Singlet, triplet, and octet axial-vector form factors of the spin $-\frac{3}{2}^+$ decuplet baryons in the chiral quark constituent model

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The axial-vector form factors of the spin $-\frac{3}{2}^+$ decuplet baryons are investigated in the chiral constituent quark model using their explicit quark spin polarizations. The quark sea arises from the chiral symmetry breaking, which results in the Goldstone bosons mediating the interaction between constituent quarks. The axial-vector form factors, which have some physical significance corresponding to the flavor singlet current, flavor isovector (triplet) current, and the flavor hypercharge axial (octet) current at zero momentum transfer, are, respectively $G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$, $G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$, and $G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$. In order to further understand the Q^2 dependence of these form factors, we have used the dipole form of parametrization. The qualitative and quantitative contribution of the quark sea has also been investigated by varying the transition probability of the chiral fluctuation.

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I. INTRODUCTION

The quest to understand the internal structure of the baryons in terms of effective degrees of freedom has stimulated both experimental and theoretical studies in the past. The measurements of polarized structure functions in the deep inelastic scattering (DIS) experiments [1–4] have thrown considerable light in describing the spin structure of the baryons. These experiments provide clear evidence that the valence quarks carry only a small fraction of its spin, and the decomposition of the baryon spin still remains to be a major unresolved issue in high energy spin physics. These composite systems are bound by increasingly strong forces for decreasing momentum transfer.

The electromagnetic current probes the Dirac and Pauli forms, whereas the isovector axial-vector current provides information on the internal structure in both strong and weak interactions. In particular, the flavor isovector (triplet) current determines precisely the neutron β decay parameters, the flavor hypercharge axial (octet) current provides the information on the hyperon β decay parameters, and the flavor singlet current can be related to the total quark spin content. The axial-vector form factors arising from the axial current are the fundamental quantities that carry the complete information in understanding the spin structure of the baryons.

The axial-vector form factors of the low-lying octet baryons are well known through their semileptonic decays over a wide region of momentum transfer squared Q^2 [5]. The axial-vector coupling constants of the nucleons can be calculated from the experimentally measured first moments and related to combinations of the spin polarizations and coupling constants. These can further be related to certain well-known sum rules derived within quantum chromodynamics (QCD). However, since the low-lying decuplet baryons decay strongly (except Ω^-), the study of their axial-vector form factors has been rather limited, and it is rather difficult to experimentally measure these quantities due to their short lifetimes. Further, the Q^2 dependence of the nucleon axial-vector form factors have been studied from the elastic scattering of neutrinos and antineutrinos [6,7] and the pion electroproduction on the proton [8] for the low-Q^2 region. Very limited data is available for the high-Q^2 region. More recently, there has been considerable progress to probe the hyperon form factors through the strangeness-production processes in the higher-energy Minerνa experiment at Fermilab [9]. These measurements will give a broader understanding by providing a refined data in a wide range of Q^2.

Since the axial-vector form factors describe the low-energy hadron phenomena in the nonperturbative regime, they can be described in models incorporating the relevant properties of QCD. There is a limitation in theoretical knowledge of these form factors, as the calculations from the first principles of QCD are still a big challenge. The lattice calculations serve as a valuable tool to determine the axial charge and form factors in a model-independent way. Even though some systematic errors still exist in the lattice calculations, a lot of refinements have been made in the recent past to remove the sources of such errors. At present, the axial-vector form factors of the nucleon are very well known in the lattice calculations [10–14], whereas lattice calculations have given some information on the properties of hyperons [15–18]. More recently, lattice QCD has provided data on the axial-vector form factors of the decuplet baryons as well [19,20]. Another powerful method is the QCD sum rules (QCDSR), where the hadron properties are estimated through the calculation of the correlation function using the operator product expansion (OPE) with Wilson coefficients and local operators [21–25]. Some more estimates have been made for the axial charges of the hyperons in the chiral perturbation theory [26–28], large Nc limit of QCD [29].

Various theoretical works have carried out the calculations for the case of decuplet baryons, including the chiral perturbation theory [30,31], Goldstone-boson-exchange relativistic constituent quark model (RCQM) [32,33], light-cone sum rules (LCSR) [34], and pertubative chiral quark model (PCQM) [35].

