterized by the ratio of the $\delta(\delta_k^n)$ to the new physics corrections δ_k $(k=N,1,2)$. The smaller the size of this ratio, the less problematic neutron distribution uncertainties become. In the case of the isotope ratios, we observe that

$$\delta(\delta_1^n)/\delta_1\approx-\left(\frac{N'}{\Delta N}\right)\times\frac{(Z\alpha)^2}{\zeta}\frac{[f_2^{N'}(x_N)\delta(\Delta x_N)+(\Delta x_N)f_2^{N''}(x_N)\delta x_N]}{2h_V^u+h_V^d}\approx\delta(\delta_2^n)/\delta_2. \tag{108}$$

In short, the relative size of the corrections induced by new physics and neutron distribution uncertainties is essentially the same, whether one employs $\mathcal{R}_1$ or $\mathcal{R}_2$. Although $\mathcal{R}_1$ is

TABLE V. Neutron distribution uncertainties in atomic parity violation. The first line gives results for ^{133}Cs and following four give results for isotope ratios. The isotope spread ΔN is taken from Ref. [28] for Cs and Ba, from Ref. [8] for Yb, and from Ref. [26] for Pb. The fourth column gives required precision in neutron radius and isotope shift in order to keep neutron distribution uncertainty below the level quoted in column three. The fifth column gives theoretical estimates of neutron distribution uncertainties.

Observable	Precision	ΔN	Requirement	Theory
$Q_W(\mathrm{Cs})$	0.35%	0	$\delta x_N \leqslant 0.05$	$\delta x_N \leqslant 0.02$[a]
$\mathcal{R}_1(\mathrm{Cs})$	0.1%	14	$\delta(\Delta x_N) \leqslant 0.0024$	$\delta(\Delta x_N) \leqslant 0.0033 \rightarrow 0.0043$[b]
$\mathcal{R}_1(\mathrm{Ba})$	0.1%	14	$\delta(\Delta x_N) \leqslant 0.0023$	$\delta(\Delta x_N) \leqslant 0.0038$[c]
$\mathcal{R}_1(\mathrm{Yb})$	0.1%	6	$\delta(\Delta x_N) \leqslant 0.0005$	
$\mathcal{R}_1(\mathrm{Pb})$	0.1%	6	$\delta(\Delta x_N) \leqslant 0.0003$	$\delta(\Delta x_N) \leqslant 0.005$[d]

[a]Reference [27].
[b]Reference [27].
[c]Reference [28].
[d]Reference [26].

more sensitive to new physics by $N'/\Delta N$ as compared to $\mathcal{R}_2$, it is also more sensitive to $\rho_n(\vec{x})$ uncertainties by the same factor.

To set the scale of $\rho_n(\vec{x})$ uncertainties, we set $x_N \approx 1$ in Eqs. (101)–(107):

$$\delta(\delta_N^n) \approx -(3/7)(Z\alpha)^2 \delta x_N, \tag{109}$$

$$\delta(\delta_1^n) \approx -(N'/\Delta N)(Z\alpha)^2[(3/7)\delta(\Delta x_N) + (1/8)\Delta x_N \delta x_N], \tag{110}$$

$$\delta(\delta_2^n) \approx -(Z\alpha)^2[(3/7)\delta(\Delta x_N) + (1/8)\Delta x_N \delta x_N]. \tag{111}$$

In general, one has $\Delta x_N \delta x_N \ll \delta(\Delta x_N)$ [26]. Consequently, we keep only the terms associated with the uncertainty in the isotope shift $\delta(\Delta x_N)$.

