PHYSICAL REVIEW C **72**, 031001(R) (2005)

ω production in pp collisions

G. Ramachandran,[1] M. S. Vidya,[2] P. N. Deepak,[3,4] J. Balasubramanyam,[5] and Venkataraya[6]
[1]*Indian Institute of Astrophysics, Koramangala, Bangalore 560034, India*
[2]*School of Physical Sciences, Jawaharlal Nehru University, New Delhi 110067, India*
[3]*Institut für Kernphysik, Forschungszentrum Jülich, D-52425 Jülich, Germany*
[4]*Physics Group, Birla Institute of Technology and Science, Pilani-Goa Campus, Zuarinagar-403726, India*
[5]*K. S. Institute of Technology, Bangalore 560062, India*
[6]*Vijaya College, Bangalore 560011, India*
(Received 3 January 2005; published 30 September 2005)

A model-independent irreducible tensor formalism that was developed earlier to analyze measurements of $\vec{p}\vec{p} \rightarrow pp\pi^{\circ}$ is extended to present a theoretical discussion of $\vec{p}\vec{p} \rightarrow pp\,\omega$ and of ω polarization in $pp \rightarrow pp\,\vec{\omega}$ and in $p\vec{p} \rightarrow pp\,\vec{\omega}$. The recent measurement of an unpolarized differential cross section for $pp \rightarrow pp\,\omega$ is analyzed by use of this theoretical formalism.

DOI: 10.1103/PhysRevC.72.031001 PACS number(s): 24.70.+s, 13.75.Cs, 13.88.+e, 25.40.Ve

Experimental study of meson production in *NN* collisions has attracted considerable interest during the past decade and a half. The early measurements of total cross section [1] for pion production were found surprisingly to be more than a factor of 5 than the theoretical predictions [2]. At c.m. energies close to threshold, the relative kinetic energies between the particles in the final state are small, and therefore an analysis involves only a few partial waves. On the other hand, a large momentum transfer is involved when an additional particle is produced in the final state, thus making the reaction sensitive to the features of the *NN* interaction at short distances where the nucleons start to overlap. When a heavier meson like ω is produced, the overlapping region corresponds [3] to a distance of about 0.2 fm. It is also known that the short-range part of the *NN* interaction is dominated by the ω exchange [4]. Consequently a variety of theoretical models have been proposed [5] not only to bridge the gap between theory and experiment, but also to test results of QCD-based discussions of the *NN* interaction. According to the Okubo-Zweig-Iizuka (OZI) rule [6], ϕ production relative to ω production is suppressed in the absence of strange quarks in the initial state. This ratio R has been measured [7] in view of the dramatic violations [8] observed in $\bar{p}p$ collisions and compared with the theoretical estimate [9] of 4.2×10^{-3} after correcting for the available phase space. We may refer to [10] for modifications of the rule. Apart from looking for the strange quark content of the nucleon in the initial state, attention has also been focused on resonance contributions [11–13] to vector-meson production in *NN* collisions. The constituent quark models [14] predict highly excited N^* states that have not been seen in πN scattering. This "missing resonance problem" [15] has also catalyzed the experimental study of ω meson production in the hope that the missing resonances may couple more strongly or even exclusively to the ωN channel in comparison with the πN channel, although ωN decay modes of resonances have not been observed [16]. Also the cross sections of vector-meson production enter as inputs into transport models for dilepton emission in heavy-ion collisions that may in turn be used to study the off-shell ω production and medium modifications of the widths and masses of the resonances [13].

Meson production in *NN* collisions involves also spin state transitions of the *NN* system, which do not occur in elastic *NN* scattering. In $pp \rightarrow pp\pi^0$, for example, the transition of the pp system at threshold is from an initial-spin triplet to a final-spin singlet state ($^3P_0 \rightarrow {}^1S_0$). Rapid advances in experimental technology have led today to high-precision measurements of spin observables [17] at several energies up to 400 MeV, employing beams of polarized protons on polarized proton targets. A conclusive theoretical interpretation of all these data has remained elusive, although the model calculations appear to do better in the case of charged pion production as compared with the neutral pion production, and the agreement even there seems to deteriorate increasingly at higher energies. It has been pointed out, both by Moskal *et al.* [5] and Hanhart [5], that the extensive experimental information available comes with a drawback that "apart from rare cases, it is difficult to extract a particular piece of information from the data."