In light of the above developments, it is clearly of great interest to find the axial-vector form factors of the spin $-\frac{3}{2}^+$ decuplet baryons, as it will provide an important test for models that attempt to describe the low-energy properties of the baryons. Their knowledge would also undoubtedly provide vital clues to the physical interpretation of the nonperturbative aspects of QCD. Based on the successes of the chiral constituent quark model (χCQM) [36] to find the axial-vector form factors of the low-lying spin $-\frac{1}{2}^+$ octet baryons [37], we use the basic idea of chiral symmetry breaking taking place at a distance scale much smaller than the confinement scale. As a consequence of this symmetry breaking, the almost massless quarks acquire a dynamical mass coupled with internal Goldstone bosons (GBs) [38–41]. As a consequence, the exchange of GBs mediates the interaction between constituent quarks. The χCQM has been successful in giving a possible solution to the proton spin problem [41], different components of the magnetic moments of octet, and decuplet baryons, including their transitions [42], factors contributing to the violation of the Gottfried sum rule [43], and the Coleman-Glashow sum rule. Further, the hyperon β decay parameters [5] and strangeness content in the nucleon [44] have also been successfully estimated. In general, the χCQM is able to provide unique and important information about the flavor and spin distributions of the quarks in the baryons.

The purpose of the present paper is to estimate the axial-vector form factors of the spin $-\frac{3}{2}^+$ decuplet baryons using the chiral constituent quark model (χCQM). In particular, we would like to phenomenologically estimate the explicit quark spin polarizations, which are directly affected by chiral symmetry breaking parameters as well as the SU(3) symmetry breaking parameters. It would be interesting to study the extent of the contribution of the quark sea arising from the Goldstone bosons, which mediate the interaction between constituent quarks. The static properties can be computed using the axial-vector current for the matrix elements, which have some physical significance. The flavor singlet current $G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$, flavor isovector (triplet) current $G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$, and the flavor hypercharge axial (octet) current $G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$ at zero momentum transfer have been investigated for the case of Δ, Σ^*, Ξ^*, and Ω baryons. Further, in order to further understand the Q^2 dependence of these form factors, we have used the dipole form of parametrization to study $G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2)$, $G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2)$, and $G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2)$. It would also be significant to analyze the extent of the contribution of the quark sea by varying the transition probability of the chiral fluctuation. Since no experimental data is available for the case of spin $-\frac{3}{2}^+$ decuplet baryons, the results can be compared with the recent available theoretical findings.

II. AXIAL-VECTOR FORM FACTORS

The axial-vector form factors can be defined using the axial-vector current constituting the quark field. We have

$$A^{\mu,a} = \bar{\mathbf{q}}(\mathbf{x})\gamma^\mu\gamma_5\frac{\lambda^a}{2}\mathbf{q}(\mathbf{x}), \tag{1}$$

where $\mathbf{q}(\mathbf{x})$ is the quark field in flavor space for $\mathbf{q} = (u, d, s)$. Here, λ^a $(a = 1, 2, \ldots 8)$ are the well known Gell-Mann matrices describing the flavor SU(3) structure of the light quarks. In the present context of axial-vector form factors, we will consider only those matrices that have diagonal representation. λ^3 corresponds to the flavor isovector (triplet) current, and λ^8 corresponds to the flavor hypercharge axial (octet) current [45,46]. In addition to these matrices, an unit matrix $\lambda^0(=\sqrt{\frac{2}{3}}I)$ can be introduced, which will correspond to the flavor singlet current.

The axial-vector form factors can be parametrized through the matrix elements of the axial-vector current and the spin $-\frac{3}{2}^+$ decuplet baryons. We have

$$\begin{aligned}\langle B^{\frac{3}{2}^+}(p', J'_z)|A^{\mu,a}|B^{\frac{3}{2}^+}(p, J_z)\rangle \\ = \bar{u}_\rho(p', J'_z)\left[\gamma^\mu\gamma_5 G^a_{\mathrm{AV}}(Q^2)\eta^{\rho\sigma} + \frac{q^\mu}{2M_B}\gamma_5 G^a_P(Q^2)\eta^{\rho\sigma}\right] \\ \times u_\sigma(p, J_z).\end{aligned} \tag{2}$$