We specify these expressions for the case of Cs, Yb, Ba, and Pb. Although no studies of cesium isotope ratios are planned at present, we include it in order to make a direct comparison between the single isotope and isotope ratios for this atom. The Yb and Ba isotopes are under study by the Berkeley and Seattle groups, respectively. We also include lead since it is one of the best understood heavy nuclei, both experimentally and theoretically. The neutron distribution uncertainties are shown for $Q_W(^{133}\mathrm{Cs})$ and for $\mathcal{R}_1$ for Cs, Yb, Ba, and Pb. In light of Eq. (108), it is sufficient to consider only $\mathcal{R}_1$. The fourth column of Table V gives the requirement on neutron distribution uncertainties for a given uncertainty in the corresponding APV observable. For $Q_W(\mathrm{Cs})$, we require $\delta(\delta_N^n)$ to be smaller than the present experimental uncertainty. For the isotope ratios, the requirement is $\delta(\delta_1^n) \leqslant 0.1\%$. In either case, the requirement must be met if the present cesium APV new physics reach is to be doubled. In the final column, we list published theoretical esimates of the corresponding neutron distribution uncertainty. The range in the case of $\mathcal{R}_1(\mathrm{Cs})$ corresponds to using the nominal error of Ref. [27] (larger value) and the spread between two models used in the calculation (lower value).

At present, there exist no reliable experimental determinations of x_N or Δx_N, so that the interpretation of APV observables must rely on nuclear theory.[5] It is conceivable that the theory uncertainty in x_N is 5% or better [26,27]. The estimate of Ref. [27] places this uncertainty closer to 2%. Consequently, one could argue that even if the atomic theory error in $Q_W(\mathrm{Cs})$ were reduced to the present experimental error, neutron distribution uncertainties should not complicate the extraction of new physics constraints. The situation regarding isotope shifts is more debatable.

Explicit studies of isotope shift uncertainties associated with $\rho_n(r)$ have been reported in Refs. [25–28]. The authors of Ref. [26] considered isotopes of lead using a variety of nuclear models and find a model spread of $\delta(\Delta x_N) \approx 0.005$, which corresponds to a 100% uncertainty in the model average for Δx_N. These authors note that the models used successfully predict the charge radii of even-even nuclei not used to fit the model parameters. The model spread is a factor of 10 larger than would be needed to keep the uncertainty in $\mathcal{R}_1(\mathrm{Pb})$ below 0.1%. Although the isotopes of lead are not presently under serious consideration for isotope ratio measurements, the scale of the model uncertainties for this well-understood set of isotopes is striking.

The authors of Ref. [27] employed two different Skyrme fits to compute Δx_N for Cs and Ba and quote an uncertainty in Δx_N of roughly 13% for the two series of isotopes (in the case of cesium, the difference in Δx_N between the two Skyrme fits is somewhat smaller than the quoted uncertainty). To our knowledge, there exist no published analyses of the ρ_n uncertainties for Yb. From the studies of Pb, Cs, and Ba, we infer that ρ_n uncertainties are presently larger than required for isotope ratio measurements to compete with those on a single isotope for yielding new physics limits.

Obtaining a sufficiently reliable computation of Δx_N remains an open problem for nuclear theory. It is argued in Ref. [26], for example, that model calculations contain a hidden uncertainty associated with the isovector surface term in

[5]Data from proton-nucleus and pion-nucleus exist for some cases, but the theoretical uncertainties are large. See, e.g., Ref. [26].

the nuclear energy functional. Changes in the coefficient of this term may signficantly affect a model calculation of Δx_N without affecting results for other observables. The authors of Ref. [27], on the other hand, considered this issue for cesium using a Skyrme interaction with two different parameter sets. For this interaction, changes in the isovector surface term larger enough to appreciably alter Δx_N also produce unacceptably large changes in binding energies. Whether the Skyrme results generalize to other interactions remains to be seen.

Given the present theoretical situation, a model-independent determination of $\rho_n(\vec{x})$ is desirable. To that end, PVES may prove useful [66,21]. Specifically, we consider a $(J^\pi,T)=(0^+,0)$ nucleus, such as ${}^{138}_{56}$Ba, noting that the isotopes of barium are under consideration for future APV isotope ratio measurements. As shown in Refs. [66,21], the PV asymmetry for $(0^+,0)$ nuclei may be written as

$$-\left[\frac{4\sqrt{2}\pi\alpha}{G_F|q^2|}\right]A_{LR}=Q_W^P+Q_W^N\frac{\int d^3x j_0(qx)\rho_n(\vec{x})}{\int d^3x j_0(qx)\rho_p(\vec{x})}. \tag{112}$$

