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# Relations Among the Structure Functions of Deep-Inelastic Neutrino-Nucleon Scattering*

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Light-cone current commutators are used to derive relations between the structure functions of the vector-vector part of neutrino-nucleon scattering and the structure functions of the vector–axial-vector part.

Reasonable expressions for the commutator of currents restricted to a lightlike surface have been derived from the quark model.[1,2] These expressions involve bilocal operators and it is possible to derive sum rules which relate the structure functions of inelastic lepton-nucleon scattering to the matrix elements of these bilocal operators.[3] In particular, if the matrix elements of the bilocal operators are written out in terms of form factors, then, by using the light-cone version[1] of the Bjorken-Johnson-Low (BJL) limit,[4] the deep-inelastic limits of the structure functions become simply Fourier transforms of the form factors. In Ref. 3 it was found that the structure functions of a scattering process involving only vector currents were related to the form factors in a one-to-one manner.

A basic assumption of the light-cone-commutator method is that the commutators are invariant under SU(3)×SU(3) symmetry. Therefore the commutator of a (conserved) axial-vector current with a vector current involves the same bilocal operators as the commutator of two vector currents and the matrix element of the commutator involves the same form factors. Deep-inelastic neutrino-nucleon scattering, for example, will have structure functions for the vector–axial-vector (*VA*) cross terms as well as the structure functions for the vector-vector or axial-vector–axial-vector terms. However, the relations between the structure functions and the form factors will involve only the form factors from vector-vector scattering. Thus the structure functions for the *VA* cross terms are all related, through the form factors, to the structure functions of the *VV* terms.

In the case where the nucleon spin is summed over, the relation is well known. It was first derived from the parton model by Llewellyn Smith[5] and has been derived from light-cone commutators by several people.[2,6,7] In the present work we shall find the relations among the spin-dependent structure functions as well. These results are interesting in that they can be compared with other models, most of which give the same answer for the spin-independent vector-vector structure functions but differ in the spin-dependent and/or vector–axial-vector structure functions.[8] In addition, if electron scattering should verify the essential correctness of the light-cone-commutator approach then these relations can provide a handle on neutrino scattering.

There are four independent invariant amplitudes for vector-spinor → vector-spinor scattering in the forward direction, two spin-dependent and two spin-independent. For spinor-vector → spinor–axial-vector scattering there are also four but only one is independent of spin.[9] The time-ordered product of two currents $J_a^\mu(x)=V_a^\mu(x)-A_a^\mu(x)$ can be written as

$$
\begin{aligned}
T_{ab}^{\mu\nu}(p,q) &\equiv i\int d^4x\, e^{iq\cdot x}\langle p,s|(J_a^\mu(x)J_b^\nu(0))_+|p,s\rangle \\
&= -\left(g^{\mu\nu}-\frac{q^\mu q^\nu}{q^2}\right)T_L^{ab}(q^2,\nu)+\left(p^\mu p^\nu-\frac{\nu}{q^2}(p^\mu q^\nu+p^\nu q^\mu)+\frac{\nu^2}{q^2}g^{\mu\nu}\right)T_2^{ab}(q^2,\nu) \\
&\quad +i\epsilon^{\mu\nu\alpha\beta}s_\alpha q_\beta T_3^{ab}(q^2,\nu)+iq\cdot s\,\epsilon^{\mu\nu\alpha\beta}p_\alpha q_\beta T_4^{ab}(q^2,\nu)+\tfrac{1}{2}i\,\epsilon^{\mu\nu\alpha\beta}p_\alpha q_\beta T_5^{ab}(q^2,\nu) \\
&\quad +q\cdot s\left(g^{\mu\nu}-\frac{q^\mu q^\nu}{q^2}\right)T_6^{ab}(q^2,\nu)+q\cdot s\left(p^\mu p^\nu-\frac{\nu}{q^2}(p^\mu q^\nu+p^\nu q^\mu)+\frac{\nu^2}{q^2}g^{\mu\nu}\right)T_7^{ab}(q^2,\nu) \\
&\quad +\left[\left(s^\mu-\frac{q\cdot s}{q^2}q^\mu\right)\left(p^\nu-\frac{\nu}{q^2}q^\nu\right)+\left(p^\mu-\frac{\nu}{q^2}q^\mu\right)\left(s^\nu-\frac{q\cdot s}{q^2}q^\nu\right)\right]T_8^{ab}(q^2,\nu) \\
&\quad +if_{abc}\frac{1}{q^2}\Gamma_c(g^{\mu\nu}\nu-q^\mu p^\nu-q^\nu p^\mu)+if_{abc}\frac{1}{q^2}\Gamma_c^A(g^{\mu\nu}q\cdot s-q^\mu s^\nu-q^\nu s^\mu),
\end{aligned}
\tag{1}
$$

where the nucleon spin, $s^\mu$, is an axial vector given in terms of the nucleon momentum $p^\mu = (E, \vec{p})$ and the direction of quantization, $\hat{n}$, as

$$s^0 = \vec{p}\cdot\hat{n}, \quad \vec{s} = m\hat{n} + \frac{\vec{p}\cdot\hat{n}}{E+m}\vec{p}. \tag{2}$$