The Rarita-Schwinger spinor $u^\rho(p, J_z)$ here represents the spinor describing the spin $-\frac{3}{2}^+$ decuplet baryon and is a tensor product between a first rank tensor and a Dirac spinor. $\eta^{\rho\sigma}$ expressed as $\eta^{\rho\sigma} = \mathrm{diag}(-1,1,1,1)$ is representing the metric tensor of Minkowski space. If $u(p)$ represents a Dirac spinor, a spin $-$ 1 field or Lorentz vector can be constructed, and the Rarita-Schwinger spinor can be described by the combination of the polarization vector of the spin $-$ 1 field and the Dirac spinor of the spin $-\frac{1}{2}$ field. In the above equation, M_B represents the baryon mass, and $u^\sigma(p, J_z)$ $[\bar{u}_\rho(p', J'_z)]$ are the Rarita-Schwinger spinors of the initial (final) spin $-\frac{3}{2}^+$ decuplet baryon states, respectively. The four momentum transfer is given in terms of the initial and final momentum $q \equiv p - p'$, and we have $Q^2 = -q^2$. The functions $G^a_{\mathrm{AV}}(Q^2)$ and $G^a_P(Q^2)$ ($a = 0$, 3, 8) are the axial and induced pseudoscalar form factors, respectively. For the present work, we will ignore the induced pseudoscalar form factors as they are not relevant.

In general, the axial-vector matrix elements are important in hadron physics as they provide a deep insight in understanding the internal spin structure [5,41]. In order to calculate the static properties of the axial-vector form factors at zero momentum transfer, these form factors can be related to the spin polarizations. Before we present the spin polarization combinations corresponding to the singlet, triplet, and octet axial-vector form factors, we first present the calculations for the spin polarizations of the constituent quark in each decuplet baryon using the chiral constituent quark model.

III. SPIN STRUCTURE OF THE DECUPLET BARYONS IN THE CHIRAL CONSTITUENT QUARK MODEL

The internal structure of the baryons can be understood using the QCD Lagrangian, which describes the dynamics of light quarks (u, d, and s). However, under the chiral transformation for the quark fields $\psi \to \gamma^5\psi$, it does not remain invariant and changes sign because of the mass terms. If the mass terms in the QCD Lagrangian are neglected, it will have global chiral symmetry of the $\mathrm{SU}(3)_L \times \mathrm{SU}(3)_R$ group. This chiral symmetry is spontaneously broken as $\mathrm{SU}(3)_L \times \mathrm{SU}(3)_R \to \mathrm{SU}(3)_{L+R}$ at around a scale of 1 GeV. As a result, a set of massless Goldstone bosons (GBs) exist and are identified as π, K, η mesons. Since the QCD Lagrangian is also invariant under the axial $U(1)$ symmetry, the existence of η' as the ninth GB is implied. Therefore, the constituent quarks and this nonet of GBs form the appropriate degrees of freedom within the region of QCD confinement scale ($\Lambda_{\mathrm{QCD}} \simeq 0.1$–$0.3$ GeV) and the chiral symmetry breaking scale $\Lambda_{\chi\mathrm{SB}}$.

The effective Lagrangian describing interaction quarks and a nonet of GBs forms the basis for the χCQM, which was introduced by Weinberg and further developed by Manohar and Georgi [36]. The underlying idea of χCQM is the fluctuation process where the GBs couple directly to the constituent quarks in the hadron interior as

$$q^{\pm} \to \mathrm{GB} + q'^{\mp} \to (q\bar{q}') + q'^{\mp}, \tag{3}$$

where $q\bar{q}' + q'$ constitute the quark sea [38,39,41]. The effective interaction Lagrangian between GBs and quarks in the leading order can be expressed as

$$\mathcal{L}_{\mathrm{int}} = -\frac{g_A}{f_\pi}\bar{\psi}\partial_\mu\Phi\gamma^\mu\gamma^5\psi. \tag{4}$$

Here, g_A is the axial-vector coupling constant. The Lagrangian be reduced to

$$\mathcal{L}_{\mathrm{int}} \approx i\sum_{q=u,d,s}\frac{m_q + m_{q'}}{f_\pi}\bar{q}'\Phi\gamma^5 q = i\sum_{q=u,d,s}P_\pi\bar{q}'\Phi\gamma^5 q, \tag{5}$$

using the Dirac equation $(i\gamma^\mu\partial_\mu - m_q)q = 0$. Here, $P_\pi(=\frac{m_q+m_{q'}}{f_\pi})$ is the coupling constant for octet of GBs, and m_q ($m_{q'}$) is the quark mass parameter. The Lagrangian of the quark-GB interaction, suppressing all the space-time structure to the lowest order, can now be expressed as