A model-independent irreducible tensor formalism [18] that has been developed to analyze measurements on $\vec{p}\vec{p} \rightarrow pp\pi^0$ at the complete kinematical double-differential level was recently [19] made use of to estimate empirically the initial singlet and triplet state contributions to the differential cross section by use of the experimental results of Meyer *et al.* [17]. The preceding theoretical formalism leads, on integration, to the relation derived earlier by Bilenky and Ryndin [20] for the total cross sections. It was also shown [21] how the irreducible tensor formalism could be utilized to effect spin filtering, in general, for any scattering or reaction process employing polarized beams of particles with arbitrary spin s_b on polarized targets with arbitrary spin s_t. The production of a heavy meson like ω at and near threshold in $\vec{p}\vec{p}$ collisions allows us to study additional spin-dependent features of *NN* interactions at much shorter distances. Unlike the pion, which is spinless, the ω has spin 1, which permits us to make observations with regard to its spin state, apart from measuring the angular distributions in polarized beam and polarized target experiments. Experimental data on total [22] and differential [23] cross sections for $pp \rightarrow pp\omega$ have already been published and proposals are underway [5,24] to study heavy meson production in

NN collisions by use of polarized beams and targets at COoler SYnchrotron (COSY).

The purpose of the present paper is to extend the earlier work [18,19] on the model-independent approach based on irreducible tensor techniques to study the spin state of the meson in $pp \to pp\vec{\omega}$ and $p\vec{p} \to pp\vec{\omega}$ as well as the double-differential cross section in the proposed polarized beam and polarized target experiments.

Let $\boldsymbol{p}_i$ denote the initial c.m. momentum, $\boldsymbol{q}$ the momentum of the meson produced with spin-parity s^π, and $\boldsymbol{p}_f$ the relative momentum $(1/2)(\boldsymbol{p}_1 - \boldsymbol{p}_2)$, between the two nucleons with c.m. momenta $\boldsymbol{p}_1$ and $\boldsymbol{p}_2$ in the final state. The double-differential cross section for meson production in c.m. may be written as

$$d^2\sigma = \frac{2\pi D}{v}\mathrm{Tr}(\boldsymbol{T}\rho^i\boldsymbol{T}^\dagger), \tag{1}$$

where D denotes the final three-particle density of states, $\boldsymbol{T}$ denotes the on-energy-shell transition matrix and $\boldsymbol{T}^\dagger$ its Hermitian conjugate, $v = 4|\boldsymbol{p}_i|/E$ at c.m energy E, and ρ^i denotes the initial spin density matrix,

$$\rho^i = \tfrac{1}{4}(1 + \boldsymbol{\sigma}_1 \cdot \boldsymbol{P})(1 + \boldsymbol{\sigma}_2 \cdot \boldsymbol{Q}), \tag{2}$$

if $\boldsymbol{P}$ and $\boldsymbol{Q}$ denote, respectively, the beam and target polarizations. The notation $\sigma(\xi, \boldsymbol{P}, \boldsymbol{Q})$ is used in [17] to denote Eq. (1). If s_i and s_f denote the initial- and final-spin states of the *NN* system, the initial and final channel spins for the reaction are s_i and S, respectively, where S can assume values $S = |s_f - s|, \ldots, (s_f + s)$. Making use of the irreducible tensor operator techniques introduced in [25], we may express $\boldsymbol{T}$ in the operator form:

$$\boldsymbol{T} = \sum_\alpha \sum_{\lambda=|s_f-s_i|}^{(s_f+s_i)} \sum_{\Lambda=|S-s_i|}^{(S+s_i)} \times ((S^s(s,0) \otimes S^\lambda(s_f, s_i))^\Lambda \cdot \mathcal{T}^\Lambda(\alpha, \lambda)), \tag{3}$$

where $\alpha = (S, s_f, s_i)$ denotes collectively the spin variables. The irreducible tensor amplitudes $\mathcal{T}_\nu^\Lambda(\alpha, \lambda)$ of rank Λ, which characterize the reaction, are given by

$$\mathcal{T}_\nu^\Lambda(\alpha, \lambda) = W(ss_f\Lambda s_i; S\lambda)[\lambda] \sum_\beta \sum_j T_{\alpha,\beta}^j\, W(s_i l_i SL; j\Lambda) \times ((Y_l(\hat{\boldsymbol{q}}) \otimes Y_{l_f}(\hat{\boldsymbol{p}}_f))^L \otimes Y_{l_i}(\hat{\boldsymbol{p}}_i))_\nu^\Lambda, \tag{4}$$

in terms of the partial-wave amplitudes

$$T_{\alpha,\beta}^j = (4\pi)^3(-1)^{L+l_i+s_i-j}[j]^2[S][s]^{-1}[s_f]^{-1} \times \langle(((ll_f)L(ss_f)S)j||T||(l_is_i)j\rangle, \tag{5}$$

which depend on E and invariant mass W of the final *NN* system. Total angular momentum j is conserved, and $\beta = (l, l_f, L, l_i)$ denotes collectively the orbital angular momentum l of the emitted meson, the initial and final relative orbital angular momenta l_i and l_f of the *NN* system, and the total orbital angular momentum L in the final state, which takes values $L = |l_f - l|, \ldots, (l_f + l)$. It may be noted that our coupling of angular momenta in the final state differs from that used by Meyer *et al.* [17] in the case of the production of a meson with spin $s = 0$. The notation $[\lambda] = \sqrt{2\lambda + 1}$ is used apart from standard notation [26]. The preceding formalism is readily extendable to arbitrary charge states of hadrons in $NN \to NNx$, where x represents a meson with isospin I_s, if we identify

