

## Contributions of vector meson photoproduction to the Gerasimov-Drell-Hearn sum rule

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An improved version of a recently developed quark model approach to vector meson photoproduction is applied to the investigation of contributions of vector meson photoproduction to the Gerasimov-Drell-Hearn (GDH) sum rule. We find that the sum rule converges at a few GeV's. Contributions to the proton channel are found to be small while to the neutron are relatively large.

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The GDH sum rule [1] connects the nucleon resonance phenomena to the nucleon's magnetic moments, which are static properties of the groundstate nucleons,

$$I_{GDH} = \int_{\nu_0}^{\infty} [\sigma_{1/2}(\nu) - \sigma_{3/2}(\nu)] \frac{d\nu}{\nu} = -\frac{2\pi^2 \alpha_e \kappa^2}{m_N^2}, \quad (1)$$

where  $m_N$  is the nucleon mass;  $\alpha_e$  is the fine structure constant;  $\kappa$  is the nucleon anomalous magnetic moment;  $\sigma_{1/2}$  and  $\sigma_{3/2}$  respectively denote photoabsorption cross sections for the nucleon and photon helicities antiparallel and parallel with each other;  $\nu_0$  denotes the threshold energy for single pion production in lab system. Experimental and theoretical studies of exclusive reactions provide another means of testing this sum rule, which can lead to better understanding of the internal degrees of freedom of nucleons. The availability of high intensive electron and photon beam facilities gives access to precise measurements of meson photoproduction. Recently, the GDH- and A2-Collaboration reported their results at the photon energies from 200 to 800 MeV [2]. Other experimental projects at JLab, MAMI, GRAAL, ELSA, and SPring8 will make it possible to test this sum rule independently and extend it to higher energies. In theory, extensive investigations [3–10] from the low-energy limit to around 2 GeV photon energies have been carried out for the light meson production channels. The most recent study by the Mainz group [10] including  $\pi N$ ,  $\eta N$ ,  $\pi\pi B$  (Born terms), and  $\pi\pi D_{13}(1535)$  showed that contributions from those channels accounted for the sum rule for the proton up to 97%, while large discrepancies were found for the neutron. Interestingly, these theoretical evaluations including the above channels led to results smaller than the absolute values for both proton and neutron, which suggests that contributions from higher energy production channels might add instead of cancel certain terms in the exclusive sum  $I_{GDH}$  in Eq. (1). In Ref. [11] contributions of the kaon photoproduction to the GDH sum rule was estimated up to 2 GeV. It was found that the kaon contribution increased the calculated value of GDH sum rule for the neutron, while might decrease that for the proton. To satisfy the GDH sum rule by exclusive study, one might have to go up to a higher energy region. On the one

hand, there might be still significant contributions from those light meson productions. On the other hand, open channels above the kaon channels could start to play a role. This is one of our motivations for investigating the contributions from the vector meson production channels, for which the thresholds are just above the kaon production. Kinematically, the light vector meson ( $\omega$  and  $\rho$ ) production is the next contribution that should be included in the exclusive calculations. Certainly, the energy suppression in the integral could have made the vector meson contribution trivial. But such an effort is by no means trivial due to the diffractive features in vector meson photoproduction. It would be a challenge for any model to show how the spin-dependent feature changes to spin-independent at high energies as required by the sum rule. A reasonable estimation of the contributions of vector meson photoproduction to the GDH sum rule will also provide a test for the model.

It should be noted that in Ref. [12] contributions above the resonance regions to the GDH sum rule were estimated using a Regge parametrization, which accounted for a large fraction of discrepancies between the sum rule and contributions from the single pion photoproduction. This result could suggest that an explicit study of higher threshold processes should be necessary.