$$\mathcal{L}_{\mathrm{int}} = P_\pi\bar{\psi}\Phi\psi. \tag{6}$$

The QCD Lagrangian is also invariant under the axial $U(1)$ symmetry, which would imply the existence of ninth GB. This breaking symmetry picks the η' as the ninth GB. The effective Lagrangian describing interaction between quarks and a nonet of GBs, consisting of an octet and a singlet, can now be expressed as

$$\mathcal{L}_{\mathrm{int}} = P_\pi\bar{\psi}\left(\Phi + P_{\eta'}\frac{\eta'}{\sqrt{3}}I\right)\psi = P_\pi\bar{\psi}(\Phi')\psi, \tag{7}$$

where P_π is the coupling constant for the octet GB, and $P_{\eta'}$ is the ratio of the coupling constants for the singlet and octet GBs. The GB field Φ' can be expressed in terms of the GBs and their transition probabilities as

$$\Phi' = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + P_\eta\frac{\eta}{\sqrt{6}} + P_{\eta'}\frac{\eta'}{\sqrt{3}} & \pi^+ & P_K K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + P_\eta\frac{\eta}{\sqrt{6}} + P_{\eta'}\frac{\eta'}{\sqrt{3}} & P_K K^o \\ P_K K^- & P_K\bar{K}^0 & -P_\eta\frac{2\eta}{\sqrt{6}} + P_{\eta'}\frac{\eta'}{\sqrt{3}} \end{pmatrix}. \tag{8}$$

From Eq. (8), all the chiral fluctuations are given in terms of the transition probabilities P_π^2, P_K^2, P_η^2, and $P_{\eta'}^2$ [38,39,41]. These parameters help us to understand the extent to which the quark sea contributes to the structure of the baryon. The transition probability P_π^2 is taken by considering the strange and nonstrange quark masses to be nondegenerate $M_s > M_{u,d}$, the transition probabilities P_K^2 and P_η^2 are taken by considering GB masses of K, η, and π to be nondegenerate $M_{K,\eta} > M_\pi$, and finally, the transition probability $P_{\eta'}^2$ is taken by considering GB masses η', K, and η to be nondegenerate $M_{\eta'} > M_{K,\eta}$. A hierarchy for the transition probabilities is important to be fixed and based on the scaling of the quark contributions, which is $\frac{1}{M_q^2}$, a constraint can be fixed as

$$P_\pi^2 > P_\pi^2 P_K^2 > P_\pi^2 P_\eta^2 > P_\pi^2 P_{\eta'}^2. \tag{9}$$

All the possible chiral fluctuations in the fluctuation process are given as

$$\begin{aligned}
u^\pm &\rightleftharpoons (d^\mp + \pi^+) + (s^\mp + K^+) + (u^\mp + \pi^0, \eta, \eta'),\\
d^\pm &\rightleftharpoons (u^\mp + \pi^-) + (s^\mp + K^0) + (d^\mp + \pi^0, \eta, \eta'),\\
s^\pm &\rightleftharpoons (u^\mp + K^-) + (d^\mp + \bar{K}^0) + (s^\mp + \eta, \eta').
\end{aligned} \tag{10}$$

The transition probability of the emission of a GB from any of the q quark, $P(q \to \mathrm{GB})$, can now be expressed in terms of the transition probabilities P_π^2, P_K^2, P_η^2, and $P_{\eta'}^2$. We have

$$P(u \to \mathrm{GB}) = P(d \to \mathrm{GB}) = \frac{P_\pi^2}{6}(9 + 6P_K^2 + P_\eta^2 + 2P_{\eta'}^2), \tag{11}$$

$$P(s \to \mathrm{GB}) = \frac{P_\pi^2}{3}(6P_K^2 + 2P_\eta^2 + P_{\eta'}^2). \tag{12}$$

