ω Production in $p\,p$ Collisions

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

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Experimental study of meson production in NN collisions has attracted considerable interest during the last 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 an analysis involves, therefore, 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 $0.2fm$. 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 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 to the theoretical estimate [9] of 4.2×10^{-3} after correcting for the available phase space. We may refer [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, 12, 13] to vector meson production in NN collisions. The constituent quark models [14] predict highly excited N^* states which 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 to 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 which 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 $pp \to pp\pi^0$, for example, the transition of the pp system at threshold is from an initial spin triplet to a final spin singlet state (${}^3P_0 \to {}^1S_0$). 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. Conclusive theoretical interpretation of all these data have remained elusive, although the model calculations appear to do better in the case of charged pion production as compared to the neutral pion production and the agreement even there seemed 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] which has been developed to analyze measurements on $\vec{p}\vec{p} \to pp\pi^0$ 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 using the experimental results of Meyer *et al.* [17]. The above 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 also apart from measuring the angular distributions in polarized beam and polarized target experiments. Experimental data on total [22] and differential [23] cross-sections for $pp \to pp\omega$ have already been published and proposals are underway [5, 24] to study heavy meson production in NN collisions using polarized beams and targets at 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

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

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. Notation $\sigma(\xi, \boldsymbol{P}, \boldsymbol{Q})$ is used in [17] to denote (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 $\mathcal{T}^\Lambda_\nu(\alpha,\lambda)$ of rank Λ, which characterize the reaction, are given by

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

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

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

in terms of the partial wave amplitudes

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

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 notations [26]. The above 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

$$T^j_{\alpha,\beta} = \sum_{I_i,I_f} C(\tfrac{1}{2}\tfrac{1}{2}I_i;\nu^i_1\nu^i_2\nu_i)C(\tfrac{1}{2}\tfrac{1}{2}I_f;\nu^f_1\nu^f_2\nu_f) \times C(I_fI_sI_i;\nu_f\nu_s\nu_i)T^{I_fI_i\,j}_{\alpha,\beta}, \tag{6}$$

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 $I_s = 0$. Pauli exclusion principle and parity conservation restrict the summations in (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 notations 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 $S^k_\nu(s_i, s'_i)$ and the initial polarization tensors

$$I^k_\nu(s_i, s'_i) = \sum_{k_1,k_2=0}^{1} F\ (P^{k_1} \otimes Q^{k_2})^k_\nu, \tag{8}$$

of rank k, using the notations $P^0_0 = Q^0_0 = 1$ and P^1_ν, Q^1_ν to denote the spherical components of $\boldsymbol{P}, \boldsymbol{Q}$ 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

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

in terms of the bilinear irreducible tensors

$$\mathcal{B}^k_\nu(s_i, s'_i) = \frac{2\pi D}{v}(\mathcal{T}^\Lambda(\alpha,\lambda) \otimes \mathcal{T}^{\dagger\Lambda'}(\alpha',\lambda'))^k_\nu, \tag{11}$$

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

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

$$\begin{aligned} \mathrm{d}^2\sigma(0,0) &= \mathrm{d}^2\sigma_0 \tfrac{1}{4}(1 - \boldsymbol{P}\cdot\boldsymbol{Q})[1 + \sqrt{3}A^0_0(11)], \quad (13) \\ \mathrm{d}^2\sigma(1,1) &= \mathrm{d}^2\sigma_0[\tfrac{1}{4}(3 + \boldsymbol{P}\cdot\boldsymbol{Q})(1 - \tfrac{1}{\sqrt{3}}A^0_0(11)) \\ &\quad + \tfrac{1}{2}((\boldsymbol{P}+\boldsymbol{Q})\cdot(\boldsymbol{A}(10) + \boldsymbol{A}(01))) \\ &\quad + ((P^1 \otimes Q^1)^2 \cdot A^2(11))], \quad (14) \\ \mathrm{d}^2\sigma(1,0) &+ \mathrm{d}^2\ \sigma(0,1) = \mathrm{d}^2\sigma_0 \\ &\quad \times[\tfrac{1}{2}((\boldsymbol{P}-\boldsymbol{Q})\cdot(\boldsymbol{A}(10) - \boldsymbol{A}(01))) \\ &\quad + ((P^1 \otimes Q^1)^1 \cdot A^1(11))], \quad (15) \end{aligned}$$

which add up to give (10) in the form

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

where the unpolarized double differential cross section

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

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

$$\mathrm{d}^2\sigma_0 A^k_\nu(k_1k_2) = \sum_{\alpha,\alpha',\Delta} F\ G\ \ \mathcal{B}^k_\nu(s_i, s'_i). \tag{18}$$

