We studied $\beta$-delayed singles-$\alpha$ spectra from $^8$Li and $^8$B decays with special emphasis on a careful calibration of the energy scale. $^8$Li and $^8$B activities were produced by $^7$Li($d$, $p$) and $^6$Li($^3$He,$n$) reactions, respectively, and deposited in thin C foils. Delayed $\alpha$'s were counted in thin, cooled Si detectors with small solid angles to reduce summing with $\beta$'s. The energy scale and detector response were calibrated with spectroscopic grade radioactive sources. We extracted the $^8$Li and $^8$B final-state continuum shapes from our spectra by using $R$-matrix analyses that included effects of lepton-recoil broadening and detector response. Our results are in excellent agreement with a recent measurement using $^8$B's implanted in a Si counter and in good agreement with our reanalysis of the older Wilkinson-Alburger $^8$Li and $^8$B data, but disagree with a recent $^8$B experiment using a coincidence technique.

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I. INTRODUCTION

The shapes of the broad $^8$Be final-state continua populated in $^8$Li and $^8$B $\beta$ decays are interesting from a nuclear physics perspective. But, more important, the $^8$B $\beta$ final-state distribution determines the intrinsic spectrum of $^8$B neutrinos that dominate the counting rates of the Kamiokande [1], Super-Kamiokande [2] and SNO [3,4] solar neutrino detectors and form a major component of the Homestake detector [5] events. Precise knowledge of this undistorted spectrum is needed to extract as much information as possible about neutrino oscillation parameters from the observed spectrum of solar neutrinos detected on Earth.

The $J^\pi = 2^+$ $^8$Li and $^8$B ground states decay to a $2^+$ continuum in $^8$Be that very rapidly breaks up into $2\alpha$ particles. The only other energetically allowed $\beta$ decays are second-forbidden transitions to $0^+$ or $4^+$ states that can be neglected. The $2^+$ continuum is populated by a $^8$Be state at 3 MeV with a width of 1.5 MeV. The interference of this broad state with other $2^+$ states at higher energies affects the shapes of the final-state continuum and therefore the shape of the $\beta$ and neutrino spectra. As a result, high-quality measurements of the final-state continuum are needed for accurate computation of the neutrino spectrum.

The $^8$Li and $^8$B final-state continua can be inferred from either $\beta$ spectra or $\beta$-delayed $\alpha$ spectra. Bahcall et al. [6] showed that the existing $^8$B singles delayed-$\alpha$ spectra [7–9] were inconsistent with one another: the spectrum peaks in the different experiments varied by $\pm 80$ keV. Because of this inconsistency, Bahcall et al. [6] used the single available $\beta$ spectrum-shape measurement [10] to infer their standard $^8$B neutrino spectrum.

We have shown [11] that lepton recoil, which had been neglected in all previous analyses of the $A = 8$ delayed-$\alpha$ shapes, perceptibly distorts the singles-$\alpha$ spectrum shapes and therefore the extracted $\beta$ final-state distributions. Accounting for lepton recoil significantly reduced the discrepancies previously noted [12–14] between the delayed-$\alpha$ shapes in $^8$Li and $^8$B decays and for the $L = 2$ phase shifts in $\alpha + \alpha$ scattering.

Recently, Ortiz et al. [15,16] used a new technique to study the $^8$B delayed-$\alpha$ spectrum with high precision. They placed a pair of oppositely facing detectors in a strong magnetic field to prevent the associated $\beta$'s from reaching the detectors and measured the summed energy of the two coincident delayed $\alpha$'s. This eliminated the first-order effect of lepton-recoil broadening. However, the magnetic field gave the coincidence $\alpha$-detection efficiency a dependence on the $^8$B excitation energy, which had to be accounted for with a Monte Carlo simulation. Still more recently, Winter et al. [17,18] made a precise study of the $^8$B delayed-$\alpha$ spectrum by a different method, but one which also measured the summed energy of the two $\alpha$'s. Winter et al. implanted 27.3 MeV $^8$B's into a thin Si counter and detected the decay positrons in a plastic scintillator. The plastic scintillator coincidence selected events where positrons traveled within 30° of the detector normal to minimize the energy deposited by the positrons. However, these two results of nominally similar precision disagreed. Although the disagreement was smaller than that noted previously by Bahcall et al., it was nevertheless significant, considering the increasing precision of the solar neutrino data.