The transition probability of the emission of a $q^\pm$ constituent quark to all the possible $q' = u, d, s$ quarks along with GBs $(q\bar{q}')$, $P(q^\pm \to q\bar{q}' + q'^\mp)$, as calculated from the Lagrangian, can be expressed as

$$\begin{aligned}
&P(u^\pm \to u\bar{q}' + q'^\mp)\\
&\quad = \frac{P_\pi^2}{6}(3 + P_\eta^2 + 2P_{\eta'}^2)u^\mp + P_\pi^2 d^\mp + P_\pi^2 P_K^2 s^\mp,
\end{aligned} \tag{13}$$

$$\begin{aligned}
&P(d^\pm \to d\bar{q}' + q'^\mp)\\
&\quad = P_\pi^2 u^\mp + \frac{P_\pi^2}{6}(3 + P_\eta^2 + 2P_{\eta'}^2)d^\mp + P_\pi^2 P_K^2 s^\mp,
\end{aligned} \tag{14}$$

$$\begin{aligned}
&P(s^\pm \to s\bar{q}' + q'^\mp)\\
&\quad = P_\pi^2 P_K^2 u^\mp + P_\pi^2 P_K^2 d^\mp + \frac{P_\pi^2}{3}(2P_\eta^2 + P_{\eta'}^2)s^\mp.
\end{aligned} \tag{15}$$

The sea quark spin distribution functions can be calculated in χCQM by substituting for every constituent quark

$$q^\pm \to P(q \to \mathrm{GB})q^\pm + P(q^\pm \to q\bar{q}' + q'^\mp). \tag{16}$$

The spin structure for the spin $-\frac{3}{2}^+$ decuplet baryons can be expressed in terms of the probabilities of the emission of a GB from any of the quark and the transition probability of the emission of any valence quark to all the possible other quarks along with GBs. We have

$$\begin{aligned}
\Delta^{++}(uuu) &= 3P(u \to \mathrm{GB})u^\pm + 3P(u^\pm \to u\bar{q}' + q'^\mp),\\
\Delta^{+}(uud) &= 2P(u \to \mathrm{GB})u^\pm + P(d \to \mathrm{GB})d^\pm + 2P(u^\pm \to u\bar{q}' + q'^\mp) + P(d^\pm \to d\bar{q}' + q'^\mp),\\
\Delta^{0}(udd) &= P(u \to \mathrm{GB})u + 2P(d \to \mathrm{GB})d + P(u \to u\bar{q}' + q'^\mp) + 2P(d \to d\bar{q}' + q'^\mp),\\
\Delta^{-}(ddd) &= 3P(d \to \mathrm{GB})d^\pm + 3P(d^\pm \to d\bar{q}' + q'^\mp),\\
\Sigma^{*+}(uus) &= 2P(u \to \mathrm{GB})u^\pm + P(s \to \mathrm{GB})s^\pm + 2P(u^\pm \to u\bar{q}' + q'^\mp) + P(s^\pm \to s\bar{q}' + q'^\mp),\\
\Sigma^{*0}(uds) &= P(u \to \mathrm{GB})u^\pm + P(d \to \mathrm{GB})d^\pm + P(s \to \mathrm{GB})s^\pm\\
&\quad + P(u^\pm \to u\bar{q}' + q'^\mp) + P(d^\pm \to d\bar{q}' + q'^\mp) + P(s^\pm \to s\bar{q}' + q'^\mp),\\
\Sigma^{*-}(dds) &= 2P(d \to \mathrm{GB})d^\pm + P(s \to \mathrm{GB})s^\pm + 2P(d^\pm \to d\bar{q}' + q'^\mp) + P(s^\pm \to s\bar{q}' + q'^\mp),\\
\Xi^{*0}(uss) &= P(u \to \mathrm{GB})u^\pm + 2P(s \to \mathrm{GB})s^\pm + P(u^\pm \to u\bar{q}' + q'^\mp) + 2P(s^\pm \to s\bar{q}' + q'^\mp),\\
\Xi^{*-}(dss) &= P(d \to \mathrm{GB})d^\pm + 2P(s \to \mathrm{GB})s^\pm + P(d^\pm \to d\bar{q}' + q'^\mp) + 2P(s^\pm \to s\bar{q}' + q'^\mp),\\
\Omega^{-}(sss) &= 3P(s \to \mathrm{GB})s^\pm + 3P(s^\pm \to s\bar{q}' + q'^\mp).
\end{aligned} \tag{17}$$