$$T_{\alpha,\beta}^j = \sum_{I_i,I_f} C\left(\tfrac{1}{2}\tfrac{1}{2}I_i; \nu_1^i\nu_2^i\nu_i\right)C\left(\tfrac{1}{2}\tfrac{1}{2}I_f; \nu_1^f\nu_2^f\nu_f\right) \times C(I_fI_sI_i; \nu_f\nu_s\nu_i)T_{\alpha,\beta}^{I_fI_i\,j}, \tag{6}$$

where I_i and I_f denote, respectively, the initial- and final-isospin quantum numbers of the *NN* system. We have $I_i = I_f = \nu_i = \nu_f = 1$, here with $I_s = 0$. The Pauli exclusion principle and parity conservation restrict the summations in Eqs. (3) and (4) to terms satisfying $(-1)^{l_i+s_i+I_i} = -1 = (-1)^{l_f+s_f+I_f}; (-1)^{l_i} = \pi(-1)^{l_f+l}$. Thus the contributing partial waves in $pp \to pp\omega$ at and near threshold may be taken as shown in Table I, where we use the same notation as in [17], viz., $S, P, D, \ldots$, for $l_i, l_f = 0, 1, 2, \ldots$, and $s, p, d, \ldots$, for $l = 0, 1, 2, \ldots$. We use $\mathcal{S}, \mathcal{P}, \mathcal{D}, \ldots$, for $L = 0, 1, 2, \ldots$, in the final state. We now express

$$\rho^i = \sum_{s_i,s_i'=0}^{1} \sum_{k=|s_i-s_i'|}^{(s_i+s_i')} (S^k(s_i, s_i') \cdot I^k(s_i, s_i')) \tag{7}$$

in terms of irreducible tensor operators $S_\nu^k(s_i, s_i')$ and the initial polarization tensors

$$I_\nu^k(s_i, s_i') = \sum_{k_1,k_2=0}^{1} F\,(P^{k_1} \otimes Q^{k_2})_\nu^k, \tag{8}$$

TABLE I. The irreducible tensor amplitudes and the partial-wave contributions to $pp \to pp\omega$ close to threshold.

$\mathcal{T}_\nu^\Lambda(\alpha, \lambda)$	l_f	l	L	s_f	S	j	l_i	s_i	$T_{\alpha,\beta}^j$	Initial pp state	Final $pp\omega$ state
$\mathcal{T}_\nu^1(101; 1)$	0	0	0	0	1	1	1	1	$T_{101;0001}^1$	$^3P_1$	$(^1Ss)^3\mathcal{S}_1$
$\mathcal{T}_\nu^1(100; 0)$	0	1	1	0	1	0	0	0	$T_{100;1010}^0$	$^1S_0$	$(^1Sp)^3\mathcal{P}_0$
	0	1	1	0	1	2	2	0	$T_{100;1012}^2$	$^1D_2$	$(^1Sp)^3\mathcal{P}_2$
$\mathcal{T}_\nu^1(110; 1)$	1	0	1	1	1	0	0	0	$T_{110;0110}^0$	$^1S_0$	$(^3Ps)^3\mathcal{P}_0$
	1	0	1	1	1	2	2	0	$T_{110;0112}^2$	$^1D_2$	$(^3Ps)^3\mathcal{P}_2$
$\mathcal{T}_\nu^2(210; 1)$	1	0	1	1	2	2	2	0	$T_{210;0112}^2$	$^1D_2$	$(^3Ps)^5\mathcal{P}_2$

of rank k, using the notation $P_0^0 = Q_0^0 = 1$ and P_ν^1, Q_ν^1 to denote the spherical components of $\boldsymbol{P}$ and $\boldsymbol{Q}$, respectively, and the factor

$$F = \tfrac{1}{2}(-1)^{k_1+k_2-k}[k_1][k_2][s_i'] \begin{Bmatrix} \frac{1}{2} & \frac{1}{2} & s_i \\ \frac{1}{2} & \frac{1}{2} & s_i' \\ k_1 & k_2 & k \end{Bmatrix}. \tag{9}$$

