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Inelastic ($1^+ \to 0^+$) electromagnetic form factor of ^{6}Li from three-body models

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Within the context of three-body (alpha particle plus two nucleons) models, and by assuming that the ground state of ^{6}He (0^+) and the lowest 0^+ excited state of ^{6}Li are members of the same isospin multiplet, the $M1$ transition form factor of ^{6}Li for the inelastic electron scattering from the ground state ($J^\pi T = 1^+0$) to the $\omega = 3.56$-MeV excited state ($J^\pi T = 0^+1$) is calculated. From the inelastic form factor evaluated at $|\mathbf{q}| = \omega$, the radiative width for the deexcitation of the 3.56-MeV 0^+ state of ^{6}Li is given. It is suggested that the calculations of this work serve as a starting point for future investigations where refinements are made to the three-body models and more details of the electromagnetic interaction, e.g., meson exchange currents, are taken into account.

I. INTRODUCTION

In a previous paper[1] (hereafter referred to as I) on the elastic electromagnetic form factors of ^{6}Li within the context of three-body models, we mentioned that at low-excitation energies the $A = 6$ nuclei can be satisfactorily thought of as being made up of an alpha particle and two nucleons. We then used the ground-state wave function of ^{6}Li derived by Lehman, Rai, and Ghovanlou[2] (LRG) to calculate the elastic electromagnetic form factors. The calculation of the inelastic transition form factor of the ^{6}Li nucleus from the ground state to the 0^+ (3.56 MeV) excited state complements our work on the elastic electron process. By assuming that the ground state of ^{6}He (0^+) and the 3.56 MeV excited state of ^{6}Li(0^+) are members of the same isospin multiplet, we can use the ^{6}He ground-state wave function given by Ghovanlou and Lehman[3] (GL) as the final-state wave function of the inelastic process. This approach has been successful in prediction[4] of the semileptonic β decay of ^{6}He. Our goals in undertaking such a calculation are similar to those in I: (1) to explain the physics behind the observed diffraction minimum in the inelastic form factor, and in particular assess the role of the underlying two-body αN and NN interactions on the inelastic form factor; (2) to establish a foundation for further, more sophisticated, but equally consistent, nonrelativistic calculations.

We shall first briefly review the existing literature on the subject. Then, in Sec. II, the derivation of the relevant equations is presented. Our results are given in Sec. III and discussion of the results follows in Sec. IV. The paper closes in Sec. V with our conclusions.

Besides calculation of the elastic form factors of ^{6}Li (see I for more details), Krasnopol'ski *et al.*[5] also have used their variational approach as applied to three-body models for the calculation of the inelastic form factor, and included the following three configurations for the excited-state wave function: $\{\lambda = l = L = 0, S = 0\}$, $\{\lambda = l = L = 1, S = 1\}$, and $\{\lambda = l = 2, L = 0, S = 0\}$, where λ and l refer to the angular momenta of the two-nucleon system and the α particle relative to the center of mass of the two-nucleon system, respectively. L is the total orbital angular momentum, and S is the total spin. With this wave function for the ^{6}Li 0^+ (3.56 MeV) state and the same ground-state wave function used in their elastic form factor work, the authors of Ref. 5 find that the experimental data for the inelastic form factor are explained rather well up to the first minimum, but for higher values of q the data are underestimated by the theory. Some of this underestimation may be attributed to the meson-exchange currents which are briefly discussed below. However, the magnitude of mesonic effects is not known, and it is not at all clear whether or not they will compensate for the underestimation.

Having measured the ($1^+ \to 0^+$) transition form factor, Bergstrom[6] tried to reproduce the experimental values by using phenomenological cluster models. Both ^{3}H-^{3}He and α-d cluster models predict a value of 8.30 eV for the radiative width for the deexcitation of the 0^+ state which is in good agreement with the experimental value ($\Gamma_{\rm expt} = 8.16 \pm 0.19$ eV).[7] The inelastic form factor is treated as a pure spin-flip transition even though, in general, convection-current terms are present. The oscillator parameters of the wave function were taken to be the same for the ground state and the excited state. The ^{3}H-^{3}He model reproduces the first minimum near the expected position, but overestimates the form factor in the first maximum region. The α-d model, on the other hand, predicts the minimum at a smaller q value and no variation of parameters consistent with the rms charge radius moves it to the right place. If the radius constraint is removed, agreement may be achieved for smaller values of the rms charge radius, and the form factor, although in agreement with the ^{3}H-^{3}He model up to the second maximum, overestimates it by a factor of 2 in that region. Both the ^{3}H-^{3}He and α-d cluster models overestimate the form factor near the peak of the first maximum. As noted by Bergstrom, this seems to be inherent in all the oscillatorlike wave functions.