We employ a recently developed quark model approach to vector meson photoproduction [13–17] for this purpose. It allows not only the study of resonance excitations in the  $\omega$ ,  $\rho$ , and  $\phi$  meson photoproduction, but also the inclusion of the diffractive contributions with a mixed Pomeron exchange model. Thus, the model can be extended to photon energies of a few GeV's. Unlike the study of various polarization observables, the most essential question arising from the sum rule study is the role played by the diffractive processes. In principle, the sum rule requires integrating the cross section difference up to infinite energy, while in vector meson production the diffractive cross sections tend to a constant value at high energies. This leaves the question of whether the spin-dependent terms can be averaged out efficiently at high energies. An analytical illustration is not available yet. However, qualitatively, based on this specific model, we can see later that this approach allows us to make an energy cut at a few GeV's, where the spin-dependent terms have been sufficiently suppressed and averaged, and the dominant term in the diffractive process has spin-independent feature. As a

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consequence, the integral of the cross section difference will converge and be reliably estimated at a few GeV's.

The transition amplitude for  $\gamma N \rightarrow VN'$  can be expressed by 12 independent helicity amplitudes:

$$\begin{aligned} H_{1\lambda_v} &= \langle \lambda_v, \lambda_f = +1/2 | T | \lambda_\gamma = +1, \lambda_i = -1/2 \rangle, \\ H_{2\lambda_v} &= \langle \lambda_v, \lambda_f = +1/2 | T | \lambda_\gamma = +1, \lambda_i = +1/2 \rangle, \\ H_{3\lambda_v} &= \langle \lambda_v, \lambda_f = -1/2 | T | \lambda_\gamma = +1, \lambda_i = -1/2 \rangle, \\ H_{4\lambda_v} &= \langle \lambda_v, \lambda_f = -1/2 | T | \lambda_\gamma = +1, \lambda_i = +1/2 \rangle, \end{aligned} \quad (2)$$

where  $T$  is the dynamic operator for this transition;  $\lambda_\gamma$  and  $\lambda_v$  ( $=0, \pm 1$ ) are helicities for the incident photon and outgoing vector meson, respectively, while  $\lambda_i$  and  $\lambda_f$  are helicities for the initial and final state nucleons. For simplicity, we have fixed the photon polarization vector as  $\epsilon_\gamma = -(1, i, 0)/\sqrt{2}$ . Parity conservation will give amplitudes for  $\epsilon_\gamma = (1, -i, 0)/\sqrt{2}$ , which are not independent from the former. For exclusive photoproduction, apart from the kinematic factors, one can see that

$$\begin{aligned} \sigma_{1/2}^v &\propto H_{1\lambda_v}^2 + H_{3\lambda_v}^2, \\ \sigma_{3/2}^v &\propto H_{2\lambda_v}^2 + H_{4\lambda_v}^2. \end{aligned} \quad (3)$$

The contribution from vector meson photoproduction is expressed as

$$I_{GDH}^v = \int_{\nu_v}^{\infty} \Delta\sigma(\nu) \frac{d\nu}{\nu} = -\frac{2\pi^2 \alpha_e \kappa_v^2}{m_N^2}, \quad (4)$$

where  $\Delta\sigma(\nu) \equiv \sigma_{1/2}^v - \sigma_{3/2}^v$  is the photoabsorption cross section difference and the total cross section for the unpolarized vector meson production is  $\sigma_{tot} = (\sigma_{1/2}^v + \sigma_{3/2}^v)/2$ ;  $\kappa_v$  denotes the contribution from vector meson production to the nucleon anomalous magnetic moment;  $\nu_v$  is the threshold energy for the vector meson in the lab system (i.e., the nucleon rest frame). For the  $\omega$  meson,  $\nu_\omega = 1.108$  GeV, while for  $\rho$ ,  $\nu_\rho = 1.086$  GeV. Using the recent quark model developed by Zhao *et al.* [13–17], we shall explicitly calculate the photoabsorption cross sections  $\sigma_{1/2}^v$  and  $\sigma_{3/2}^v$ .

In this model, there are three processes that contribute to the transition amplitudes for the neutral vector meson photoproduction: (i) the  $s$ - and  $u$ -channel resonance excitations and the nucleon pole terms; (ii) the  $t$ -channel light meson exchange (i.e., pion exchange in  $\omega$  production and  $\sigma$  meson exchange in  $\rho$  production); and (iii) the  $t$ -channel Pomeron exchange in the neutral vector meson production (i.e.,  $\gamma N \rightarrow \omega N$  and  $\gamma N \rightarrow \rho^0 N$ ).