Since $|Q_W^P/Q_W^N|\ll 1$, and since $\rho_p(\vec{x})$ is generally well determined from parity conserving electron scattering, A_{LR} is essentially a direct ''meter'' of the Fourier transform of $\rho_n(\vec{x})$. At low momentum-transfer ($qR_{N,P}\ll 1$) this expression simplifies:

$$-\left[\frac{4\sqrt{2}\pi\alpha}{G_F|q^2|}\right]A_{LR}\approx\frac{N}{Z}\left[1+\frac{q^2}{6}\left(\frac{R_P^2}{Z}-\frac{R_N^2}{N}\right)\right] \tag{113}$$

so that a determination of R_N is, in principle, attainable from A_{LR}.[6]

In a realistic experiment PVES experiment, one does not have $qR_{N,P}\ll 1$; larger values of q are needed to obtain the requisite precision for reasonable running times [66,21]. In Ref. [66], it was shown that a 1% determination of $\rho_n(\vec{x})$ for ^{208}Pb is experimentally feasible for $q\sim 0.5$ fm^{-1} with reasonable running times. An experiment with barium is particularly attractive. If the barium isotopes are used in future APV measurements as anticipated by the Seattle group, then a determination of $\rho_n(\vec{x})$ for even one isotope could reduce the degree of theoretical uncertainty for neutron distributions along the barium isotope chain. Moreover, the first excited state of ^{138}Ba occurs at 1.44 MeV. The energy resolution therefore required to guarantee elastic scattering from this nucleus is well within the capabilities of the Jefferson Lab.

The foregoing discussion illustrates general features of, and presents order of magnitude estimates for, neutron distribution effects in APV. A more complete analysis of q_n using realistic atomic wave functions will be required to translate PVES information on $\rho_n(q)$ into useful input for APV calculations. Indeed, the function $f(x)$ which weights $\rho_n(\vec{x})$ in Eq. (97) is not the same as the Bessel function $j_0(qx)$ which weights ρ_n in the asymmetry. Evidently, a determination of $\rho_n(q)$ over some range in q will be required. Assuming $\rho_n(\vec{x})$ can be sufficiently well determined for a single isotope, it remains to be seen how tightly such a determination constraint nuclear theory calculations of $\rho_n(\vec{x})$ along the isotope chain or elsewhere in the periodic table. A detailed treatment of these issues lies beyond the scope of the present study.

B. Hadronic form factors

From the form of Eq. (31), it is clear that a precise determination of Q_W from A_{LR} requires sufficiently precise knowledge of the form factor term $F(q)$. This term is presently under study at a variety of accelerators, with the hope of extracting information on the strange quark matrix element $\langle N(p')|\bar{s}\gamma_\mu s|N(p)\rangle$. The latter is parametrized by two form factors $G_E^{(s)}$ and G_E^p. The other form factors which enter $F(q)$ are known with much greater certainty than are the strange quark form factors. A separation of Q_W from $F(q)$ requires at least one forward angle measurement [35]. The kinematics must be chosen so as to minimize the importance of $F(q)$ relative to Q_W while keeping the statistical uncertainty in the asymmetry sufficiently small. These competing kinematic requirements—along with the desired uncertainty in Q_W—dictate the maximum uncertainty in $F(q)$ which can be tolerated. Since $A_{LR}(^1H)$ generally manifests the greatest sensitivity to new physics, we illustrate the form factor considerations for PV ep scattering.

Since $Q_{\rm EM}^p=1$, the ep asymmetry has the form

$$A_{LR}=a_0\tau[Q_W^P+F^p(q)], \tag{114}$$

where $a_0\approx -3.1\times 10^{-4}$ and $\tau=|q^2|/4m_N^2$. The form factor contribution is given at the tree level in the standard model by [35,21]

$$F^p(\tau)=-[G_E^p(G_E^n+G_E^{(s)})+\tau G_M^p(G_M^n+G_M^{(s)})]/[(G_E^p)^2+\tau(G_M^p)^2], \tag{115}$$

where $G_{E,M}^{p,n}$ denote the proton or neutron Sachs electric or magnetic form factors. Since $Q_{\rm EM}^n=0$ and since the proton carries no net strangeness, both G_E^n and $G_E^{(s)}$ must vanish at $\tau=0$. Consequently, we may write $F^p(\tau)$ as

$$F^p(q)=\tau B(\tau). \tag{116}$$

For the purposes of this discussion, it is useful to write

$$B(\tau)=B_0(\tau)+\delta B(\tau), \tag{117}$$

where B gives the contribution from G_E^p, G_E^n, G_M^p, and G_M^n while δB contains the contributions from $G_E^{(s)}$ and $G_M^{(s)}$ as well as from τ-dependent higher-order electroweak corrections, as discussed in the next section. As noted in Ref. [35],