This paper reports a precise measurement of the delayed-$\alpha$ continuum in $^8$B and $^8$Li decays performed by using a conventional singles-$\alpha$ technique (rather than the summed-$\alpha$ techniques employed in Refs. [15] and [17,18]). Our data were taken and analyzed after the publication of Ref. [15] but before that of Ref. [17].

II. EXPERIMENTAL PROCEDURE

A. Apparatus

Figure 1 shows a schematic diagram of our apparatus; details can be found in Ref. [19]. The $^8$Li and $^8$B activities were produced via $^7$Li($d$, $p$) and $^6$Li($^3$He,$n$) reactions at

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bombarding energies of 1.0 and 5.5 MeV, respectively. The University of Washington tandem accelerator fitted with a terminal ion source delivered \( \sim 16 \mu\text{A} \) 1.0 MeV d and \( \sim 12 \mu\text{A} \) 5.5 MeV \(^3\text{He}\) beams through a 3.2 mm diameter collimator. These beams struck spinning targets to reduce the target heating; the recoiling \(^8\text{Li}\) and \(^8\text{B}\) nuclei were implanted into 11.1-mm-diameter carbon catcher foils located directly behind the target. The catcher foils for the \(^8\text{Li}\) and \(^8\text{B}\) runs had nominal thicknesses of 10 and 20 \( \mu\text{g/cm}^2 \), respectively. The target for the \(^8\text{Li}\) run consisted of \(^\text{nat}\text{LiF}\) evaporated onto the downstream side of a copper foil for better heat conductivity.

The catcher foils were attached to the ends of a 37-cm-radius arm that rotated 180° to cycle the foils from the bombardment station to the counting station and back again. Targets were bombarded for 2.00 s. The beam was then switched off by a fast magnet during the rotation and counting periods, which had durations of 0.15 and 5.00 s, respectively. The catcher foils were viewed on opposite sides by a pair of detectors (E counters) that were tightly collimated (\( \Delta \Omega/4\pi = 2.2 \times 10^{-3} \)) to minimize the energy summing of the \( \alpha \)’s with the preceding \( \beta \)’s.

Two additional 75-\( \mu \text{m} \)-thick \( \text{Si}\) surface barrier (SSB) detectors (E counters) that were tightly collimated (\( \Delta \Omega/4\pi = 2.2 \times 10^{-3} \)) to minimize the energy summing of the \( \alpha \)’s with the preceding \( \beta \)’s.

B. Delayed-\( \alpha \) data collection

The two E counters and the two 45° counters fed timing single-channel analyzers with thresholds set just above the noise. Data acquisition was triggered by an output of any of the single-channel analyses, in which case all detectors were read out. Counting rates were low enough that random coincidences were negligible. For each event we recorded the time interval between the start of the counting cycle, i.e., the arrival of a catcher foil at the counting station, and a signal in any one of our detectors. This was done by reading a 1 kHz clock that was zeroed every time fresh radioactivity was positioned at the counting station.

III. CALIBRATIONS

A. Energy scale and detector response

We placed special emphasis on the energy calibration. The linearity and zero point of our energy scale were determined with a precision pulser. The absolute scale factor was calibrated with high-resolution, spectroscopic-grade \(^{148}\text{Gd}\), \(^{239}\text{Pu}\), and \(^{241}\text{Am}\) \( \alpha \) sources [20] that were periodically placed in front of the detectors without breaking vacuum. The energies of these \( \alpha \) decays were taken from Ref. [21].

We fitted the \( \alpha \) source peaks (see Fig. 2) with an analytic function consisting of a Gaussian folded through two low-energy exponential tails:

\[
R(E, E') = \sum_{i=1}^{2} A_i \exp \left[ \frac{(E - E') - \sigma^2}{2\lambda_i^2} \right] \times \text{erfc} \left( \frac{E - E' + \sigma^2/\lambda_i}{\sqrt{2}\sigma} \right),
\]

where \( E \) and \( E' \) are the nominal and observed energies of the peak, \( \sigma = \text{FWHM}/(8 \log 2) \), \( \lambda_{1(2)} \) are exponential decay lengths (\( \lambda_1 < \lambda_2 \)), and \( \text{erfc} \) is the complement of the incomplete error function. The normalization coefficients are \( A_1 = 1/(1 + r) \) and \( A_2 = r/(1 + r) \) with \( r \) the relative area of tail 2 compared with tail 1. We used two measures of the peak position: \( E_c \) and \( E_{1c} \). The centroid of the entire peak is \( E_c = E - A_1\lambda_1 - A_2\lambda_2 \), while the centroid of the peak including the effect of tail 1 alone is \( E_{1c} = E - A_1\lambda_1 \). We found (see below) that tail 2 depended much more strongly on