In the helicity frame, the  $s$ -channel resonance excitation amplitude can be expressed as the product of the resonance electromagnetic excitation helicity amplitude  $A_{\Lambda_i}^\chi$  ( $\Lambda_i = 1/2, 3/2$ ) and its vector meson decay amplitude  $A_{\Lambda_f}^v$  ( $\Lambda_f = 1/2, 3/2$ ) for the transverse polarization and  $S_{\Lambda_f}^v$  ( $\Lambda_f = 1/2$ ) for the longitudinal polarization. For a resonance of the

$SU(6) \otimes O(3)$  quark model with spin  $J$ , its 12 independent transition amplitudes can be written

$$\begin{aligned} H_{11}^J &= d_{1/2,3/2}^J(\theta) A_{1/2}^v A_{3/2}^\chi, \\ H_{10}^J &= d_{-1/2,3/2}^J(\theta) S_{-1/2}^v A_{3/2}^\chi \\ &= (-1)^{J_v} d_{1/2,3/2}^J(\pi + \theta) S_{1/2}^v A_{3/2}^\chi, \\ H_{1-1}^J &= d_{-3/2,3/2}^J(\theta) A_{-3/2}^v A_{3/2}^\chi \\ &= (-1)^{J_v+1} d_{3/2,3/2}^J(\pi + \theta) A_{3/2}^v A_{3/2}^\chi, \\ H_{21}^J &= d_{1/2,1/2}^J(\theta) A_{1/2}^v A_{1/2}^\chi, \\ H_{20}^J &= d_{-1/2,1/2}^J(\theta) S_{-1/2}^v A_{1/2}^\chi \\ &= (-1)^{J_v} d_{1/2,1/2}^J(\pi + \theta) S_{1/2}^v A_{1/2}^\chi, \\ H_{2-1}^J &= d_{-3/2,1/2}^J(\theta) A_{-3/2}^v A_{1/2}^\chi \\ &= (-1)^{J_v} d_{3/2,1/2}^J(\pi + \theta) A_{3/2}^v A_{1/2}^\chi, \\ H_{31}^J &= d_{3/2,3/2}^J(\theta) A_{3/2}^v A_{3/2}^\chi, \\ H_{30}^J &= d_{1/2,3/2}^J(\theta) S_{1/2}^v A_{3/2}^\chi, \\ H_{3-1}^J &= d_{-1/2,3/2}^J(\theta) A_{-1/2}^v A_{3/2}^\chi \\ &= (-1)^{J_v} d_{1/2,3/2}^J(\pi + \theta) A_{1/2}^v A_{3/2}^\chi, \\ H_{41}^J &= d_{3/2,1/2}^J(\theta) A_{3/2}^v A_{1/2}^\chi \\ &= -d_{1/2,3/2}^J(\theta) A_{3/2}^v A_{1/2}^\chi, \\ H_{40}^J &= d_{1/2,1/2}^J(\theta) S_{1/2}^v A_{1/2}^\chi, \\ H_{4-1}^J &= d_{-1/2,1/2}^J(\theta) A_{-1/2}^v A_{1/2}^\chi \\ &= (-1)^{J_v} d_{1/2,1/2}^J(\pi + \theta) A_{-1/2}^v A_{1/2}^\chi, \end{aligned} \quad (5)$$

where in some of the above equations, parity conservation allows us to relate  $A_{-\Lambda_f}^v$  with  $A_{\Lambda_f}^v$  for each  $SU(6) \otimes O(3)$  state with spin  $J$  which then decays into a vector meson and a nucleon with relative angular momentum  $J_v$ ,

$$A_{-\Lambda_f}^v(J) = (-1)^{1/2-J-J_v} A_{\Lambda_f}^v(J), \quad (6)$$

where the factor  $1/2$  denotes the spin of the final state nucleon. The parity of such a state is  $(-1)^N$ , where  $N$  is the main quantum number of the harmonic oscillator shell. Meanwhile, the final state system has a parity  $(+1)(-1)^{J_v}$  which is determined by the parity of the nucleon, vector meson, and their relative angular momentum. Equivalence of these two parities gives  $(-1)^N = (-1)^{J_v+1}$ , and thus