The terms with $T_L^{ab}(q^2, \nu)$ and $T_2^{ab}(q^2, \nu)$ are the spin-independent, $VV+AA$ contribution; $T_3^{ab}(q^2, \nu)$ and $T_4^{ab}(q^2, \nu)$ mark the spin-dependent, $VV+AA$ contribution. The term with $T_5^{ab}(q^2, \nu)$ is the only spin-independent, $VA+AV$ part, while $T_6^{ab}(q^2, \nu)$, $T_7^{ab}(q^2, \nu)$, and $T_8^{ab}(q^2, \nu)$ give the spin-dependent, $VA+AV$ terms.

The final two terms in (1) are the $t$-channel poles, where we have defined $\Gamma_c$ and $\Gamma_c^A$ in terms of the matrix elements of the vector and axial-vector currents,

$$\langle p|V_c^\mu(0)|p\rangle = p^\mu \Gamma_c, \tag{3a}$$

$$\langle p|A_c^\mu(0)|p\rangle = s^\mu \Gamma_c^A. \tag{3b}$$

The cross section in the laboratory frame for the scattering of a neutrino with a nucleon whose spin is aligned parallel (+) or antiparallel (−) to the direction of motion of the incoming neutrino is written in terms of the imaginary parts, $W_i(q^2, \nu)$, of the $T_i(q^2, \nu)$ as

$$\begin{aligned}\frac{d\sigma^{\nu p(\pm)}}{dE'd\Omega} = \frac{G^2}{16\pi^3}\frac{E'}{mE}\big[&-q^2W_L(q^2,\nu) + (E^2+E'^2+\tfrac{1}{2}q^2)m^2W_2(q^2,\nu) \pm mq^2(E+E'\cos\theta)W_3(q^2,\nu)\\ &\pm m^2q^2(E+E')(E-E'\cos\theta)W_4(q^2,\nu) - \tfrac{1}{2}mq^2(E+E')W_5(q^2,\nu) \mp mq^2(E-E'\cos\theta)W_6(q^2,\nu)\\ &\mp m^3(E-E'\cos\theta)(E^2+E'^2+\tfrac{1}{2}q^2)W_7(q^2,\nu) \mp 2m^2EE'(1+\cos\theta)W_8(q^2,\nu)\big],\end{aligned} \tag{4}$$

where $E$ is the energy of the original neutrino, $E'$ is the energy of the final electron, $\theta$ is the scattering angle, $m$ is the mass of the nucleon, $G$ is the Fermi coupling constant, $q^2 = -2EE'(1-\cos\theta)$, and $\nu = m(E-E')$. The electron mass has been set equal to zero.[10] In writing (4) we have neglected the Cabibbo angle (since $\sin\theta_C \approx 0.05 \approx 0$) and so have only the $\Delta S = 0$ cross section. For the scattering of an antineutrino we replace $W_i(\nu, q^2)$ by $-W_i(-\nu, q^2)$ for $i = L, 2, 4, 8$ and by $+W_i(-\nu, q^2)$ for $i = 3, 5, 6, 7$.

It is easy to see, by writing (4) in terms of the variables $\omega = -q^2/2\nu$ and $x = \nu/mE$, that each term in (4) increases linearly with $E$ as $E$ goes to infinity. Such behavior is well known for the spin-independent terms.

To derive the relations we use the light-cone BJL theorem,[1,3]

$$T_{ab}^{\mu\nu}(p,q) \underset{q^-\to\infty}{\sim} \text{polynomial} - \frac{1}{q^-}\int dx^- d^2x_\perp e^{iq^+x^-} e^{-i\vec{q}_\perp\cdot\vec{x}_\perp}\langle p|[J_a^\mu(x), J_b^\nu(0)]|p\rangle, \tag{5}$$

with the commutators[1]

$$\begin{aligned}[V_a^+(x), V_b^+(0)]_{x^+=0} &= [A_a^+(x), A_b^+(0)]_{x^+=0}\\ &= if_{abc}V_c^+(0)\delta(x^-)\delta^2(x_\perp),\end{aligned} \tag{6}$$