![](https://example.com/filename.png)  
**FIG. 1.** (Color online) Schematic diagram of our apparatus.

![](https://example.com/filename.png)  
**FIG. 2.** (Color online) Lineshape fit of the \(^{148}\text{Gd}\) peak in detector 1. The dashed curves show the contributions of the short and long tails, while the solid curve shows their sum.
FIG. 3. Alpha particle spectra and lineshape fits for the $^{148}$Gd and $^{239}$Pu and $^{241}$Am sources. The $^{241}$Am line and the $^{239}$Pu line are triplets.

FIG. 4. Source thickness measurement showing the $^{148}$Gd peak position as a function of the source rotation angle. These data were taken with detector 1.

E counters were quite different: $58.7 \pm 4.9$ and $41.2 \pm 1.8 \mu g/cm^2$, respectively, where we assumed that the dead layers were gold. The nominal gold dead layer specified by the manufacturer is $40 \mu g/cm^2$.

We used the $\alpha$-source data, along with the zero-energy point determined with a precision pulser, to establish the energy calibration. The actual $\alpha$ energies deposited in the detectors, accounting for energy losses in the sources and detector dead layers, were fitted as linear functions of the $E_c$ positions of the peaks (which had very small uncertainties). The fit residuals were $\pm 3.0$ and $\pm 2.7$ keV for detectors E1 and E2, respectively (see Fig. 6). Figure 7 shows that the gain of our system was quite stable, typically varying by $0.04\%$ over 7 days of data taking. We therefore assign systematic errors of $\pm 3$ keV and $\pm 6$ keV to our $\alpha$ particle and excitation energies, respectively.

B. Catcher foil thickness

A PIN detector and $^{148}$Gd alpha source permanently mounted in the vacuum chamber allowed us to periodically monitor the catcher foil thickness by rotating the arm by $90^\circ$ to place the foil between the source and the PIN detector. Typical measured foil thicknesses for the $^8$Li ($^8$B) run were $14.5 (23.5) \mu g/cm^2$, increasing by $4.5\% (2.8\%)$ after 18 h of bombardment.

FIG. 5. Measurement of the dead layer of detector 1, showing the $^{148}$Gd peak position as a function of the source position.
IV. DATA REDUCTION AND ANALYSIS

A. Decay-time spectra

We generated E-counter decay-time spectra by gating on pulse heights just above the valley between the falling β background and the rising α’s and sorted the results according to the time signal from the 1 kHz clock that was reset every bombardment cycle. Results are shown in Fig. 8. These time spectra were fitted with a decaying exponential plus a constant to account for any long-lived β activities. The best-fit exponential time constants were in good agreement with tabulated values [22] for 8Li and 8B decays.

B. Delayed-α energy spectra

We sorted the delayed-α data into early (first 2 s of the counting period) and late (last 2 s of the counting period) spectra and subtracted the late-time spectra from the early-time spectra to minimize the low-energy background from long-lived β activity. We then used the β events in our E counters in coincidence with α’s in our 45° counters to subtract the low-energy β backgrounds from the decay of 8Li and 8B themselves. The β-subtracted spectra were first gain shifted to place the pulser peak in the same channel for all runs (the pulser peaks were quite stable, drifting by no more than 0.02% over the course of week long runs). We then converted channel numbers into delayed-α energies deposited into the counter by using the calibrations based on a linear fit of the zero-energy point and the three α-source peaks, taking into account the source thicknesses and detector dead layers. We then corrected for the energy loss of the delayed α’s in the catcher foil and detector dead layer. Finally, individual runs for each detector were added and then rebinned into excitation-energy spectra with 100 keV-wide bins, taking into account that the 8Be ground state is unbound by 92 keV. The results are shown in Fig. 9.