Using known properties [25] of the irreducible tensor operators and standard Racah techniques, we have

$$d^2\sigma = \sum_{\alpha,\alpha',\Delta,k} G\ \left(I^k(s_i,s_i')\cdot\mathcal{B}^k(s_i,s_i')\right), \tag{10}$$

in terms of the bilinear irreducible tensors

$$\mathcal{B}_\nu^k\ (s_i,s_i') = \tfrac{2\pi D}{v}\left(\mathcal{T}^\Lambda(\alpha,\lambda)\otimes\mathcal{T}^{\dagger\Lambda'}(\alpha',\lambda')\right)_\nu^k \tag{11}$$

of rank k and the geometrical factors

$$G = \delta_{s_f s_f'}[s_f]^2[s_i][s]^2(-1)^{\lambda+\lambda'+\Lambda'}[\lambda][\Lambda][\lambda'][\Lambda'] \times W(s\lambda\Lambda' k;\Lambda\lambda')W(s_i' k s_f\lambda; s_i\lambda'), \tag{12}$$

where $\mathcal{T}_\nu^{\dagger\Lambda}(\alpha,\lambda)$ and the complex conjugates $\mathcal{T}_\nu^\Lambda(\alpha,\lambda)^*$ of Eq. (4) are related through $\mathcal{T}_\nu^{\dagger\Lambda}(\alpha,\lambda) = (-1)^\nu\mathcal{T}_{-\nu}^\Lambda(\alpha,\lambda)^*$ and $\Delta = (\lambda,\lambda',\Lambda,\Lambda')$.

Defining the partial contributions to $d^2\sigma$ through $d^2\sigma = \sum_{s_i,s_i'} d^2\sigma(s_i,s_i')$ and using Eq. (8), we have

$$d^2\sigma(0,0) = d^2\sigma_0\tfrac{1}{4}(1-\boldsymbol{P}\cdot\boldsymbol{Q})\big[1+\sqrt{3}A_0^0(11)\big], \tag{13}$$

$$\begin{aligned} d^2\sigma(1,1) = d^2\sigma_0\big[&\tfrac{1}{4}(3+\boldsymbol{P}\cdot\boldsymbol{Q})\big(1-\tfrac{1}{\sqrt{3}}A_0^0(11)\big) \\ &+\tfrac{1}{2}((\boldsymbol{P}+\boldsymbol{Q})\cdot(\boldsymbol{A}(10)+\boldsymbol{A}(01))) \\ &+((P^1\otimes Q^1)^2\cdot A^2(11))\big], \end{aligned} \tag{14}$$

$$\begin{aligned} &d^2\sigma(1,0)+d^2\sigma(0,1) \\ &\quad= d^2\sigma_0\big[\tfrac{1}{2}((\boldsymbol{P}-\boldsymbol{Q})\cdot(\boldsymbol{A}(10)-\boldsymbol{A}(01))) \\ &\qquad+((P^1\otimes Q^1)^1\cdot A^1(11))\big], \end{aligned} \tag{15}$$

which add up to give Eq. (10) in the form

$$\begin{aligned} d^2\sigma = d^2\sigma_0[1&+\boldsymbol{P}\cdot\boldsymbol{A}(10)+\boldsymbol{Q}\cdot\boldsymbol{A}(01) \\ &+\sum_{k=0}^{2}((P^1\otimes Q^1)^k\cdot A^k(11))], \end{aligned} \tag{16}$$

where the unpolarized double-differential cross section,

$$d^2\sigma_0 = \frac{1}{4}\sum_{\alpha,\lambda,\Lambda}(-1)^\Lambda[s_f]^2[s]^2[\Lambda]\mathcal{B}_0^0(s_i,s_i), \tag{17}$$

is denoted as $\sigma_0(\xi)$ in [17]. The beam and target analyzing powers $\boldsymbol{A}(01)$ and $\boldsymbol{A}(10)$ are represented by the irreducible tensors $A_\nu^1(10)$ and $A_\nu^1(01)$, respectively, and the spin correlations by $A_\nu^k(11)$ of rank $k = 0, 1, 2$. We have

$$d^2\sigma_0 A_\nu^k(k_1k_2) = \sum_{\alpha,\alpha',\Delta} FG\mathcal{B}_\nu^k(s_i,s_i'). \tag{18}$$

Our $A_\nu^k(k_1k_2)$ are given, in terms of the notation of Meyer *et al.* [17], by

$$A_0^1(10) = A_{z0}(\xi), \quad A_0^1(01) = A_{0z}(\xi),$$

$$\begin{aligned} A_{\pm1}^1(10) &= \mp\tfrac{1}{\sqrt{2}}[A_{x0}(\xi)\pm iA_{y0}(\xi)], \\ A_{\pm1}^1(01) &= \mp\tfrac{1}{\sqrt{2}}[A_{0x}(\xi)\pm iA_{0y}(\xi)], \\ A_0^0(11) &= -\tfrac{1}{\sqrt{3}}[A_\Sigma(\xi)+A_{zz}(\xi)], \\ A_0^1(11) &= -\tfrac{i}{\sqrt{2}}A_\Xi(\xi), \\ A_{\pm1}^1(11) &= \tfrac{1}{2}[(A_{xz}(\xi)-A_{zx}(\xi))\pm i(A_{yz}(\xi)-A_{zy}(\xi))], \\ A_0^2(11) &= \tfrac{1}{\sqrt{6}}[2A_{zz}(\xi)-A_\Sigma(\xi)], \\ A_{\pm1}^2(11) &= \mp\tfrac{1}{2}[(A_{xz}(\xi)+A_{zx}(\xi))\pm i(A_{yz}(\xi)+A_{zy}(\xi))], \\ A_{\pm2}^2(11) &= \tfrac{1}{2}[A_\Delta(\xi)\pm i(A_{xy}(\xi)+A_{yx}(\xi))], \end{aligned} \tag{19}$$

where $A_{ij}(\xi)$, $i, j = 0, x, y, z$ are the same as in Eq. (4) of [17] and A_Σ, A_Δ, and A_Ξ are defined by Eq. (5) of [17].