[6]In a realistic analysis of A_{LR} for heavy nuclei, the effects of electron wave distortion must be included in the analysis of A_{LR}. For a recent distorted wave calculation, see Ref. [67].

TABLE VI. Conditions for a new physics search with PV elastic ep scattering. The first 12 lines give conditions for determination of Q_W^P with the precision listed in column four asssuming 10 msr solid angle detector, $\mathcal{L}=5\times10^{38}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1}$, and 100% beam polarization. The corresponding statistical uncertainty in the asymmetry is listed in column three, while the required precision in $B(\tau)$ is given in column five. The final three lines give present present and prospective determinations of $B(\tau)$.

θ	τ	$\delta A_{LR}/A_{LR}$	$\delta Q_W^P/Q_W^P$	$\delta B/B_0$
12.3 °	0.018	0.063[a]	0.112	0.145
		0.045[b]	0.08	0.103
6 °	0.004	0.087[a]	0.103	0.59
		0.062[b]	0.073	0.42
6 °	0.008	0.059[a]	0.08	0.227
		0.041[b]	0.056	0.161
6 °	0.012	0.043[a]	0.065	0.123
		0.033[b]	0.046	0.087
6 °	0.018	0.032[a]	0.057	0.074
		0.023[b]	0.04	0.052
6 °	0.024	0.025[a]	0.05	0.05
		0.018[b]	0.035	0.035
12.3 °	0.14	0.16[c]	0.032[e]	0.196
12.3 °	0.14	0.05[c]	0.032[e]	0.065
35 °	0.07	0.04[d]	0.032[e]	0.057

[a]1000 h of running time.
[b]2000 h of running time.
[c]Jefferson Laboratory (see Refs. [15,69]).
[d]Mainz (see Ref. [70]).
[e]The uncertainty in Q_W^P in the last three lines is computed using the hadronic uncertainty of Ref. [34].

any determination of Q_W^P must be made at such low τ that only $B(\tau=0)$ enters the analysis. The experimental problem is to measure $B(\tau)$ with sufficient precision over a sufficient range of τ such that $\tau\delta B(0)$ smaller than the desired uncertainty in Q_W^P in a low-τ measurement.

In the first 12 lines of Table VI, we summarize the conditions for several prospective determinations of Q_W^P with a measurement of $A_{LR}(^1\mathrm{H})$. The last three lines summarize existing or planned forward angle determinations of $B(\tau)$. For both sets of measurements, the third column gives the statistical uncertainty in the asymmetry, assuming a solid angle of 10 msr, a luminosity $\mathcal{L}=5\times10^{38}\ \mathrm{cm}^{-2}\ \mathrm{s}^{-1}$, and 100% beam polarization for various running times and kinematics [21]. The fourth column gives the corresponding experimental uncertainty in Q_W^P. The requirements to keep the error in Q_W^P from $B(\tau)$ smaller than the statistical uncertainty are given in the fifth column. For the second set of measurements (final three rows), only the standard model uncertainty in Q_W^P is listed. The dominant uncertainty arises from hadronic loops appearing in the Z-γ mixing tensor [34]. The uncertainty associated with the experimental value of $\sin^2\theta_W$ is about a factor of 3 smaller than the hadronic uncertainty [33]. We note that measurements at $\theta=6°$ would require the development of new beam optics for the CEBAF detectors; such developments appear technically feasible [16,68]. The choices for τ in the first 12 lines correspond roughly to CEBAF beam energies.