II. DERIVATION OF EQUATIONS

The transverse transition form factor is experimentally extracted from the cross section for inelastic electron scattering from ^{6}Li. In order to express the cross section in a general form and to see how the transverse form factor is defined from the nuclear current, we begin with the relativistic expression for the transition current of the inelastic process. Assuming that the nucleus is a "spin-1" object which changes its spin value in the inelastic scattering process to zero spin, the nature of the current is limited further by imposing Hermiticity, current conservation, parity, and time-reversal invariance. These requirements reduce the number of independent transition form factors to only one:

$$J^{\mu}_{fi}=[Ze/(V\sqrt{2\mathcal{E}_i 2\mathcal{E}_f})]F_1(q_\rho^2)\epsilon^{\mu}_{\sigma\lambda\rho}q^\sigma p^\lambda S_i^\rho , \quad (1)$$

where $\mathcal{E}_i$ and $\mathcal{E}_f$ are the initial and final energies of the nucleus, V is the normalization volume, and $\epsilon_{\mu\sigma\lambda\rho}$ is the completely antisymmetric tensor of rank 4. The p^λ and q^σ are given in terms of the initial (p_i^μ) and final (p_f^μ) four-momenta of the nucleus by $(p_i^\lambda+p_f^\lambda)$ and $(p_f^\sigma-p_i^\sigma)$, respectively. S_i designates the polarization vector of a freely moving particle of "spin 1."

Taking the target nucleus to be unpolarized, we can use

$$\sum_{\substack{\text{polarization}\\ \text{of the}\\ \text{nucleus}}} J^{\mu}_{fi}J^{\nu\dagger}_{fi}=\frac{Z^2e^2}{V^2 2\mathcal{E}_i 2\mathcal{E}_f}|F_1(q_\rho^2)|^2\{-q_\rho^2 p^\mu p^\nu-p_\rho^2 q^\mu q^\nu+(p_\rho q^\rho)(p^\mu q^\nu+p^\nu q^\mu)-[(p_\rho q^\rho)^2-p_\rho^2 q_\rho^2]g^{\mu\nu}\} , \quad (2)$$

where $q_\rho^2\equiv q_\rho q^\rho$. The usual form[6] for the inelastic cross section in the laboratory frame assuming one-photon exchange is then given by

$$\frac{d\sigma}{d\Omega_{e_f}}=\left[\frac{Z^2\alpha^2\cos^2(\theta/2)}{4E_i^2\sin^4(\theta/2)}\right]\left[\frac{1}{1+\frac{2E_i}{m}\sin^2(\theta/2)}\right] \times|F_T(q_\rho^2)|^2\left[-\frac{q_\rho^2}{2q^2}+\tan^2(\theta/2)\right] , \quad (3)$$

where $q^2\equiv\mathbf{q}\cdot\mathbf{q}$. In Eq. (3), θ is the angle between the incident and outgoing momenta, α is th fine-structure constant, E_i is the incident electron energy, and m is the ^{6}Li mass. F_T can then be related to F_1 by

$$F_T(q_\rho^2)=\sqrt{2/3}\,2|\mathbf{q}|F_1(q_\rho^2) . \quad (4)$$

In any other frame, the cross section will have a longitudinal form factor piece as well as a transverse one. However, in that case the two form factors are not independent.

The derivation of the inelastic form factor within the nonrelativistic framework of the three-body model begins by noting that only the transverse part of the constituent current enters in the transition operator. (The arguments are exactly the same as the ones given for the elastic magnetic form factor of ^{6}Li in Ref. 1.) Therefore, one needs to calculate the matrix element:

$$\langle {}^6\text{He}(0^+)|\epsilon_\lambda\cdot\mathbf{J}|{}^6\text{Li}(1^+)\rangle ,$$

where

$$\epsilon_\lambda\cdot\mathbf{J}=i\sum_{j=1}^{3}e^{i\mathbf{q}\cdot\mathbf{r}_j}\left[e_jF^j_{\text{ch}}(q^2)\frac{P^{[1]}_{j\lambda}}{M_j}+\bar{\delta}_{3j}\mu_j q\lambda\frac{F^j_{\text{mag}}(q^2)}{2M_j}\sigma^{[1]}_{j\lambda}\right] . \quad (5)$$