$$[V_a^+(x), A_b^+(0)]_{x^+=0} = if_{abc}A_c^+(0)\delta(x^-)\delta^2(x_\perp), \tag{7}$$

$$\begin{aligned}[V_a^+(x), V_b^-(0)]_{x^+=0} &= [A_a^+(x), A_b^-(0)]_{x^+=0}\\ &= if_{abc}V_c^-(0)\delta(x^-)\delta^2(x_\perp)\\ &\quad - \tfrac{1}{2}i f_{abc}\{\partial_-[\epsilon(x^-)\delta^2(x_\perp)\mathcal{V}_c^-(x|0)] + \tfrac{1}{2}\partial_i[\epsilon(x^-)\delta^2(x_\perp)\mathcal{V}_c^i(x|0)] - \tfrac{1}{2}\partial_i\epsilon^{ij}[\epsilon(x^-)\delta^2(x_\perp)\bar{\mathcal{A}}_{jc}(x|0)]\}\\ &\quad - \tfrac{1}{2}i d_{abc}\{\partial_-[\epsilon(x^-)\delta^2(x_\perp)\bar{\mathcal{V}}_c^-(x|0)] + \tfrac{1}{2}\partial_i[\epsilon(x^-)\delta^2(x_\perp)\bar{\mathcal{V}}_c^i(x|0)] + \tfrac{1}{2}\partial_i\epsilon^{ij}[\epsilon(x^-)\delta^2(x_\perp)\mathcal{A}_{jc}(x|0)]\},\end{aligned} \tag{8}$$

$$\begin{aligned}[V_a^+(x), A_b^-(0)]_{x^+=0} &= [A_a^+(x), V_b^-(0)]_{x^+=0}\\ &= if_{abc}A_c^-(0)\delta(x^-)\delta^2(x_\perp)\\ &\quad - \tfrac{1}{2}i f_{abc}\{\partial_-[\epsilon(x^-)\delta^2(x_\perp)\mathcal{A}_c^-(x|0)] + \tfrac{1}{2}\partial_i[\epsilon(x^-)\delta^2(x_\perp)\mathcal{A}_c^i(x|0)] - \tfrac{1}{2}\epsilon^{ij}\partial_i[\epsilon(x^-)\delta^2(x_\perp)\bar{\mathcal{V}}_{cj}(x|0)]\}\\ &\quad - \tfrac{1}{2}i d_{abc}\{\partial_-[\epsilon(x^-)\delta^2(x_\perp)\bar{\mathcal{A}}_c^-(x|0)] + \tfrac{1}{2}\partial_i[\epsilon(x^-)\delta^2(x_\perp)\bar{\mathcal{A}}_c^i(x|0)] + \tfrac{1}{2}\epsilon^{ij}\partial_i[\epsilon(x^-)\delta^2(x_\perp)\mathcal{V}_{jc}(x|0)]\},\end{aligned} \tag{9}$$

where $\mathcal{V}_c^\mu(x|0)$, $\bar{\mathcal{V}}_c^\mu(x|0)$, $\mathcal{A}_c^\mu(x|0)$, and $\bar{\mathcal{A}}_c^\mu(x|0)$ are the bilocal operators defined in Ref. 3. The form factors are defined by

$$\langle p|\mathcal{V}_c^\mu(x|0)|p\rangle = p^\mu V_1^c(x^2, x\cdot p) + x^\mu V_2^c(x^2, x\cdot p), \tag{10}$$

$$\langle p|\mathcal{A}_c^\mu(x|0)|p\rangle = s^\mu A_1^c(x^2, x\cdot p) + p^\mu x\cdot s A_2^c(x^2, x\cdot p) + x^\mu x\cdot s A_3^c(x^2, x\cdot p) \tag{11}$$

with similar definitions for $\bar{\mathcal{V}}_c^\mu(x|0)$ in terms of $\bar{V}_i^c(x^2, x\cdot p)$ and $\bar{\mathcal{A}}_c^\mu(x|0)$ in terms of $\bar{A}_i^c(x^2, x\cdot p)$.