$$A_{-\Lambda_f}^v(J) = (-1)^{(1/2-J)-(N-1)} A_{\Lambda_f}^v(J). \quad (7)$$

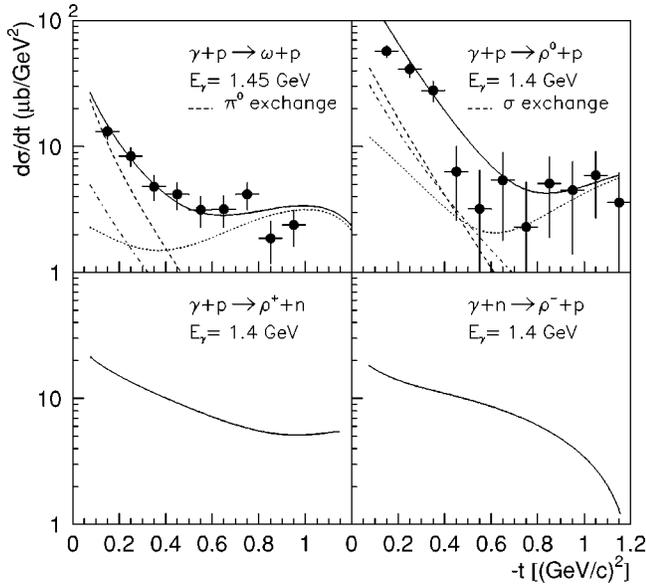


FIG. 1. Differential cross sections for four isospin channels. The solid curves denote full-model calculations, while the dotted and dot-dashed curves denote the exclusive  $s$ - and  $u$ -channel processes and Pomeron exchanges, respectively. The dashed curve denotes pion exchange for  $\omega$  production and  $\sigma$  exchange for  $\rho^0$  production, respectively. Data come from Ref. [21].

So far, the dynamics of the reaction has not been explicitly involved. Note that the change of the rotation functions has been taken into account in Eq. (5). The symmetric feature for each state of  $SU(6) \otimes O(3)$  quark model can be seen clearly. The nodal structure is determined by the interfering amplitudes, therefore can be studied in an explicit model.

Equations (5) and (7) add some interesting relations to the 12 independent transition amplitudes. For example, apart from the rotation function and phase factors, we can see that  $\mathcal{H}_{1\lambda_v}^J = \mathcal{H}_{3(-\lambda_v)}^J$ , and  $\mathcal{H}_{2\lambda_v}^J = \mathcal{H}_{4(-\lambda_v)}^J$ , where  $\mathcal{H}$  denotes the product of  $A^v(S^v)$  and  $A^J$  in Eq. (5). Thus, the corresponding amplitudes have the same dynamical parts. Note that the phase factor and the rotation function carry information about the resonance partial waves and symmetries, and the dynamical part carries information about the resonance structures, which are generally spin-dependent. Therefore, when we add the amplitudes for resonance excitations together and calculate the cross section differences, interferences among resonances of different quark model representations will arise. At low energies, such interferences produce nonzero contributions to the sum rule. We show in Fig. 1 the spin-dependent phenomena near threshold with the study of cross sections for four reactions. We shall see below that at the energy very close to threshold, the strong  $D_{33}(1700)$  will lead to an overestimation of the total cross sections for  $\gamma p \rightarrow \rho^0 p$ . However, instead of presenting details for the study of resonance excitations, our special attentions will be paid to their high energy behavior. We expect that at high energies, the spin-dependent terms will die out, while spin-independent terms will dominate and result in convergence of the cross section difference.