Here, $\bar{\delta}_{3j}=(1-\delta_{3j})$, $F^j_{\text{ch}}(q^2)$ and $F^j_{\text{mag}}(q^2)$ refer to the charge and magnetic form factors and M_j is the mass of the jth particle. The ^{6}He and ^{6}Li wave functions are those of GL and LRG that have been derived on the basis of three-body models of the $A=6$ system.[3,2] In the absence of any physical quantity by which the transverse inelastic form factor can be normalized at $|\mathbf{q}|=0$, one may set the scale of the form factor by comparing the cross section derived from the nonrelativistic (model-dependent) nuclear current, in the impulse approximation, with the general expression in Eq. (3). With this in mind, the exponential factor in the operator of Eq. (5) is expanded (see I) to give the multipole contributions to the inelastic form factor.

An important distinction between this case and the elastic magnetic form factor is that here the final-state wave function has total angular momentum $J=0$ and is antisymmetric with respect to the exchange of the two nucleons, excluding the isospin part.

The above considerations plus that of parity and time-reversal invariance eliminate all but a few terms, which do not vanish by any other symmetry arguments. The form factor in momentum space is then given by

$$F_T(q^2)=\frac{\sqrt{2/3}}{3}\left[\frac{iq}{2M\sqrt{3}}\right]\left[\frac{-i\sqrt{6}}{q\sqrt{4\pi}}[F^p_{\text{ch}}(q^2)-F^n_{\text{ch}}(q^2)][\tilde{\Psi}^{[0]}(\mathbf{p},\mathbf{k})|[Y^{[1]}(\hat{q})\times P^{[1]}_1]^{[1]}|\Psi^{[1]}(\mathbf{p}+2\mathbf{q}/3,\mathbf{k}-\mathbf{q}/2)] + [\mu_pF^p_{\text{mag}}\cdot(q^2)-\mu_nF^n_{\text{mag}}\cdot(q^2)]\frac{1}{\sqrt{4\pi}} \times[\tilde{\Psi}^{[0]}(\mathbf{p},\mathbf{k})|[Y^{[0]}\times\sigma^{[1]}_1]^{[1]}-\frac{1}{\sqrt{2}}[Y^{[2]}(\hat{q})\times\sigma^{[1]}_1]^{[1]}|\Psi^{[1]}(\mathbf{p}+2\mathbf{q}/3,\mathbf{k}-\mathbf{q}/2)]\right] , \quad (6)$$

where we note that the argument of F_T in the nonrelativistic framework is q^2 rather than q_ρ^2. The radiative width of the ($0^+ \rightarrow 1^+$) transition can be found from the transition amplitude to be

$$\Gamma_{0^+ \rightarrow 1^+} = 54\alpha\omega |F_T(|\mathbf{q}|=\omega)|^2 , \quad (7)$$

where $\omega = 3.562$ MeV is the excitation energy of the 0^+ state in ^{6}Li and α is the fine-structure constant. Detailed expressions for the wave functions can be found in Refs. 2 and 3.

In the present calculation the contributions of the meson-exchange currents are ignored. However, a power-counting argument shows that the lowest-order meson-exchange-current effects can be included in a nonrelativistic calculation such as this one in a consistent fashion. The lowest-order isovector meson-exchange current contribution to the $1^+ \rightarrow 0^+$ inelastic transition form factor is of $O(1/M)$.[8,9] Given the fact that the electromagnetic current operator used in this calculation is of $O(1/M)$ and the nonrelativistic wave functions used are of $O(1)$, it is easy to see that the lowest-order mesonic effects can be consistently included in a nonrelativistic calculation of the transition form factor. Although a calculation of the mesonic effects will not be considered here, it is hoped that the present calculation will set the groundwork for later investigation of the magnitude of the mesonic effects for the $1^+ \rightarrow 0^+$ inelastic form factor of ^{6}Li within the framework of the three-body model.