At a $\vec{p}\vec{p}$ facility similar to Polarized INternal Target Experiment (PINTEX) at Indiana University Cyclotron Facility (IUCF), but with sufficiently high energies E, it should therefore be possible to determine Eqs. (13)–(15) individually, apart from Eqs. (16) and (17). It is interesting to note from Table I that only the $\mathcal{T}_\nu^1(101;1)$ from the initial state $^3P_1$ contribute to Eq. (14) and hence $|T_{101;0001}^1|^2$ can be determined empirically, whereas Eq. (15) obtains contributions to the interference of $T_{101;0001}^1$ with all the other five singlet amplitudes, which by themselves determine Eq. (13). Moreover, we note that $d^2\sigma_0$ given by Eq. (17) may itself be decomposed into $\sum_{s_i,m_i}{}^{2s_i+1}(d^2\sigma_0)_{m_i}$, where

$$^1(d^2\sigma_0)_0 = \frac{d^2\sigma_0}{4}\big[1+\sqrt{3}A_0^0(11)\big], \tag{20}$$

$$^3(d^2\sigma_0)_0 = \frac{d^2\sigma_0}{4}\big[1-\tfrac{1}{\sqrt{3}}A_0^0(11)-\tfrac{2\sqrt{2}}{\sqrt{3}}A_0^2(11)\big], \tag{21}$$

$$^3(d^2\sigma_0)_{\pm1} = \frac{d^2\sigma_0}{4}\big[1-\tfrac{1}{\sqrt{3}}A_0^0(11)+\sqrt{\tfrac{2}{3}}A_0^2(11)\big], \tag{22}$$

which represent physically the double-differential cross section for $pp\to pp\omega$ from the initial-spin states $|00\rangle$ and $|1m\rangle$, $m = 0, \pm1$. Clearly, measurements of $\sigma_0(\xi)$, A_{zz}, and A_Σ are sufficient to determine Eqs. (20)–(22) individually.

Finally, we may characterize the state of polarization of the ω meson in $pp\to pp\vec{\omega}$ by the density matrix ρ^s, whose elements are given by

$$\rho_{\mu\mu'}^s = \frac{2\pi D}{v}\frac{1}{4}\sum_{s_f}\sum_{m_f}\langle s\mu; s_f m_f|\boldsymbol{T}\boldsymbol{T}^\dagger|s\mu'; s_f m_f\rangle. \tag{23}$$

When ρ^s is expressed in the standard [27] form,

$$\rho^s = \frac{1}{2s+1}\sum_{k=0}^{2s}(\tau^k\cdot t^k), \tag{24}$$

in terms of $\tau_\nu^k \equiv S_\nu^k(s,s)$, the Fano statistical tensors t_ν^k are given by

$$\begin{aligned} t_\nu^k = \frac{1}{4}&\sum_{\alpha,\lambda,\Lambda,\Lambda'}(-1)^{\lambda-s}[s_f]^2[s]^3[\Lambda][\Lambda'] \\ &\times W(s\Lambda s\Lambda';\lambda k)\mathcal{B}_\nu^k(s_i,s_i), \end{aligned} \tag{25}$$

at the double-differential level. It may be noted that ρ^s is unnormalized so that Eq. (25) with $k = 0$ leads to Eq. (17).

The vector and tensor polarizations of ω (with $s = 1$) are readily obtained by setting $k = 1, 2$ respectively, in Eq. (25).

It is worth noting that one may measure the Fano statistical tensors t^k_ν by looking at the decay $\omega \to \pi^0\gamma$ [16], with a branching ratio of 8.92 %. The angular distribution of circularly polarized radiation emitted by polarized ω is proportional to

$$I_p(\theta_\gamma, \varphi_\gamma) = \sum_{k=0}^{2} \frac{1}{[k]} C(11k; p, -p) F_k(\theta_\gamma, \varphi_\gamma), \quad (26)$$

where $p = \pm 1$ correspond, respectively, to left and right circular polarizations as defined by Rose [26] and

$$F_k(\theta_\gamma, \varphi_\gamma) = \sum_{q=-k}^{k} (-1)^q t^k_q Y_{k-q}(\theta_\gamma, \varphi_\gamma), \quad (27)$$

where $(\theta_\gamma, \varphi_\gamma)$ denote the polar angles of the direction of γ emission in the same frame of reference in which t^k_q are given. If no observation is made on the polarization of the radiation, the intensity is proportional to

$$\sum_p I_p(\theta_\gamma, \varphi_\gamma) = \sum_{k=0,2} \frac{2}{[k]} C(11k; 1, -1) F_k(\theta_\gamma, \varphi_\gamma), \quad (28)$$

from which it is clear that the tensor polarization can be measured from the anisotropy of the angular distribution. On the other hand, the circular polarization asymmetry,

$$\Sigma(\theta_\gamma, \varphi_\gamma) = I_{-p}(\theta_\gamma, \varphi_\gamma) - I_p(\theta_\gamma, \varphi_\gamma) = \sqrt{2} F_1(\theta_\gamma, \varphi_\gamma), \quad (29)$$

enables the measurement of vector polarization.