The scaling limits ($\nu\to\infty$, $q^2\to\infty$, $-q^2/2\nu\equiv\omega$ fixed) of the structure functions are

$$W_L^{ab}(q^2, \nu) \to \frac{-1}{2\omega}\left(F_L^{ab}(\omega) + \frac{1}{q^2}G_L^{ab}(\omega)\right), \tag{12a}$$

$$\nu W_i^{ab}(q^2, \nu) \to F_i^{ab}(\omega) + \frac{1}{q^2}G_i^{ab}(\omega), \quad i = 2, 3, 5, 6, 8 \tag{12b}$$

$$\nu^2 W_i^{ab}(q^2, \nu) \to F_i^{ab}(\omega) + \frac{1}{q^2}G_i^{ab}(\omega), \quad i = 4, 7. \tag{12c}$$

Using $F_i^{ab}(\omega) = F_i^{(ab)}(\omega) + iF_i^{[ab]}(\omega)$ we have our results:

$$F_L^{ab}(\omega) = 0, \tag{13}$$

$$F_2^{[ab]}(\omega) = \omega f_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} V_1^c(0, \alpha), \tag{14a}$$

$$F_5^{(ab)}(\omega) = 2d_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} V_1^c(0, \alpha), \tag{14b}$$

$$F_2^{(ab)}(\omega) = i\omega d_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \bar{V}_1^c(0, \alpha), \tag{15a}$$

$$F_5^{[ab]}(\omega) = 2if_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \bar{V}_1^c(0, \alpha), \tag{15b}$$

$$\omega F_6^{ab}(\omega) + F_8^{ab}(\omega) = 0, \tag{16}$$

$$F_3^{[ab]}(\omega) = \tfrac{1}{2} if_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \bar{A}_1^c(0, \alpha), \tag{17a}$$

$$F_8^{(ab)}(\omega) = i\omega d_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \bar{A}_1^c(0, \alpha), \tag{17b}$$

$$F_3^{(ab)}(\omega) = \tfrac{1}{2}d_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} A_1^c(0, \alpha), \tag{18a}$$

$$F_8^{[ab]}(\omega) = \omega f_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} A_1^c(0, \alpha), \tag{18b}$$

$$F_4^{[ab]}(\omega) = \tfrac{1}{2}i f_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \alpha\bar{A}_2^c(0, \alpha), \tag{19a}$$

$$F_7^{(ab)}(\omega) + F_8^{(ab)}(\omega) = i\omega d_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \alpha\bar{A}_2^c(0, \alpha), \tag{19b}$$

$$F_4^{(ab)}(\omega) = \tfrac{1}{2}d_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \alpha A_2^c(0, \alpha), \tag{20a}$$

$$F_7^{[ab]}(\omega) + F_8^{[ab]}(\omega) = \omega f_{abc}\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha} \alpha A_2^c(0, \alpha), \tag{20b}$$

and

$$G_L^{ab}(\omega) = -8i\omega^3\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha}\alpha[d_{abc}\bar{V}_2^c(0, \alpha) + f_{abc}V_2^c(0, \alpha)], \tag{21}$$

$$\omega G_6^{ab}(\omega) + G_8^{ab}(\omega) = -4i\omega^3\int_{-\infty}^{\infty} d\alpha\, e^{-i\omega\alpha}\alpha^2 \times[d_{abc}\bar{A}_3^c(0, \alpha) + f_{abc}A_3^c(0, \alpha)]. \tag{22}$$

Equation (13) follows from the fact that there is no $q$-number Schwinger term on the right-hand side of (6).[1] In the same way if there were a Schwinger term in (7) of the form $\partial_-\partial_-[P^{ab}(x|0) \times\epsilon(x^-)\delta^2(x_\perp)]$, where $P^{ab}(x|0)$ is a pseudoscalar whose matrix element can be written as $x\cdot s \times P^{ab}(x^2, x\cdot p)$, then the right-hand side of (16) would be the Fourier transform of $x\cdot pP^{ab}(0, x\cdot p)$.

Equations (13), (14), (15), (17a), (18a), (19a), (20a), and (21) have been derived before.[1,2,3] Equations (16), (17b), (18b), (19b), (20b), and (22) are new. These can be used to derive many sum rules. One that is well known[2,5-7] used (15) to relate $F_2(\omega)$ for electron-nucleon scattering with $F_5(\omega)$ for neutrino-nucleon scattering (neglecting the Cabibbo angle):

$$6[F_2^{ep}(\omega) - F_2^{en}(\omega)] = \omega[F_5^{\nu n}(\omega) - F_5^{\nu p}(\omega)]. \tag{23}$$

In the same way (18) and (20) give

$$6\omega[F_3^{ep}(\omega) - F_3^{en}(\omega)] = [F_8^{\nu n}(\omega) - F_8^{\nu p}(\omega)], \tag{24}$$

$$6\omega[F_4^{ep}(\omega) - F_4^{en}(\omega)] = [F_7^{\nu n}(\omega) + F_8^{\nu n}(\omega) - F_7^{\nu p}(\omega) - F_8^{\nu p}(\omega)]. \tag{25}$$