III. RESULTS

For the benefit of the reader, we summarize the five potential models used in our calculation before the actual presentation of the results (for a full account see Refs. 2 and 3). In the "simple model," the tensor force is not included in the NN interaction and only the $P_{3/2}$ component of the αN interaction is present. The $P_{1/2}$, $P_{3/2}$, and $S_{1/2}$ components of the αN interaction are present in the remaining four models. This enables us to see the effects of the inclusion of the $S_{1/2}$ component of the αN interaction clearly. The $P_{1/2}$ component of the αN interaction is weak enough that any major observable difference between the simple model and the remaining four models is due to the presence of the $S_{1/2}$ part of the interaction in those models. Furthermore, the remaining models are labeled "full repulsive" or "projected bound state" according to the representation of the $S_{1/2}$ αN interaction by a fully repulsive potential or an attractive potential (from which a Pauli "forbidden" bound state has been removed by a projection method),[10] respectively. In order to assess the effect of the tensor force in the NN interaction, each of the above models is employed with (0%) or (4%) tensor component.

The general procedures for the algebra and numerical computation of the inelastic form factor are similar to the elastic form factors and are explained in Ref. 1. Here we simply state the results.

In this calculation of the inelastic form factor, we drop the convection part of the current operator because of its complexity relative to the expected smallness of its contribution. Here, as in the elastic magnetic form factor (see I), the contribution of

$$[\sigma^{[1]} \times Y^{[2]}(\hat{q})]^{[1]}$$

is found to be negligible, leaving only the term with

$$[\sigma^{[1]} \times Y^{[0]}(\hat{q})]^{[1]} .$$

In Table I, the values for the radiative width in the $0^+ \rightarrow 1^+$ transition in ^{6}Li are presented for different models. As mentioned in Sec. II [see Eq. (7)], the radiative width is related to the value of the form factor at $|\mathbf{q}| = \omega$ where $\omega = 3.562$ MeV or $|\mathbf{q}| = 0.0181$ fm^{-1}.

Plots of the inelastic form factor $F_T(q^2)$ for different models and for different fits (*np* or *nn*) (the labels "*np* best fit" or "*nn* best fit" refer to the type of NN interaction) in the ^{6}He wave function are presented in Figs. 1–3. All models with the exception of the simple model show diffraction minima between $q = 2.0$ and 2.5 fm^{-1}. The agreement in the second lobe with experimental data is somewhat better in comparison with the magnetic elastic form factor.

Experimentally the $1^+ \rightarrow 0^+$ inelastic form factor has been measured up to the value of $q \approx 3.0$ fm^{-1} by Bergstrom *et al.*[6] It exhibits a minimum at about $q \approx 1.20-1.50$ fm^{-1} and two maxima on either side of it at about $q \approx 0.6$ and 1.9 fm^{-1}. For the analysis of this work Bergstrom *et al.*'s results will be used (see Fig. 1). The parametrizations of the isovector nucleon form factors that appear in Eq. (6) were taken from the work of Höhler *et al.*[11]

IV. DISCUSSION

The plots presented in Figs. 1–3 reflect the following general features of the different models with respect to $F_T(q^2)$. (1) All models except for the simple model predict a diffraction minimum at q values approximately

TABLE I. Radiative width for de-excitation of the 0^+ state $\Gamma(0^+ \rightarrow 1^+) = 54\alpha\omega |F_T(|q|=\omega)|^2$ with $\omega = 3.562$ MeV and $\alpha = 1/137.04$. Experimental result (Ref. 7): $\Gamma(\text{expt}) = 8.16 \pm 0.19$ eV.

Model	$\Gamma_{0^+ \rightarrow 1^+}$ (eV)
Simple model	6.69
Repulsive 4% (*np* best fit)	6.65
Projected bound state 0% (*np* best fit)	6.94
Projected bound state 4% (*nn* best fit)	6.71
Projected bound state 4% (*np* best fit)	6.79