If the polarization of the ω meson is measured with a nucleon polarized initially, we may express

$$t^k_\nu = \sum_{\nu'} \mathcal{D}(k, \nu; 1, \nu') P^1_{\nu'}; \quad k = 1, 2 \quad (30)$$

in terms of the spin transfers

$$\mathcal{D}(k, \nu; 1, \nu') = \sum_\zeta H C(1\Lambda'' k; \nu'\nu''\nu) \mathcal{B}^{\Lambda''}_{\nu''}(s_i, s'_i), \quad (31)$$

where $\zeta \equiv (\alpha, \alpha', \lambda, \lambda', \Lambda, \Lambda', \Lambda'', k')$ and

$$\begin{aligned} H = &-\frac{1}{8}\sqrt{\frac{3}{2}} (-1)^{\lambda+\lambda'+k'-k} [s]^3 [s_f]^2 [s_i][s'_i][\lambda] \\ &\times [\lambda'][\Lambda][\Lambda'][\Lambda''][k']^2 W(s\lambda k' 1; \Lambda\lambda') \\ &\times W(s'_i 1 s_f \lambda; s_i \lambda') W(1\Lambda k \Lambda'; k'\Lambda'') \\ &\times W(s\lambda' k \Lambda'; k' s) W\left(s'_i \tfrac{1}{2} 1 \tfrac{1}{2}; \tfrac{1}{2} s_i\right) \end{aligned} \quad (32)$$

if the beam is polarized. If the target is polarized, we may replace P^1_ν with Q^1_ν in Eq. (30) and attach a factor $(-1)^{s'_i - s_i}$ to H.

Denoting the six $T^j_{\alpha\beta}$ in Table I serially as $T_1 - T_6$, the irreducible tensor amplitudes $\mathcal{T}^\Lambda_\nu(\alpha, \lambda)$ that describe $pp \to pp\omega$ close to threshold are explicitly given by

$$\mathcal{T}^1_\nu(101; 1) = \frac{1}{24\pi^{3/2}} T_1 \delta_{\nu 0}, \quad (33)$$

$$\mathcal{T}^1_\nu(100; 0) = \frac{1}{12\pi} \left[T_2 + \frac{3\nu^2 - 2}{\sqrt{10}} T_3 \right] Y_{1\nu}(\hat{\boldsymbol{q}}), \quad (34)$$

$$\mathcal{T}^1_\nu(110; 1) = \frac{1}{12\pi} \left[T_4 + \frac{3\nu^2 - 2}{\sqrt{10}} T_5 \right] Y_{1\nu}(\hat{\boldsymbol{p}}_f), \quad (35)$$

$$\mathcal{T}^2_\nu(210; 1) = \frac{1}{20\sqrt{6}\pi} \nu(4 - \nu^2)^{1/2} T_6 Y_{1\nu}(\hat{\boldsymbol{p}}_f), \quad (36)$$

as $Y_{l_i m_i}(\hat{\boldsymbol{p}}_i) = ([l_i]/\sqrt{4\pi})\delta_{m_i 0}$, if we choose the beam direction as the z axis. All the observables considered in the preceding discussion are readily evaluated by use of Eqs. (33)–(36) in terms of the six partial-wave amplitudes and the angles characterizing $\boldsymbol{q}$ and $\boldsymbol{p}_f$. The unpolarized differential cross section measured in [23] is readily evaluated after Eq. (17) is integrated with respect to $d\Omega_{p_f} d\epsilon$, where $\epsilon = W - 2M$, and we have

$$d\sigma_0 = a_0 + a_2 \cos^2\theta, \quad (37)$$

where a_0 derives contributions from all the irreducible tensor amplitudes, whereas $\mathcal{T}^1_\nu(100; 0)$ alone, which produces the meson in the p wave, contributes to a_2. The existing data [23] are in good agreement with the form of Eq. (37), which hence provides clear evidence for the presence of the initial-spin singlet amplitude $\mathcal{T}^1_\nu(100; 0)$ given by Eq. (34) in addition to the initial-spin triplet threshold amplitude $\mathcal{T}^1_\nu(101; 1)$ given by Eq. (33). If we can assume that the contribution of $\mathcal{T}^1_\nu(110; 1)$ and $\mathcal{T}^2_\nu(210; 1)$ is small or negligible, a_0 and a_2 involve the bilinear combinations $[|T_1|^2 + 3|T_2 + \frac{1}{\sqrt{10}} T_3|^2]$ and $[\,|T_3|^2 - 2\sqrt{10}\Re(T_2 T_3^*)]$ of the partial-wave amplitudes duly integrated with respect to ϵ. If we measure not only the angular distribution of ω but also its energy, we can dispense with the integration with respect to ϵ.