Certainly, the convergence of the cross section difference

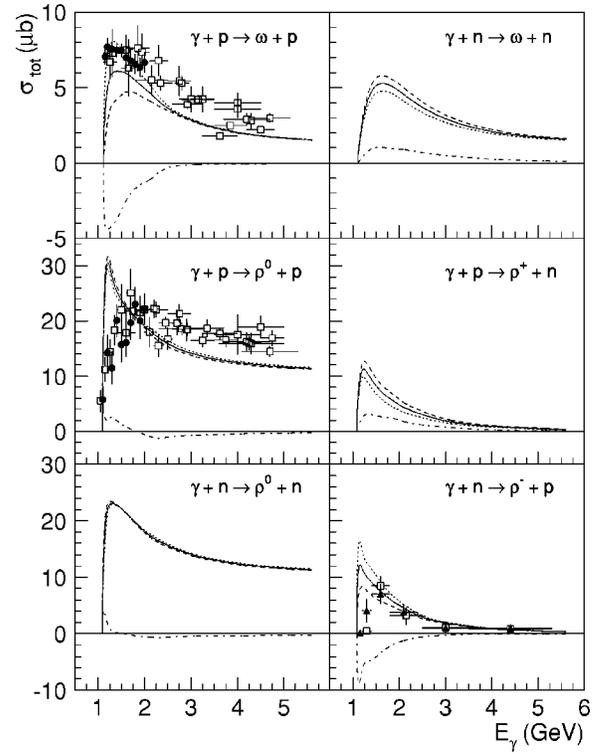


FIG. 2. Total cross sections and cross section differences for the  $\omega$  and  $\rho$  meson photoproduction. The solid, dashed, dotted, and dot-dashed curves denote the calculations for  $\sigma_{tot}[(\sigma_{1/2} + \sigma_{3/2})/2]$ ,  $\sigma_{1/2}$ ,  $\sigma_{3/2}$ , and  $\Delta\sigma(=\sigma_{1/2} - \sigma_{3/2})$ , respectively. Data come from Refs. [21–23].

is not obvious in such an approach. However, several basic aspects can be outlined. Firstly, the quark model wave functions guarantee the theory to be unitary when all the baryon resonances are included. The spatial integrals provide form factors which are proportional to  $e^{-(\mathbf{k}^2 + \mathbf{q}^2)/6\alpha^2}$  in the transition amplitudes. This factor results in the disappearance of resonance contributions at high energies, which corresponds to high  $|\mathbf{k}|$  and  $|\mathbf{q}|$ . For higher excited states, as shown in Ref. [14], in the process of photon and meson coupling to the same quark, the form factor for the harmonic oscillator shell  $N$  is

$$F_{Nl}(\mathbf{k}, \mathbf{q}) = \frac{1}{(N-l)!} \left( \frac{\mathbf{k} \cdot \mathbf{q}}{3\alpha^2} \right)^{N-l} e^{-(\mathbf{k}^2 + \mathbf{q}^2)/6\alpha^2}, \quad (8)$$

while in the process of photon and meson coupling to different quarks,

$$F'_{Nl}(\mathbf{k}, \mathbf{q}) = \frac{1}{(N-l)!} \left( -\frac{1}{2} \right)^N \left( \frac{\mathbf{k} \cdot \mathbf{q}}{3\alpha^2} \right)^{N-l} e^{-(\mathbf{k}^2 + \mathbf{q}^2)/6\alpha^2}, \quad (9)$$

where  $N \geq l$  when summed over all permitted  $N$  and  $l$ . Apparently, an additional factor  $(-1/2)^N$  suppresses the process that the photon and meson couple to different quarks at high energies. Meanwhile, for a given  $N$ , the dominant contribution comes from terms with  $l=0$ .

The dying-out trend governed by the form factor provides a possible way to bypass difficulty arising from resonance excitations. A simple analytical argument can be made by assuming that those degenerate resonances are on their mass shells and have the same width. Then, the transition amplitudes can be factorized out as

$$H_{a\lambda_\nu} = S_{a\lambda_\nu}(\mathbf{k}, \mathbf{q}) \sum_{N=0}^{\infty} F_{N0}(\mathbf{k}, \mathbf{q}) \\ = S_{a\lambda_\nu}(\mathbf{k}, \mathbf{q}) e^{-(\mathbf{k}-\mathbf{q})^2/6a^2}, \quad (10)$$