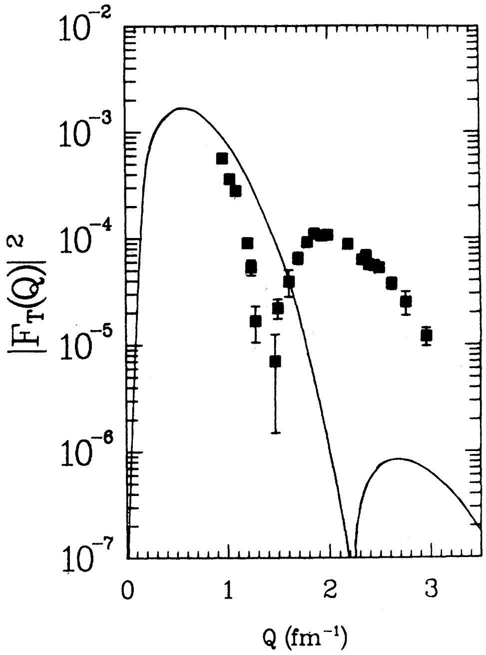

FIG. 1. The inelastic form factor ($1^+ \rightarrow 0^+$) of ^{6}Li for the projected model [(4%) for ^{6}Li, *np* best fit for ^{6}He], and the experimental data of Ref. 6.

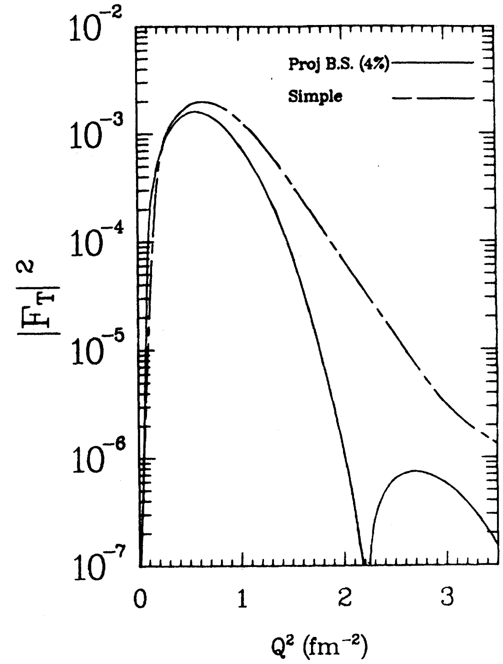

FIG. 3. The inelastic form factor ($1^+ \rightarrow 0^+$) of ^{6}Li. Comparison of the projected bound state (4%), and the simple models (the ^{6}He wave function with the *np* best fit).

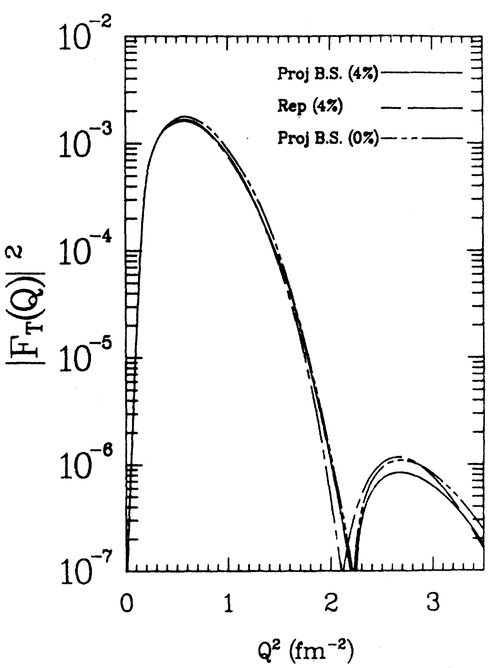

FIG. 2. The inelastic form factor ($1^+ \rightarrow 0^+$) of ^{6}Li. Comparison of the projected bound state (4%), projected bound state (0%), and the full repulsive (4%) models (the ^{6}He wave function with the *np* best fit).

0.8–1 fm^{-1} larger than the experiment. (2) The absence of the tensor force leads to slightly larger values of $|F_T(q^2)|^2$ in the region of the secondary maximum (Fig. 2). (3) The projected bound state and the full repulsive models, although basically indistinguishable in terms of their prediction of the shape of $|F_T(q^2)|^2$, are slightly distinguishable in terms of their prediction of the magnitude of $|F_T(q^2)|^2$ in the region of the secondary maximum (Fig. 2). (4) The diffraction minimum originates from the Pauli repulsion between the alpha particle and each nucleon as made manifest through the $S_{1/2}$ αN interaction. The importance of the $S_{1/2}$ interaction in the αN system for the production of the diffraction minimum is seen in Fig. 3, where it is observed that the simple model (the only model in which the $S_{1/2}$ component of the αN interaction is not included) does not render a diffraction minimum. Finally, though not shown in the figures, it is found that $F_T(q^2)$ is not sensitive to whether the $^1S_0$ *NN* interaction originates from the *nn* or *np* fits.