The results entries in Table VI illustrate the trade-offs between kinematics, desired precision in Q_W^P, and required precision in $B(\tau)$. For a given scattering angle θ, increasing τ decreases the statistical uncertainty in A_{LR} but increases the contribution from $\tau B(\tau)$. The latter increase has two effects. First, it reduces the relative contribution of Q_W^P, making it more difficult to match the fractional uncertainty in A_{LR} with the desired uncertainty in Q_W^P. Second, it imposes more stringent requirements on knowledge of $B(\tau)$. Consequently, it may be desireable to go to slightly longer running times and lower τ. Comparing the two possible measurements at $\theta=6°$, for example, we see that a 1000 h measurement at $\tau=0.024$ yields a 2.5% statistical uncertainty in A_{LR} but only a 5% uncertainty in Q_W^P. Moreover, the required precision on B is slightly more stringent than will be obtained with any of the current PVES measurements (last three lines). However, a 2000 h $\theta=6°$ experiment at $\tau=0.018$ yields a 4% determination of Q_W^P for a 2.3% measurement of A_{LR} while imposing similar requirements on δB. An even more precise determination of Q_W^P would require reduction in the hadronic uncertainty entering the standard model radiative corrections.

We emphasize that the entries in Table VI are intended as illustrative benchmarks. The optimal kinematics for a precise determination of Q_W^P require a detailed analysis of acutal experimental conditions at different laboratories. We also emphasize that the measured uncertainty in B at higher τ (last three lines of Table VI) does not necessarily translate into the same uncertainty at the lower τ needed for new physics searches. For example, the strange quark form factors may not scale with τ in the same way as the nucleon EM form factors. Hence, it is likely that measurements of $B(\tau)$ over a range of kinematics will be needed to sufficiently constrain its value at the photon point (see, e.g., Refs. [35,21]). A detailed analysis of this issue would constitute a critical component of an experimental proposal.

C. Dispersion corrections

The foregoing discussion has implicitly relied upon a first Born approximation of the electroweak amplitudes contributing to low-energy PV. A realistic analysis of precision observables must take into account contributions beyond the first Born amplitude. In the case of electron scattering, these contributions are generally divided into two classes: Coulomb distortion of plane wave electron wave functions and dispersion corrections. The former can be treated accurately for electron scattering using distorted wave methods. Results of such a treatment are reported in Ref. [67]. The dispersion correction, however, has proven less tractable.

The leading dispersion correction (DC) arises from diagrams of Fig. 2, where the intermediate state nucleus or hadron lives in any one of its excited states. More generally, box diagrams such as those of Fig. 2 can be treated exactly for scattering of electrons from pointlike hadrons. When at least one of the exchanged bosons is a photon, the amplitude is prone to infrared enhancements. For elastic PV scattering of

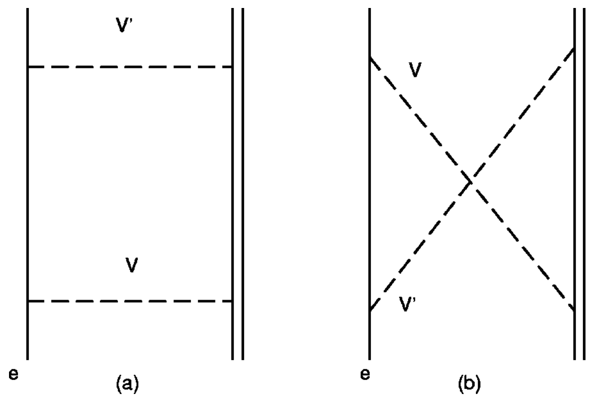

FIG. 2. Two vector boson exchange dispersion corrections. Here V and V' denote γ, Z^0, or $W^{\pm}$. Double vertical line denotes hadronic target.

an electron from a pointlike proton, for example, the Z-γ amplitude contains infrared enhancement factors such as $\ln|s|/M_Z^2$, where s is the ep c.m. energy [29]. Such factors can enhance the scale of the amplitude by as much as an order of magnitude over the nominal $\mathcal{O}(\alpha)$ scale. Consequently, one might expect box graph amplitudes which depend on details of hadronic or nuclear structure to be a potential source of theoretical error in the analysis of precision electroweak observables.