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

$$\begin{aligned} A^1_0(10) &= A_{z0}(\xi),\ A^1_0(01) = A_{0z}(\xi), \\ A^1_{\pm1}(10) &= \mp\tfrac{1}{\sqrt{2}}[A_{x0}(\xi) \pm iA_{y0}(\xi)], \\ A^1_{\pm1}(01) &= \mp\tfrac{1}{\sqrt{2}}[A_{0x}(\xi) \pm iA_{0y}(\xi)], \\ A^0_0(11) &= -\tfrac{1}{\sqrt{3}}[A_\Sigma(\xi) + A_{zz}(\xi)], \\ A^1_0(11) &= -\tfrac{i}{\sqrt{2}}A_\Xi(\xi), \\ A^1_{\pm1}(11) &= \tfrac{1}{2}[(A_{xz}(\xi) - A_{zx}(\xi)) \pm i(A_{yz}(\xi) - A_{zy}(\xi))], \\ A^2_0(11) &= \tfrac{1}{\sqrt{6}}[2A_{zz}(\xi) - A_\Sigma(\xi)], \\ A^2_{\pm1}(11) &= \mp\tfrac{1}{2}[(A_{xz}(\xi) + A_{zx}(\xi)) \pm i(A_{yz}(\xi) + A_{zy}(\xi))], \\ A^2_{\pm2}(11) &= \tfrac{1}{2}[A_\Delta(\xi) \pm i(A_{xy}(\xi) + A_{yx}(\xi))], \end{aligned} \tag{19}$$

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

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

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

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

$${}^3(\mathrm{d}^2\sigma_0)_{\pm1} = \tfrac{\mathrm{d}^2\sigma_0}{4}[1 - \tfrac{1}{\sqrt{3}}A^0_0(11) + \sqrt{\tfrac{2}{3}}A^2_0(11)], \tag{22}$$

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 (20) to (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

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

Expressing ρ^s 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 $\tau^k_\nu \equiv S^k_\nu(s,s)$,the Fano statistical tensors t^k_ν are given by

$$t^k_\nu = \frac{1}{4}\sum_{\alpha,\lambda,\Lambda,\Lambda'}(-1)^{\lambda-s}[s_f]^2[s]^3[\Lambda][\Lambda'] \times W(s\Lambda s\Lambda';\lambda k)\mathcal{B}^k_\nu(s_i,s_i), \quad (25)$$

at the double differential level. It may be noted that ρ^s is unnormalized so that (25) with $k=0$ leads to (17). The vector and tensor polarizations of ω (with $s=1$) are readily obtained by setting $k=1,2$ respectively in (25).

It is worth noting that the Fano statistical tensors t^k_ν may be measured 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 measurement of vector polarization.

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

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

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

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

if the beam is polarized. If the target is polarized, we may replace P^1_ν by Q^1_ν in (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 to T_6, the irreducible tensor amplitudes $\mathcal{T}^\Lambda_\nu(\alpha,\lambda)$ which describe $pp\to pp\omega$ close to threshold are explicitly given by

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

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

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

$$\mathcal{T}^2_\nu(210;1) = \frac{1}{20\sqrt{6}\pi}\nu(4-\nu^2)^{1/2}T_6Y_{1\nu}(\hat{\boldsymbol{p}}_f), \quad (36)$$

since $Y_{l_i m_i}(\hat{\boldsymbol{p}}_i) = ([l_i]/\sqrt{4\pi})\delta_{m_i 0}$, if we choose the beam direction as the z-axis. All the observables considered in the above discussion are readily evaluated using (33) to (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 integrating (17) with respect to $\mathrm{d}\Omega_{p_f}\mathrm{d}\epsilon$, where $\epsilon = W-2M$, and we have

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

where a_0 derives contributions from all the irreducible tensor amplitudes, while $\mathcal{T}^1_\nu(100;0)$ alone, which produces the meson in $p-$wave, contributes to a_2. The existing data [23] is in good agreement with the form (37), which hence provides clear evidence for the presence of the initial spin singlet amplitude $\mathcal{T}^1_\nu(100;0)$ given by (34) in addition to the initial spin triplet threshold amplitude $\mathcal{T}^1_\nu(101;1)$ given by (33). If we can assume the contribution of $\mathcal{T}^1_\nu(110;1)$ and $\mathcal{T}^2_\nu(210;1)$ to be 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 $[\,|T_3|^2-2\sqrt{10}\Re(T_2T_3^*)]$ of the partial wave amplitudes duly integrated with respect to ϵ. If one measures not only the angular distribution of ω but also its energy, the integration with respect to ϵ can be dispensed with.

Integrating the right hand side of (18) with respect to $\mathrm{d}\Omega_{p_f}$ and equating it to $\mathrm{d}\sigma_0 A^k_\nu(k_1k_2)$ defines the analyzing powers at the d^3q level. It is interesting to note that the Wigner $9j$ symbol in (9) ensures that the initial spin triplet amplitude (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 $|T_2+\frac{1}{\sqrt{10}}T_3|^2$ using the above expression 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 d^3q level are also obtained on integration of (25) or (30) with respect to $\mathrm{d}\Omega_{p_f}$. There is as yet no data available on any of the spin observables.

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