Integrating the right-hand side of Eq. (18) with respect to $d\Omega_{p_f}$ and equating it to $d\sigma_0 A^k_\nu(k_1 k_2)$ defines the analyzing powers at the d^3q level. It is interesting to note that the Wigner $9j$ symbol in Eq. (9) ensures that the initial-spin triplet amplitude [Eq. (33)] alone contributes to $A^2_0(11)$, a measurement of which determines $|T_1|^2$. Knowledge of $|T_1|^2$ leads to a determination of $|T_2 + \frac{1}{\sqrt{10}} T_3|^2$ with the preceding expression used for a_0. Moreover, it is interesting to note that $\boldsymbol{A}(10) - \boldsymbol{A}(01)$ or $\boldsymbol{A}(11)$ are proportional to the interference of the initial-spin triplet amplitude $\mathcal{T}^1_\nu(101; 1)$ with the initial-spin singlet amplitude $\mathcal{T}^1_\nu(100; 0)$. This leads to a bilinear involving T_1 with $T_2 + \frac{1}{\sqrt{10}} T_3$. Likewise, t^k_ν at the d^3q level are also obtained on integration of Eq. (25) or (30) with respect to $d\Omega_{p_f}$. There is as yet no data available on any of the spin observables.

We thank the referee for helpful suggestions. G.R. is grateful to B.V. Sreekantan, R. Cowsik, and J.H. Sastry for facilities provided for research at the Indian Institute of Astrophysics; M.S.V. acknowledges support of the Department of Science and Technology, India; P.N.D. is thankful to the Alexander von Humboldt Foundation for the award of a Fellowship; and J.B. acknowledges encouragement for research given by T. G. S. Moorthy and the management of K.S. Institute of Technology.

[1] H. O. Meyer *et al.*, Phys. Rev. Lett. **65**, 2846 (1990); Nucl. Phys. **A539**, 633 (1992); S. Stanislaus, D. Horvath, D. F. Measday, A. J. Noble, and M. Salomon, Phys. Rev. C **41**, R1913 (1990); E. Korkmaz *et al.*, Nucl. Phys. **A535**, 637 (1991); A. Bondar *et al.*, Phys. Lett. **B356**, 8 (1995); W. W. Daehnick *et al.*, Phys. Rev. Lett. **74**, 2913 (1995).

[2] D. S. Koltun and A. Reitan, Phys. Rev. **141**, 1413 (1966); M. E. Schillaci, R. R. Silbar, and J. E. Young, Phys. Rev. Lett. **21**, 711 (1968); Phys. Rev. **179**, 1539 (1969).

[3] K. Nakayama, Matter Mater. **11**, 119 (2002); K. Tsushima and K. Nakayama, Phys. Rev. C **68**, 034612 (2003).

[4] R. Machleidt, Adv. Nucl. Phys. **19**, 189 (1989).

[5] H. Machner and J. Haidenbauer, J. Phys. G. **25**, 231(R) (1999); P. Moskal, M. Wolke, A. Khoukaz, and W. Oelert, Prog. Part. Nucl. Phys. **49**, 1 (2002); C. Hanhart, Phys. Rep. **397**, 155 (2004).

[6] S. Okubo, Phys. Lett. **B5**, 165 (1963); G. Zweig, CERN Report 8419/TH412 (1964); I. Iizuka, Prog. Theor. Phys. Suppl. **37-38**, 21 (1966).

[7] F. Balestra *et al.*, Phys. Rev. Lett. **81**, 4572 (1998); F. Balestra *et al.*, Phys. Rev. C **63**, 024004 (2001).

[8] C. Amsler, Rev. Mod. Phys. **70**, 1293 (1998); M. G. Sapozhnikov, Nucl. Phys. **A655**, 151c (1999).

[9] H. J. Lipkin, Phys. Lett. **B60**, 371 (1976); J. Ellis, M. Karliner, D. E. Kharzeev, and M. G. Sapozhnikov, Nucl. Phys. **A673**, 256 (2000).

[10] S. Okubo, Phys. Rev. D **16**, 2336 (1977); H. J. Lipkin, Int. J. Mod. Phys. E **1**, 603 (1992); V. P. Nomokov and M. G. Sapozhnikov, Phys. Part. Nucl. **34**, 94 (2003).