A. Quark spin polarizations

The explicit quark spin distributions of the spin $-\frac{3}{2}^+$ decuplet baryons can be evaluated using the matrix elements for the spin structure, which are, in general, defined as follows [38]

$$\widehat{B^{\frac{3}{2}^+}} \equiv \langle B^{*\frac{3}{2}^+} | \mathcal{N}_{q^+q^-} | B^{*\frac{3}{2}^+} \rangle, \tag{18}$$

where $|B^{*\frac{3}{2}^+}\rangle$ is the SU(6) wave function giving the spin and flavor structure of the baryon (detailed in Ref. [47]), and $\mathcal{N}_{q^+q^-}$ is the number operator calculating the explicit quarks with polarization up or down,

$$\mathcal{N}_{q^+q^-} = n_{u^+}u^+ + n_{u^-}u^- + n_{d^+}d^+ + n_{d^-}d^- + n_{s^+}s^+ + n_{s^-}s^-. \tag{19}$$

In the above equation, the number of the quarks polarized in the up (down) direction are given as $n_{q^+}(n_{q^-})$, respectively, for each light quark u, d, and s.

The wave function for the ground state of decuplet baryons can be expressed in terms of the appropriate flavor and spin parts using the symmetry principles. We have

$$|B^{*\frac{3}{2}^+}\rangle \equiv \left|10, \frac{3^+}{2}\right\rangle = \varphi^s \chi^s, \tag{20}$$

where the spin wave functions (χ^s) for the case of spin $-\frac{3}{2}^+$ decuplet baryons are expressed as

$$\chi^s = \uparrow\uparrow\uparrow. \tag{21}$$

The flavor wave functions φ^s for the spin $-\frac{3}{2}^+$ decuplet baryons of the types $B^{\frac{3}{2}^+}(q_1q_1q_1)$, $B^{\frac{3}{2}^+}(q_1q_1q_2)$, and $B^{\frac{3}{2}^+}(q_1q_2q_3)$ are, respectively, expressed as

$$\begin{aligned}
\phi^s_{B^{\frac{3}{2}^+}} &= q_1q_1q_1, \\
\phi^s_{B^{\frac{3}{2}^+}} &= \frac{1}{\sqrt{3}}(q_1q_1q_2 + q_1q_2q_1 + q_2q_1q_1), \\
\phi^s_{B^{\frac{3}{2}^+}} &= \frac{1}{\sqrt{6}}(q_1q_2q_3 + q_1q_3q_2 + q_2q_1q_3 \\
&\quad + q_2q_3q_1 + q_3q_1q_2 + q_3q_2q_1).
\end{aligned} \tag{22}$$

The quark spin polarizations are basically defined as the difference of the quark polarized in the up direction and the quark polarized in the down direction. For the spin $-\frac{3}{2}^+$ decuplet baryons, we have

$$\Delta q_{B^{\frac{3}{2}}} = q^+_{B^{\frac{3}{2}}} - q^-_{B^{\frac{3}{2}}}. \tag{23}$$

Using Eqs. (18) and (20), we can calculate for any member of the spin $-\frac{3}{2}^+$ decuplet baryon, the number of q quarks polarized in up (down) direction [38,39]

$$q^{+(-)} = \langle \varphi^s\chi^s | n_{u^+}(n_{u^-}) | \varphi^s\chi^s \rangle. \tag{24}$$

For the spin $-\frac{3}{2}^+$ decuplet baryons of the types $B^{\frac{3}{2}^+}(q_1q_1q_1)$, $B^{\frac{3}{2}^+}(q_1q_1q_2)$, and $B^{\frac{3}{2}^+}(q_1q_2q_3)$, we, respectively, have

$$q_1^+ = 3, \qquad q_1^- = 0, \tag{25}$$

$$q_1^+ = 2, \quad q_1^- = 0, \quad q_2^+ = 1, \quad q_2^- = 0, \tag{26}$$

$$\begin{aligned}
&q_1^+ = 1, \qquad q_1^- = 0, \qquad q_2^+ = 1, \\
&q_2^- = 0, \qquad q_3^+ = 1, \qquad q_3^- = 0.
\end{aligned} \tag{27}$$

The constituent quark spin polarizations can be calculated from the above equations, whereas the sea quark spin polarizations can be calculated using Eqs. (16) and (17) by substituting for every constituent quark. The quark spin polarizations for spin $-\frac{3}{2}$ decuplet baryons have been presented in Table I. Even though the expressions for spin polarizations have already been derived and presented in Ref. [39], we present them here for the sake of completeness since the numerical values for the axial-vector form factors will be calculated from them. A closer look at the expressions of these quantities reveals that the constant factors represent the constituent quark polarizations. On the other hand, the factors with transition probability P_π^2 represent the contribution from the quark sea in general (with or without SU(3) symmetry breaking). These expressions are calculated in line with the methodology adopted in Ref. [39] and are in agreement with their results if the same values of parameters are taken.