where  $S_{a\lambda_\nu}(\mathbf{k}, \mathbf{q})$  is a sum of all the spin-dependent and independent parts. Note that, we neglect the process of photon and meson coupling to different quarks, and take the leading contribution of  $l=0$ . The second line of Eq. (10) could be regarded as a good approximation for the resonance excitations at  $\nu \rightarrow \infty$ . The behavior of the exponent, despite its angular-dependence, suppresses the amplitudes with increasing energies, especially at backward direction. This feature is very helpful since those spin-dependent terms generally have large effects at large angles. Remember that  $S_{a\lambda_\nu}(\mathbf{k}, \mathbf{q})$  contains both spin-dependent and independent terms, a simple survey over these amplitudes indeed suggests that the spin-dependent terms are strongly suppressed at high energies. It can be seen quite explicitly that the spin-independent term  $(\boldsymbol{\epsilon}_\nu \times \mathbf{q}) \cdot (\boldsymbol{\epsilon}_\gamma \times \mathbf{k}) = \boldsymbol{\epsilon}_\nu \cdot \boldsymbol{\epsilon}_\gamma \mathbf{q} \cdot \mathbf{k} - \boldsymbol{\epsilon}_\nu \cdot \mathbf{k} \boldsymbol{\epsilon}_\gamma \cdot \mathbf{q}$  plays a dominant role, where  $\boldsymbol{\epsilon}_\nu$  is the transverse polarization vector for the vector meson. This feature is the first aspect that guarantees the convergence of the integral at a few GeV's.

The spin-dependence of the Pomeron exchange in the cross section difference calculations is also not obvious due to interferences among different spin operators. Fortunately, however, some typical features of the Pomeron exchange model [18–20] can help us understand its behavior at high energies. Since the Pomeron, which accounts for the diffractive process, is rather like a charge conjugation  $C = +1$  vector meson, this feature means that its dominant contribution is located at small  $|t|$ . The longitudinal amplitude becomes negligible at high energies. As discussed in detail in Ref. [15], the dominant term is proportional to  $\boldsymbol{\epsilon}_\gamma \cdot \boldsymbol{\epsilon}_\nu$  in the c.m. system, which is spin-independent. When  $\nu \rightarrow \infty$ , we can immediately see that its contribution to the cross section difference  $\Delta\sigma(\nu)$  is zero, i.e.,  $H_{1\lambda_\nu}^2 + H_{3\lambda_\nu}^2 = H_{2\lambda_\nu}^2 + H_{4\lambda_\nu}^2$ . In other words, the Pomeron exchange contribution to the GDH sum rule becomes negligible at high energies, although it is dominant over all other processes.

Explicitly, the equivalence,  $H_{1\lambda_\nu}^2 + H_{3\lambda_\nu}^2 = H_{2\lambda_\nu}^2 + H_{4\lambda_\nu}^2$ , can be satisfied in the exclusive  $\pi^0$  (for  $\omega$  and  $\phi$ ) and  $\sigma$  exchange (for  $\rho^0$ ) due to the feature that no spin carried by the exchanged pion and  $\sigma$  meson. Undoubtedly, the interferences from the resonance excitations at low energies will violate the equivalence. However, beyond the resonance region, all these contributing processes will be dominated by the spin-independent terms, which will result in vanishing of  $\Delta\sigma(\nu)$ . Numerically, we find that when  $\nu \approx 6$  GeV,  $\Delta\sigma(\nu)$  becomes negligible. Above the energy, the effects from spin-

TABLE I. Contributions of vector meson photoproduction to the GDH sum rule in comparison with other exclusive channels. Contributions from  $\pi N, \eta N, \pi\pi B$  (Born terms), and  $\pi\pi D$  [ $D_{13}(1520)$  resonance] are from Mainz group study [10]; kaon channel contributions are from Ref. [11]; experimental data are from Ref. [2].