Data on the electromagnetic $(\gamma\gamma)$ dispersion correction for ep scattering is in general agreement with the scale predicted by theoretical calculations. The situation regarding electron scattering from nuclei, however, is less satisfying. Recent data for ^{12}C(e,e') taken at MIT-Bates and NIHKEF disagree dramatically with nearly all published calculations (for a more detailed discussion and references, see Ref. [71]). An experimental determination of any electroweak DC ($V=\gamma$, $V'=W^{\pm},Z^0$) is unlikely, and reliance on theory to compute this correction is unavoidable. As we show below, the corresponding theoretical uncertainty is far less problematic for a determination of Q_W from PVES than for the extraction of information on the strange quark form factors.

To this end, it is convenient to write the (V,V') DC as a correction $R_{VV'}$ to the tree level EM and PV neutral current ampltitudes [71]

$$M_{\rm EM}=M_{\rm EM}^{\rm tree}[1+R_{\gamma\gamma}+\cdots], \tag{118}$$

$$M_{\rm NC}^{\rm PV}=M_{\rm NC}^{\rm PV,\,tree}[1+R_{VV'}+\cdots], \tag{119}$$

where the ellipses denote other higher order corrections to the tree level amplitude. Because $M_{\rm EM}^{\rm tree}\propto 1/q^2$ while the $\gamma\gamma$ amplitude contains no pole at $q^2=0$, $R_{\gamma\gamma}$ has the general structure

$$R_{\gamma\gamma}(q^2)=q^2\tilde{R}_{\gamma\gamma}(q^2), \tag{120}$$

where $\tilde{R}_{\gamma\gamma}(q^2)$ describes the q^2 dependence of the $\gamma\gamma$ amplitude and $\tilde{R}_{\gamma\gamma}(0)$ is finite. Since the tree level NC amplitude contains no pole at $q^2=0$, however, the PV DC's do not vanish at $q^2=0$. Using Eqs. (118)–(120) and expanding the PV corrections in powers of q^2 we obtain

$$\frac{A_{LR}}{a_0\tau}=Q_W[1+R_{WW}(0)+R_{ZZ}(0)+R_{Z\gamma}(0)],+\tilde{F}(q) \tag{121}$$

where we replace the form factor $F(q)$ appearing in Eq. (31) by an effective form factor $\tilde{F}(q)$:

$$\tilde{F}(q)=F(q)+q^2[R'_{WW}(0)+R'_{ZZ}(0)+R'_{Z\gamma}(0)-\tilde{R}_{\gamma\gamma}(q^2)+\cdots], \tag{122}$$

with $F(q)$ containing the dependence on hadronic form factors as before.

From Eq. (121) we observe that the entire $\gamma\gamma$ DC, as well as the subleading q^2 dependence of the WW, ZZ, and $Z\gamma$ DC's, contribute to A_{LR} as part of an effective form factor term $\tilde{F}(q)$. Since $F(q)\sim q^2$ for low-$|q^2|$ at forward angles, the DC contributions entering Eq. (121) will be experimentally constrained along with $F(q)$ when the form factor term $\tilde{F}(q)$ is kinematically separated from the weak charge term. Consequently, an extraction of Q_W from A_{LR} does not require theoretical computations of the $\gamma\gamma$ DC or of the subleading q^2 dependence of the other DC's. A determination of the strange-quark form factors, however, does require such theoretical input.

In order to constrain possible new physics contributions to Q_W, a standard model theoretical calculation of $R_{WW}(0)$, $R_{ZZ}(0)$, and $R_{Z\gamma}(0)$ is necessary. The theoretical uncertainty associated with $R_{WW}(0)$ and $R_{ZZ}(0)$ is small, since box diagrams involving the exchange are dominated by hadronic intermediate states having momenta $p\sim M_W$. These contributions can be reliably treated perturbatively. The $R_{Z\gamma}(0)$ correction, however, is infrared enhanced and displays a greater sensitivity to the low-lying part of the nuclear and hadronic spectrum. Fortunately, the sum of diagrams 2(a) and 2(b) conspire to suppress this contribution by $g_V^e=-1+4\sin^2\theta_W$. This feature was first shown in Ref. [72] for the case of APV. Here, we summarize the argument as it applies to scattering.