C. Analysis of the delayed-α spectra

We analyzed our excitation-energy distributions by using essentially the same procedure (and notation) we used in Ref. [11]. We employed Warburton’s version [14] of the
However, the mean energy shifts are almost negligible—in $^{8}\text{B}$ ($^{6}\text{Li}$) decay $\langle \Delta E \rangle$ is 3.3 (2.9), 2.1 (1.8), and 1.1 (0.9) keV for $^{8}\text{Be}$ excitation energies of 3, 6, and 9 MeV respectively—and were ignored in our analysis. The folded $R$-matrix function was then fitted to our data, the adjustable parameters being the energy and width of the 3-MeV level, the width of the background level, the GT matrix elements feeding the 3 MeV level ($M_1$), the $T=0$ component of the doublet ($M_{2+3}$), and the background state ($M_0$). The energies and widths of the 16 MeV doublet levels were fixed at their accepted values [22]. The fits to our $^{8}\text{Li}$ and $^{8}\text{B}$ delayed-$\alpha$ spectra are shown in Fig. 9, and their $R$-matrix parameters are listed in Tables II and III below.

### D. Final-state distributions and their uncertainties

Our $^{8}\text{Li}$ and $^{8}\text{B}$ $\beta$-decay final-state distributions as functions of $E_x$ in $^{8}\text{Be}$ are shown in Table I and in Figs. 10, 11, and 12 below. These were computed from the best-fit $R$-matrix parameters given in Table III below. The uncertainties have two sources: errors arising from the $R$-matrix parametrization of our spectra, and uncertainties in the energy calibration of the spectra.

For a fixed $R$-matrix matching radius, the uncertainty from the $R$-matrix parametrization (computed by using the full covariance matrix of the fit parameters) was negligible compared with the systematic error described below. (The different $R$-matrix parameters were highly correlated, so one cannot infer the uncertainties in the final-state distribution from the diagonal covariance elements alone.) However, a significant error from the $R$-matrix parametrization arises because of the ambiguity in the proper choice of matching radius. We chose a reasonable range of matching radii to be between $R^{-} = 4.2$ fm and $R^{+} = 4.8$ fm by requiring the total $\chi^2$ not to vary more than 1 unit from its minimum value at $R^{0} = 4.5$ fm. We computed the corresponding final-state distributions $(dN/dE_x)_{R^-}$ and $(dN/dE_x)_{R^+}$ and compared...
them with \((dN/dE_x)_R\) to obtain the \(R\)-matrix systematic uncertainty. This systematic error is in principle asymmetric because both \((dN/dE_x)_{R^-}\) and \((dN/dE_x)_{R^+}\) could lie on the same side of \((dN/dE_x)_R\), the final-state distribution inferred from the \(R = 4.5\) fm analysis (recall that all three final-state distributions must have the same area).

The energy-calibration systematic error was found by rebinning our delayed-\(\alpha\) data with high and low energy calibrations differing by \(\pm 1\sigma\) from the central calibration, taking account of the full covariance matrix from fitting the energy-calibration data. We then made \(R\)-matrix analyses of these high and low spectra, which again had the same