Proton	$I_{GDH}(\mu b)$	Neutron	$I_{GDH}(\mu b)$
$\gamma p \rightarrow \pi^0 p$	-150	$\gamma n \rightarrow \pi^0 n$	-154
$\gamma p \rightarrow \pi^+ n$	-21	$\gamma n \rightarrow \pi^- p$	+30
$\gamma p \rightarrow \eta p$	+15	$\gamma n \rightarrow \eta n$	+10
$\gamma p \rightarrow \pi\pi B$	-30	$\gamma n \rightarrow \pi\pi B$	-35
$\gamma p \rightarrow \pi\pi D$	-15	$\gamma n \rightarrow \pi\pi D$	-15
$\gamma p \rightarrow K^+ \Lambda$	+1.66	$\gamma n \rightarrow K^0 \Lambda$	-4.78
$\gamma p \rightarrow K^+ \Sigma^0$	+1.53	$\gamma n \rightarrow K^+ \Sigma^-$	+1.59
$\gamma p \rightarrow K^0 \Sigma^+$	+0.83	$\gamma n \rightarrow K^0 \Sigma^0$	+1.21
$\gamma p \rightarrow \omega p$	-2.01	$\gamma n \rightarrow \omega n$	+0.93
$\gamma p \rightarrow \rho^0 p$	+0.05	$\gamma n \rightarrow \rho^0 n$	-0.05
$\gamma p \rightarrow \rho^+ n$	+2.22	$\gamma n \rightarrow \rho^- p$	-2.93
sum of above	-196.72	sum of above	-168.03
GDH	-205	GDH	-233
exp. results	-210 <sup>a</sup>	exp. results	not available

<sup>a</sup>In Ref. [2],  $I_{GDH} = -226 \pm 5(\text{stat}) \pm 12(\text{syst}) \mu b$  was reported in the energy range  $200 < \nu < 800$  MeV, while taking into account the missing contributions from  $\nu < 200$  MeV, a deduced value  $-210 \mu b$  was estimated.

dependent terms will only make a change to the fourth decimal place in the integrals. This gives confidence of cutting off the photon energies at a few GeV's for our model.

In Fig. 2, we report the calculations of the cross section differences of six isospin channels. The parameters are extracted from the  $\omega$  production channel. We do not attempt to fit data for the  $\rho$  productions. This requires more subtle considerations. In fact, the prediction gives an overall agreement with the data, which suggests that isospin conservation has been roughly kept for the  $\omega$  and  $\rho$  mesons. A sensible feature arising between the  $\omega$  and  $\rho$  production is that isospin 3/2 resonances will contribute in the  $\rho$  meson production but be eliminated in the  $\omega$  production, as required by isospin conservation. Meanwhile, the quark model symmetry eliminates those states of quark model representation [70, 48] from contribution in the proton target reactions [24]. This feature results in an interchanging of the relative positions between the dashed ( $\sigma_{1/2}$ ) and dotted curves ( $\sigma_{3/2}$ ) in  $\gamma p \rightarrow \omega p$  and  $\gamma n \rightarrow \omega n$ . Similar phenomena can be also seen in  $\gamma p \rightarrow \rho^+ n$  and  $\gamma n \rightarrow \rho^- p$ . It is still worth noting that the  $D_{33}(1700)$  is found to play an important role in the  $\rho$  meson production near threshold. Details of this will be reported and discussed elsewhere.

The contribution of the  $\omega$  and  $\rho$  channels to the GDH sum rule are listed in Table I to compare with other channel contributions predicted by other studies. The neutral  $\rho$  production has a relatively smaller contribution to the proton and neutron sum rule, while contributions from the  $\omega$  and  $\rho^\pm$  channels are more than one order of magnitude larger. How-

ever, when we sum over those reaction channels for the proton and neutron, we find that the overall magnitude of the contribution to the proton ( $I_{GDH}^v = +0.26\mu\text{b}$ ) is much smaller than to the neutron ( $I_{GDH}^v = -2.05\mu\text{b}$ ). In particular, the sign of the sum ( $I_{GDH}^v$ ) suggests that the contribution to the proton sum rule will cancel a small number of previous results [10], while the contribution to the neutron will add a relatively larger one. This trend is fairly consistent with the exclusive studies of Ref. [10].

In conclusion, we have evaluated the contributions of vector meson photoproduction to the GDH sum rule using a quark model with an effective Lagrangian. Although more detailed study at resonance region is needed, we show that the cross section differences converge with the increasing energy, thus, leads to a reasonable energy cut at 6 GeV. The

main contributions of vector meson photoproduction are found from the resonance region. Although the total values are small, their corrections to the sum rule are shown to be in the right direction. This study provides a test for our approach for vector meson photoproduction though more accurate data from experiments are needed to constrain the model. Nevertheless, it provides some insights into the relevance of the exclusive vector meson photoproduction to the GDH sum rule.

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