[11] K. Nakayama, A. Szczurek, C. Hanhart, J. Haidenbauer, and J. Speth, Phys. Rev. C **57**, 1580 (1998); K. Nakayama, J. W. Durso, J. Haidenbauer, C. Hanhart, and J. Speth, *ibid.* **60**, 055209 (1999); K. Nakayama, H. F. Arellano, J. W. Durso, and J. Speth, *ibid.* **61**, 024001 (1999); K. Nakayama, J. Haidenbauer, and J. Speth, *ibid.* **63**, 015201 (2000); Nucl. Phys. **A689**, 402c (2001); K. Nakayama, J. Speth, and T.-S. H. Lee, Phys. Rev. C **65**, 045210 (2002); K. Tsushima and K. Nakayama, Nucl. Phys. **A721**, 633 (2003).

[12] A. I. Titov, B. Kämpfer, and V. V. Shklyar, Phys. Rev. C **59**, 999 (1999); N. Kaiser, *ibid.* **60**, 057001 (1999); A. I. Titov, B. Kämpfer, and B. L. Reznik, Eur. Phys. J. A **7**, 543 (2000); V. Yu. Grishina, L. A. Kondratyuk, and M. Büscher, Phys. At. Nucl. **63**, 1824 (2000).

[13] A. Faessler, C. Fuchs, and M. I. Krivoruchenko, Phys. Rev. C **61**, 035206 (2000); C. Fuchs, M. I. Krivoruchenko, H. L. Yadav, A. Faessler, B. V. Martemyanov, and K. Shekhter, *ibid.* **67**, 025202 (2003); A. Faessler, C. Fuchs, M. I. Krivoruchenko, and B. V. Martemyanov, *ibid.* **70**, 035211 (2004); C. Fuchs, A. Faessler, D. Cozma, B. V. Martemyanov, and M. I. Krivoruchenko, Nucl. Phys. **A755**, 499 (2005).

[14] N. Isgur and G. Karl, Phys. Rev. D **18**, 4187 (1978); **19**, 2653 (1979); N. Isgur, *ibid.* **23**, 817 (1981); R. Koniuk and N. Isgur, *ibid.* **21**, 1868 (1980); Phys. Rev. Lett. **44**, 845 (1980); S. Capstick and N. Isgur, Phys. Rev. D **34**, 2809 (1986).

[15] T. Barns and H. P. Morsch (Eds.), *Baryon Excitations*, Lectures of COSY workshop held at Forschungszentrum Jülich, 2–3 May 2000, ISBN 3-89336-273-8 (2000); C. Carlson and B. Mecking, in *Proceedings of the International Conference on the Structure of Baryons*, BARYONS 2002, Newport News, VA, 3–8 March 2002 (World Scientific, Singapore, 2003); S. A. Dytman and E. S. Swanson, *Proceedings of NSTAR 2002 Workshop*, Pittsburg, PA, 9–12 October 2002 (World Scientific, Singapore, 2003).

[16] S. Eidelman *et al.*, Phys. Lett. **B592**, 1 (2004).

[17] H. O. Meyer *et al.*, Phys. Rev. C **63**, 064002 (2001); W. W. Daehnick *et al.*, *ibid.* **65**, 024003 (2002).

[18] G. Ramachandran, P. N. Deepak, and M. S. Vidya, Phys. Rev. C **62**, 011001(R) (2000); G. Ramachandran and P. N. Deepak, *ibid.* **63**, 051001(R) (2001).

[19] P. N. Deepak and G. Ramachandran, Phys. Rev. C **65**, 027601 (2002).

[20] S. M. Bilenky and R. M. Ryndin, Phys. Lett. **6**, 217 (1963).

[21] G. Ramachandran and P. N. Deepak, J. Phys. G. **26**, 1809 (2000).

[22] F. Hibou *et al.*, Phys. Rev. Lett. **83**, 492 (1999).

[23] S. Abd El-Samad *et al.* (COSY-TOF Collaboration), Phys. Lett. **B522**, 16 (2001).

[24] F. Rathmann *et al.*, Study of heavy meson production in *NN* collisions with polarized beam and target, letter of intent to COSY-PAC No. 81 (Jülich, 1999); Progress Report on Intermediate Energy Spin Physics 213; We-Hereaus Seminar, Jülich, 18–20 November 1998, edited by F. Rathmann, W. T. H. van Oers, and C. Wilkin (Jülich, 1998); F. Rathmann *et al.*, Czech. J. Phys. **52**, 319c (2002).

[25] G. Ramachandran and M. S. Vidya, Phys. Rev. C **56**, R12 (1997).

[26] M. E. Rose, *Elementary Theory of Angular Momentum* (Wiley, New York, 1957), republished by Dover, NY, 1995; D. A. Varshalovich, A. N. Moskalev, and V. K. Khersonskii, *Quantum Theory of Angular Momentum* (World Scientific, Singapore, 1988).

[27] G. R. Satchler *et al.*, in *Proceedings of the Third International Symposium on Polarization Phenomena in Nuclear Reactions*, edited by H. H. Barschall and W. Haeberli (University of Wisconsin Press, Madison, WI, 1970).