The dominant contributions to the loop integrals for diagrams 2(a) and 2(b) arise when external particle masses and momenta are neglected relative to the loop momentum l_μ. In this case, the integrands from the two loop integrals sum to give

$$\bar{u}[\gamma_\alpha \not{l}\gamma_\beta(g_V^e+g_A^e\gamma_5)-\gamma_\beta(g_V^e+g_A^e\gamma_5)\not{l}\gamma_\alpha]uT^{\alpha\beta}(l)D(l^2)$$
$$=2i\epsilon_{\alpha\lambda\beta\mu}l^\lambda\bar{u}\gamma^\mu(g_V^e\gamma_5+g_A^e)uT^{\alpha\beta}(l)D(l^2), \tag{123}$$

where

$$T^{\alpha\beta}(l)=\int d^4x\ \ e^{il\cdot x}\langle 0|T\{J_{\rm EM}^\alpha(x)J_{\rm NC}^\beta(0)\}|0\rangle, \tag{124}$$

$D(l^2)$ contains the electron and gauge boson propagators when external momenta and masses are neglected relative to

l_μ, and $J^\alpha_{\rm EM}$ and $J^\beta_{\rm NC}$ are the hadronic electromagnetic and weak neutral currents, respectively. The terms in Eq. (123) which transform similar to pseudoscalars are those containing the EM current and either (a) both the axial currents $\bar{u}\gamma^\mu\gamma_5 u$ and $J^{\beta 5}_{\rm NC}$ or (b) both the vector currents $\bar{u}\gamma^\mu u$ and $J^\beta_{\rm NC}$. The former has the coefficient $g_V^e=-1+4\sin^2\theta_W$ and the latter has a the coefficient $g_A^e=1$. The dependence of these terms on the spatial currents is given by [$\lambda=0$ in Eq. (123)]

$$g_V^e \ \text{term:} \sim \bar{u}\vec{\gamma}\gamma_5 u(\vec{J}_{\rm EM}\times\vec{J}^5_{\rm NC}), \tag{125}$$

$$g_A^e \ \text{term:} \sim \bar{u}\vec{\gamma} u(\vec{J}_{\rm EM}\times\vec{J}_{\rm NC}). \tag{126}$$

The hadronic part of the g_V^e term transforms as a polar vector, so that this term contributes to the $A(e)\times V(\text{had})$ amplitude. The hadronic part of the g_A^e terms, on the other hand, transfors as an axial vector, yielding a contribution to the $V(e)\times A(\text{had})$ amplitude. Hence, only the g_V^e term contributes to Q_W term in the asymmetry.

Since $g_V^e\sim-0.1$, the contribution $R_{Z\gamma}(0)$ in Eq. (121) is suppressed. For scattering from (0^+,0) nuclei, then, $R_{Z\gamma}(0)\sim\mathcal{O}(\alpha/10)$, while for PV ep scattering, $R_{Z\gamma}\sim\mathcal{O}(\alpha)$. Since 1% and 5–10 % determinations of $Q_W(0^+,0)$ and $Q_W(p)$, respectively, are needed to constrain new physics scenarios, large theoretical uncertainties in $R_{Z\gamma}(0)$ should not be problematic. A similar statement applies to APV, for which contributions to $R_{Z\gamma}$ from excited nuclear states have yet to be computed. Whether these contributions can be reliably computed at the 0.3% level remains to be evaluated.

VI. CONCLUSIONS

The prospects for future, precise measurements of low-energy PV observables is promising. In addition to the approved PV Möller experiment and planned APV isotope measurements, a precise measurement of A_{LR} for PV electron-proton or electron-nucleus scattering at Jefferson Laboratory appears feasible. Depending on the degree of experimental and theoretical precision realized in each case, future measurements could improve upon the present cesium APV new physics sensitivity by a factor of 2. At the same time, such studies would complement future new physics searches at high-energy colliders. Indeed, while high-energy studies are particularly sensitive to the mass scale Λ associated with new interactions, low-energy PV probes the coupling-to-mass ratio, g/Λ. For new physics scenarios in which g is fixed (e.g., LR symmetric gauge theories or fermion compositeness), even the present cesium APV bounds on Λ exceed those obtained from the Tevatron or LEP2. Taken together, high-energy and low-energy PV measurements provide a powerful, combined probe of physics at the TeV scale.