<table>
<thead>
<tr>
<th>(E_x)</th>
<th>(dN/dE_x \pm \sigma_E \pm \sigma_R)</th>
<th>(E_x)</th>
<th>(dN/dE_x \pm \sigma_E \pm \sigma_R)</th>
<th>(E_x)</th>
<th>(dN/dE_x \pm \sigma_E \pm \sigma_R)</th>
<th>(E_x)</th>
<th>(dN/dE_x \pm \sigma_E \pm \sigma_R)</th>
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<tbody>
<tr>
<td>0.2</td>
<td>0.005 (\pm 0.000 \pm 0.000)</td>
<td>0.4</td>
<td>0.088 (\pm 0.001 \pm 0.002)</td>
<td>0.6</td>
<td>0.472 (\pm 0.004 \pm 0.013)</td>
<td>0.8</td>
<td>1.534 (\pm 0.010 \pm 0.036)</td>
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<tr>
<td>0.4</td>
<td>1.467 (\pm 0.452 \pm 0.452)</td>
<td>4.6</td>
<td>127.239 (\pm 0.389 \pm 0.436)</td>
<td>4.8</td>
<td>111.598 (\pm 0.345 \pm 0.396)</td>
<td>5.0</td>
<td>98.735 (\pm 0.315 \pm 0.345)</td>
</tr>
<tr>
<td>0.6</td>
<td>8.6 (\pm 18.673 \pm 0.093 \pm 0.059)</td>
<td>9.0</td>
<td>15.638 (\pm 0.088 \pm 0.058)</td>
<td>9.2</td>
<td>14.290 (\pm 0.085 \pm 0.059)</td>
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<td>13.043 (\pm 0.082 \pm 0.059)</td>
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<tr>
<td>0.8</td>
<td>12.8 (\pm 2.077 \pm 0.025 \pm 0.019)</td>
<td>13.0</td>
<td>1.815 (\pm 0.024 \pm 0.017)</td>
<td>13.4</td>
<td>1.366 (\pm 0.021 \pm 0.020)</td>
<td>13.6</td>
<td>1.175 (\pm 0.019 \pm 0.020)</td>
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<td>1.0</td>
<td>14.0 (\pm 0.852 \pm 0.017 \pm 0.018)</td>
<td>1.2</td>
<td>29.640 (\pm 0.052 \pm 0.265)</td>
<td>1.6</td>
<td>58.426 (\pm 0.233 \pm 0.073)</td>
<td>2.0</td>
<td>50.945 (\pm 0.084 \pm 0.202)</td>
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<td>10.0 (\pm 9.843 \pm 0.070 \pm 0.050)</td>
<td>1.4</td>
<td>71.102 (\pm 0.261 \pm 0.168)</td>
<td>1.8</td>
<td>8.935 (\pm 0.065 \pm 0.046)</td>
<td>2.0</td>
<td>8.097 (\pm 0.061 \pm 0.042)</td>
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<td>1.4</td>
<td>14.2 (\pm 0.717 \pm 0.016 \pm 0.017)</td>
<td>1.6</td>
<td>64.342 (\pm 0.248 \pm 0.112)</td>
<td>1.8</td>
<td>9.599 (\pm 0.014 \pm 0.017)</td>
<td>2.0</td>
<td>9.459 (\pm 0.013 \pm 0.016)</td>
</tr>
<tr>
<td>1.8</td>
<td>14.4 (\pm 0.446 \pm 0.011 \pm 0.013)</td>
<td>2.0</td>
<td>53.202 (\pm 0.218 \pm 0.067)</td>
<td>2.4</td>
<td>10.4 (\pm 5.963 \pm 0.047 \pm 0.042)</td>
<td>2.6</td>
<td>11.0 (\pm 5.363 \pm 0.043 \pm 0.042)</td>
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<td>48.556 (\pm 0.201 \pm 0.060)</td>
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<td>44.394 (\pm 0.185 \pm 0.050)</td>
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<td>40.645 (\pm 0.170 \pm 0.072)</td>
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<td>4.836 (\pm 0.053 \pm 0.039)</td>
<td>3.4</td>
<td>6.615 (\pm 0.052 \pm 0.040)</td>
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<td>5.963 (\pm 0.047 \pm 0.042)</td>
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<tr>
<td>3.2</td>
<td>10.5 (\pm 0.326 \pm 0.010 \pm 0.012)</td>
<td>3.8</td>
<td>7.326 (\pm 0.056 \pm 0.039)</td>
<td>4.0</td>
<td>9.599 (\pm 0.014 \pm 0.017)</td>
<td>4.2</td>
<td>9.459 (\pm 0.013 \pm 0.016)</td>
</tr>
<tr>
<td>4.0</td>
<td>14.6 (\pm 0.495 \pm 0.013 \pm 0.016)</td>
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<td>10.5 (\pm 0.326 \pm 0.010 \pm 0.012)</td>
<td>4.6</td>
<td>9.459 (\pm 0.013 \pm 0.016)</td>
</tr>
</tbody>
</table>

FIG. 11. Comparison of the \(^8\)B final-state continuum shapes we extract from \(R\)-matrix analyses of our \(\alpha\) spectrum (solid curve) and previous spectrum of Wilkinson and Alburger [7] (dashed). Both shapes have been normalized to the same total area. The agreement is remarkably good.