The radiative width values (Table I) for the deexcitation of the 0^+ (3.56 MeV) state of ^{6}Li lie below the experimental values by approximately 15–19%. This shows that even at low q values (where F_T is evaluated for the calculation of Γ) the agreement with experiment is not satisfactory.

What are the possible sources of discrepancies? As mentioned in Sec. II, in the calculation of the inelastic form factor of ^{6}Li one can in a consistent manner include the lowest-order meson-exchange current operator contributions. It is hoped that the present calculation would serve as a first step toward future calculations of the meson-exchange-current contributions. In particular, it would be interesting to see if the radiative width discrepancy is mainly due to the mesonic contributions or the convection-current contributions. Calculations of Dubach *et al.* show that the meson-exchange effects change the form-factor values at the tail of the second maximum and are only very important in the region of the second and the third maxima.[12] Based on their estimate, the values of Γ would not change significantly upon the inclusion of the meson-exchange-current operators. On the other hand, the phenomenological shell-model wave functions used by Dubach *et al.* are limited in that they assume the valence nucleons reside only in the p shell for the ground state and 3.562 MeV excited state of ^{6}Li, whereas the three-body model wave functions contain all possible j-j coupling components for the valence nucleons.[13] Therefore, one might expect a greater contribution from meson-exchange currents in the three-body framework than in the model of Dubach *et al.* Of course, it may be that both convection-current and meson-exchange-current contributions must be included before a better agreement with experiment is reached even for low q values of the inelastic form factor.

How do the present results compare to those of Krasnopol'ski *et al.*?[5] In the 0–1 fm^{-1} region of q, the results are comparable. However, beyond 1 fm^{-1}, the curve of Ref. 5 falls more rapidly to produce a diffraction minimum in $F_T(q)$ at roughly the value of q indicated by the data, i.e., $q \sim 1.3-1.4$ fm^{-1}. Moreover, their value of $F_T(q)$ at the secondary maximum is only an order of magnitude below the data rather than the 2 orders of magnitude in the present calculation. At this stage it is difficult to pinpoint the source of the differences. The authors of Ref. 5 do not give the details of the electromagnetic operator used. Also, unlike the present work, where the $J^\pi T = 0^+ 1$ excited state is a particle-stable bound state as in nature, the $0^+ 1$ excited state is unbound in the work of Ref. 5 and is approximated as a bound state (binding energy $+1.29$ MeV) in their variational framework. Though the authors of Ref. 5 limit the angular-momentum configurations in their wave functions, which we do not, they do have a better representation of the two-nucleon interaction through the Reid soft core potential than our rank-one separable forms. This latter aspect might explain some of the difference at the higher values of q. In fact, improvement of our models by use of higher-rank separable interactions to better represent the NN interaction, especially the short-range repulsion, should be the next step before consideration of the long-range one-pion-exchange-type meson-exchange currents. Of course, both the approach of Ref. 5 and the present one treat the alpha particle as elementary. After these aspects are carefully considered, if discrepancies between theory and experiment still persist, the role of the alpha particle's structure must be investigated.

V. CONCLUSION

Three-body wave functions of ^{6}He and ^{6}Li have been used to calculate the $1^+ \to 0^+$ (3.56 MeV) inelastic form factor of ^{6}Li. Once the underlying two-body interactions are parametrized, the calculations are performed without further parametrization at the three-body level. The following insights are extracted from the results of our calculation: (1) The $S_{1/2}$ component of the αN interaction is responsible for the production of the observed diffraction minimum. (2) The form-factor values are only slightly sensitive to the particular representation of the $S_{1/2}$ interaction in the αN subsystem. (3) The tensor force in the NN interaction has a definite effect on the form-factor values, especially in the region of the secondary maximum.

With these important insights with regard to the physics of the problem, we have removed some ambiguities about our understanding of the $1^+ \to 0^+$ inelastic process in ^{6}Li. We have also given a framework for further extensions of the present calculation so that the meson-exchange current contributions can be incorporated in an unambiguous and consistent fashion. Inclusion of the convection-current operator would push our quantitative understanding still further. These extensions, in addition to improvement of the underlying two-body interactions and inclusion of the Coulomb interaction, are the limits of our capabilities for the present, relatively sophisticated, nonrelativistic calculations.

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