Equation (18a) implies[3]

$$d_{abc}\Gamma_c^A = \frac{2}{\pi}\int_0^1 d\omega\, F_3^{(ab)}(\omega), \tag{26}$$

where $\Gamma_c^A$ is the matrix element of the axial-vector current as defined in (3b). From (18b), (20b), and (16) we have

$$f_{abc}\Gamma_c^A = \frac{1}{\pi}\int_0^1 \frac{d\omega}{\omega} F_8^{[ab]}(\omega) \tag{27a}$$

$$= -\frac{1}{\pi}\int_0^1 d\omega\, F_6^{[ab]}(\omega) \tag{27b}$$

$$= -\frac{1}{\pi}\int_0^1 \frac{d\omega}{\omega} F_7^{[ab]}(\omega). \tag{27c}$$

The Dashen–Fubini–Gell-Mann sum rule in the deep-inelastic region follows from (14a)[11]:

$$\int_0^1 d\omega \frac{1}{\omega} F_2^{[ab]}(\omega) = \pi f_{abc} \Gamma^c , \qquad (28)$$

where $\Gamma^c$ is defined in (3a). Now we also get a sum rule from (14b)[12]:

$$\int_0^1 d\omega\, F_5^{(ab)}(\omega) = 2\pi d_{abc} \Gamma^c . \qquad (29)$$

In the same way the other sum rules derived in Ref. 3 for $F_2^{ab}(\omega)$, $F_3^{ab}(\omega)$, and $F_4^{ab}(\omega)$ now have analogs in terms of $F_5^{ab}(\omega)$, $F_6^{ab}(\omega)$, $F_7^{ab}(\omega)$, and $F_8^{ab}(\omega)$.

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[9]In this paper we are only interested in the deep-inelastic (scaling) region where, presumably, all masses can be neglected and all currents are conserved. The number of invariant amplitudes is determined by counting helicity amplitudes for spin-1, spin-$\frac{1}{2}$ scattering into spin 1, spin $\frac{1}{2}$. There are 12 independent helicity amplitudes, four of which remain in the forward direction; $T_{1,\frac{1}{2};1,\frac{1}{2}}$, $T_{0,\frac{1}{2};0,\frac{1}{2}}$, $T_{-1,\frac{1}{2};-1,\frac{1}{2}}$, and $T_{1,\frac{1}{2};0,-\frac{1}{2}}$. When we sum over the spin of the nucleon we have two amplitudes for the vector-vector case, $T_{1,\frac{1}{2};1,\frac{1}{2}} + T_{1,-\frac{1}{2};1,-\frac{1}{2}} = T_{1,\frac{1}{2};1,\frac{1}{2}} + T_{-1,\frac{1}{2};-1,\frac{1}{2}}$ and $T_{0,\frac{1}{2};0,\frac{1}{2}} + T_{0,-\frac{1}{2};0,-\frac{1}{2}} = 2T_{0,\frac{1}{2};0,\frac{1}{2}}$; but only one amplitude for the vector-axial-vector case because $T_{0,\frac{1}{2};0,\frac{1}{2}} + T_{0,-\frac{1}{2};0,-\frac{1}{2}} = 0$.

[10]The fact that we have treated the axial-vector current as being conserved and written all the terms in (1) in a gauge-invariant manner does not change the cross section, Eq. (4), as long as the electron mass is neglected.

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# High-Energy Limit of Lowest-Order Weak Amplitudes with Electromagnetic Radiative Corrections in the Ladder Approximation

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The asymptotic lowest-order weak amplitude for the process $\nu_\mu + \bar{\nu}_e \to \mu^- + e^+$ is found to be modified by a factor $\exp[-(\alpha/2\pi)\ln t]$ on account of radiative corrections in the ladder approximation. This is in conformity with the result obtained earlier by Li in the perturbation theory.

## I. INTRODUCTION

It is well known that cross sections for the purely leptonic weak scattering processes, such as $\nu_\mu + e^- \to \mu^- + \nu_e$, $\nu_\mu + \bar{\nu}_e \to \mu^- + e^+$, etc., increase quadratically with the center-of-mass energy of the system, in the lowest order in the weak interaction in the current×current theory.[1] Recently Li[2] has investigated the effect of electromagnetic radiative corrections to the high-energy behavior of lowest-order weak amplitudes in perturbation theory. He finds that the damping thus obtained, though calculated to all orders in $\alpha$, does not produce a significant change in the behavior of these amplitudes at the available experimental energy resolutions. His method essentially consists in replacing all the fermion propagators in the matrix element for the $n$th-order Feynman diagram by