FIG. 12. (Color online) Comparison of \(^8\)B final-state continuum shapes from our work (narrow black curve) to recent results from Ortiz et al. [16] (broad blue band) and Winter et al. [18] (narrow light-red band). The shaded bands show \(\pm 1\sigma\) error bands of the \(R\)-matrix final-state distribution. All shapes have been normalized to the same total area. Our results and those of Winter et al. overlap to such a degree that they are nearly indistinguishable in this plot. Our \(\pm 1\sigma\) error band was computed as described in the text. We generated the error band for the Ortiz et al. data from our \(R\)-matrix fits to their \(\pm 1\sigma\) spectra. The Winter et al. error band was generated by distorting the energy scale with a multiplicative factor of 1 \(+ (0.275\%)\) added in quadrature with a constant offset of 3 keV as described in Ref. [18].
TABLE II. Comparison of $^8$B $\beta$-decay $R$-matrix parameters extracted from this and other recent works. Energies and widths are in keV, GT matrix elements are in $\mu N$. $E_2$ and $E_3$ and the reduced widths $\gamma^2_2$ and $\gamma^2_3$ were fixed at values derived from Ref. [22]. All analyses use the $R$-matrix formalism of Ref. [11].

<table>
<thead>
<tr>
<th></th>
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</thead>
<tbody>
<tr>
<td>$E_1$</td>
<td>3037 ± 5 1</td>
<td>3063 ± 4</td>
<td>3012 ± 12</td>
<td>3043</td>
</tr>
<tr>
<td>$\gamma^2_2$</td>
<td>1075 ± 6 3</td>
<td>1148 ± 6</td>
<td>1179 ± 18</td>
<td>1087</td>
</tr>
<tr>
<td>$M_1$</td>
<td>-0.1449 ± 0.0006 ± 0.0001</td>
<td>-0.1507 ± 0.0005</td>
<td>-0.1476 ± 0.0015</td>
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<td>16626</td>
<td>16626</td>
<td>16626</td>
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<tr>
<td>$\gamma^2_3$</td>
<td>10.89</td>
<td>10.89</td>
<td>10.89</td>
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<td>7.46</td>
<td>7.42</td>
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<tr>
<td>$M_{2+3}$</td>
<td>3.152 ± 0.158 ± 0.009</td>
<td>2.377 ± 0.100</td>
<td>3.555 ± 0.430</td>
<td>2.423</td>
</tr>
<tr>
<td>$E_4$</td>
<td>37000</td>
<td>37000</td>
<td>37000</td>
<td>37000</td>
</tr>
<tr>
<td>$\gamma^2_5$</td>
<td>7099 ± 276 ± 12</td>
<td>5660 ± 242</td>
<td>7756 ± 681</td>
<td>5619</td>
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<tr>
<td>$M_4$</td>
<td>-0.254 ± 0.022 ± 0.001</td>
<td>-0.133 ± 0.020</td>
<td>-0.332 ± 0.025</td>
<td>-0.132</td>
</tr>
</tbody>
</table>

The total number of counts as the spectrum based on the best-fit energy calibration. Energy-calibration systematic errors were obtained from the $(dN/dE_\gamma)_e^+$ and $(dN/dE_\gamma)_e^-$ final-state distributions inferred from $R$-matrix fits to the high and low spectra. These systematic errors were again, in principle, asymmetrical. For simplicity, we adopted symmetric measures of the $R$-matrix and energy-calibration uncertainties, in each case the larger of the two ±1σ asymmetric uncertainties. We recommend that the $R$-matrix and energy-calibration uncertainties be added in quadrature to obtain the total uncertainty in the final-state distribution.

E. Comparison with previous delayed-α final-state distributions

1. Wilkinson-Alburger data

Figures 10 and 11 compare our new measurements of the final-state distributions in $^8$Li and $^8$B decays with the corresponding quantities from our reanalysis [11] of the Wilkinson-Alburger thin-catcher data. We cannot show realistic error bands for this comparison because the covariance matrix of the energy calibration for the Wilkinson-Alburger data is not available. Nevertheless, the reanalyzed Wilkinson-Alburger data agree well with our results, being well within the ±60 keV uncertainty in $E_\alpha$ (±30 keV in $E_\alpha$ quoted in Ref. [14]). The concordance of the $^8$B results is particularly impressive. During the course of this work, we discovered an error in the fitting program used in Ref. [11]. Correcting this error gave $R$-matrix parameters, shown in Table II, that differed slightly from the values quoted in Ref. [11].