As the discussion of Secs. III and IV illustrates, no single low-energy PV process is equally sensitive to every new physics scenario. For example, APV on a single isotope is strongly sensitive to new isoscalar interactions but much less transparent to new isovector heavy physics. Similarly, elastic PV ep scattering constitutes the most sensitive probe of new e-q physics (for a given experimental precision) except for scenarios in which new ep couplings are fortuitously suppressed (e.g., left-right symmetric or E_6 models). In addition, each low-energy process encounters its own brand of theoretical uncertainties which may limit the interpretation of a given measurement in terms of new physics.

We conclude that the most thorough search for new physics using low-energy PV would require a program of measurements drawing upon the complementarity of different processes. Here we summarize the elements of this complementarity

(a) Both ep asymmetry $A_{LR}(^1\text{H})$ and the isotope ratios $\mathcal{R}_i$ are sensitive to the *same* combination of new e-q interactions [see ΔQ_W^P of Eqs. (7),(25)–(26)].

(b) A 2–3 % determination of $A_{LR}(^1\text{H})$ or a 0.1% determination of $\mathcal{R}_1$ would nearly double the present $Q_W(\text{Cs})$ sensitivity for some scenarios (e.g., fermion compositeness and leptoquarks) but not others (e.g., right-handed neutral gauge bosons). Moreover, either of these PVES or isotope ratio measurements would, together with the present cesium APV result, afford a separate determination of new e-u and e-d interactions.

(c) The planned measurement of the Möller asymmetry will provide the best test of lepton compositeness of any electroweak observable, exceeding the $\Lambda_{ij}(ee)$ bounds from LEP by nearly an order of magnitude. However, PV ee scattering is 100 times less sensitive to leptoquark and R-parity-violating SUSY interactions than are semileptonic PV observables.

(d) The bounds on extra gauge bosons obtained from a one percent determination of $A_{LR}(N\rightarrow\Delta)$ and $A_{LR}(0^+,0)$ could exceed those derived from either (a), (d), or cesium APV.

Evidently, at least one additional measurement—in addition to the cesium APV and planned Möller experiments—is necessary to provide the complete range of low-energy information on new neutral current interactions.

From the standpoint of the *interpretation* of PV measurements, the Möller asymmetry provides the theoretically cleanest probe of new physics. The relevant theoretical uncertainties in this case are those associated with hadronic contributions to the Z-γ mixing tensor [34] and with the (small) scattering backgrounds [17]. Neither source of uncertainty appears to be problematic for the extraction of new physics limits from $A_{LR}(ee)$.

The interpretation of semileptonic observables, however, requires improved input from atomic, nuclear, and hadron structure theory. The most challenging theory issues lie with the APV observables. A reduction in the cesium atomic theory uncertainty by a factor of 4 would make it comparable to the present experimental error. In this case, the cesium new physics sensitivity would improve by a factor of 2. Whether or not such an improvement in the atomic theory can be achieved is an open question. In the case of APV isotope ratios, the attainment of the new physics sensitivity discussed above may require an experimental determination of $\rho_n(r)$ using PVES. At present, there exist no published estimates of the isotope shifts in ρ_n for Yb. Estimates for cesium and barium suggest that the theoretical isotope shift uncertainty may be about two times larger than desirable for future new physics searches. In the absence of improved nuclear theory input, measurements of the $\mathcal{R}$ will provide more information on nuclear structure than on new electroweak physics. A precise determination of ρ_n using PVES, however, may sufficiently constrain model calculations so as to significantly reduce the theoretical isotope shift uncertainty.

The theoretical issues entering the interpretation of semileptonic PVES appear less formidable. The dominant corrections to the Q_W term of the asymmetry—including both hadronic form factors and the $\gamma\gamma$ dispersion correction—are measurable in principle. The remaining hadron and nuclear structure-dependent corrections are fortuitously suppressed. Consequently, the primary challenge in peforming new physics searches with PVES will be experimental.

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