2. Recent Notre Dame and LBL-ANL results

Two new high-precision measurements of the the $^8$B delayed-α continuum have recently been reported. Although both results were based on detecting the summed energies of the two α’s, so that the first-order effects of lepton recoil were absent, the experimental techniques were quite different. Ortiz et al. [15,16] at Notre Dame produced the $^8$B with the same reaction we used, but mounted their catcher foils on a chain that carried them into a strong magnetic field that swept away the positrons, avoiding the positron summing problem. But

![Figure 13](image.png)

FIG. 13. (Color online) Detailed comparison of our $^8$B final-state continuum shape (narrower black) to that of Winter et al. [18] (broader red). Our central value has been subtracted from both data sets. The shaded regions show the ±1σ error bands of the two data sets.
experiments, we had to account for the first-order effects of lepton recoil, but this can be done with little uncertainty. As a result we quote the smallest errors of the three measurements.

Figure 12 and Table II compare our \(^8\)B final-state distribution to LBL-ANL and Notre Dame results. Our results agree beautifully with Winter et al. Both we and Winter et al. find that the peak in the final-state distribution is narrower and occurs about 60 keV higher in \(E_s\) than in the Notre Dame result. Figure 13 shows a detailed comparison of our work to that of Winter et al.

V. CONCLUSION

The increasingly precise solar neutrino data, particularly from the SNO detector, place a premium on high-quality measurements of the unoscillated or intrinsic spectrum of \(^8\)B neutrinos, as the oscillation parameters are sensitive to the energy-dependent distortion of the spectrum shape. This in turn requires accurate knowledge of the shape of the very broad \(^8\)Be final-state continuum fed in the decays.

Two measurements [15,17] of this continuum were recently reported, with comparably high quoted precisions. Both measurements detected the summed energies of the two alphas, so the first-order effects of lepton recoil were absent. However, the two experiments involved very different systematic effects.

This paper reports an entirely different measurement based on measuring the energies of only one of the two alphas, which, of course, had completely different systematic effects from Refs. [15,17]. Beta summing was negligible, but we had to account for lepton recoil. Our main systematic effects were due to the energy resolutions of the detectors and calibration sources. Our data were taken and analyzed after the publication of Ref. [15] but before the appearance of Ref. [17]. The delay in publication resulted from our slowness in computing the \(^8\)B neutrino spectrum from our final-state continuum measurements. Winter et al.’s calculation of the neutrino spectrum, including a careful treatment of recoil-order effects, has superseded our attempts, so we report only our measurement of the final-state continuum.

The outstanding agreement between our work and that of Refs. [17,18] and our reanalysis of Ref. [7]) is remarkable. For the first time, there is a concordance, at a high level of precision, between three independent measurements of the final-state continuum in \(^8\)B decay performed by two entirely different techniques. This provides a robust basis for calculating the intrinsic neutrino spectrum from \(^8\)B activity in the Sun. Because of the excellent agreement of our final-state distribution and that of Refs. [17,18], there is no need for us to present a neutrino spectrum derived from our work, as it would not differ significantly from that of Winter et al.

Finally, it is noteworthy that the parameters of the lowest \(2^+\) level in \(^8\)Be inferred from our \(^8\)Li and \(^8\)B delayed-\(\alpha\) data are consistent within errors, and both are quite close to the parameters inferred from the \(\alpha + \alpha\) phase shifts (see Table III). The excitation energies agree within the quoted uncertainties. Although the width extracted from scattering data is about 100 keV less than the \(\beta\)-decay value, this could well be an artifact of a simplifying assumption inherent in the \(R\)-matrix description of the delayed-\(\alpha\) spectra, namely that the Gamow-Teller matrix elements are strictly independent of excitation energy. This agreement, which is now substantially closer than that observed by Warburton [14], further weakens Barker’s argument for a low-lying intruder state [12,13] in \(^8\)Be.

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