B. Axial-vector form factors at $Q^2 = 0$

As discussed earlier, the axial-vector form factors at zero momentum transfer can be related to the different combinations of explicit quark spin polarizations. The singlet, triplet, and octet axial-vector form factors can be defined using the spin polarizations of the constituent quark calculated for each decuplet baryon using the chiral constituent quark model. We have

$$\begin{aligned}
G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0) &= \Delta u_{B^{*\frac{3}{2}}} + \Delta d_{B^{*\frac{3}{2}}} + \Delta s_{B^{*\frac{3}{2}}}, \\
G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(0) &= \Delta u_{B^{*\frac{3}{2}}} - \Delta d_{B^{*\frac{3}{2}}}, \\
G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(0) &= \Delta u_{B^{*\frac{3}{2}}} + \Delta d_{B^{*\frac{3}{2}}} - 2\Delta s_{B^{*\frac{3}{2}}}.
\end{aligned} \tag{28}$$

In order to numerically calculate the above form factors, we have to first fix the symmetry breaking parameters of χCQM. We have already defined the transition probabilities P_π^2, P_K^2, P_η^2, and $P_{\eta'}^2$, which give the extent of quark sea contribution to the structure of the baryon. In the present work, we will use the set of parameters as obtained by carrying out a best fit analysis for the case of spin and flavor distribution functions of spin $-\frac{1}{2}^+$ octet baryons since no experimental data is available for the case of spin $-\frac{3}{2}^+$ decuplet baryons. To find the upper and lower limits of the parameters, a very gross analysis was carried out using the quantities whose experimental measurements are

TABLE I. Quark spin polarizations for the spin $-\frac{3}{2}^+$ decuplet baryons in terms of the constituent quark polarizations (constant factors) and quark sea polarizations (with transition probability P_π^2).

Decuplet baryons $B^{*\frac{3}{2}^+}$	$\Delta u_{B^{\frac{3}{2}}}$	$\Delta d_{B^{\frac{3}{2}}}$	$\Delta s_{B^{\frac{3}{2}}}$
$\Delta^{++}(uuu)$	$3-P_\pi^2(6+3P_K^2+P_\eta^2+2P_{\eta'}^2)$	$-3P_\pi^2$	$-3P_\pi^2P_K^2$
$\Delta^{+}(uud)$	$2-P_\pi^2(5+2P_K^2+\frac{2}{3}P_\eta^2+\frac{4}{3}P_{\eta'}^2)$	$1-P_\pi^2(4+P_K^2+\frac{1}{3}P_\eta^2+\frac{2}{3}P_{\eta'}^2)$	$-3P_\pi^2P_K^2$
$\Delta^{o}(udd)$	$1-P_\pi^2(4+P_K^2+\frac{1}{3}P_\eta^2+\frac{2}{3}P_{\eta'}^2)$	$2-P_\pi^2(5+2P_K^2+\frac{2}{3}P_\eta^2+\frac{4}{3}P_{\eta'}^2)$	$-3P_\pi^2P_K^2$
$\Delta^{-}(ddd)$	$-3P_\pi^2$	$3-P_\pi^2(6+3P_K^2+P_\eta^2+2P_{\eta'}^2)$	$-3P_\pi^2P_K^2$
$\Sigma^{*^+}(uus)$	$2-P_\pi^2(4+3P_K^2+\frac{2}{3}P_\eta^2+\frac{4}{3}P_{\eta'}^2)$	$-P_\pi^2(P_K^2+2)$	$1-2P_\pi^2(2P_K^2+\frac{2}{3}P_\eta^2+\frac{1}{3}P_{\eta'}^2)$
$\Sigma^{*^o}(uds)$	$1-P_\pi^2(3+2P_K^2+\frac{1}{3}P_\eta^2+\frac{2}{3}P_{\eta'}^2)$	$1-P_\pi^2(3+2P_K^2+\frac{1}{3}P_\eta^2+\frac{2}{3}P_{\eta'}^2)$	$1-2P_\pi^2(2P_K^2+\frac{2}{3}P_\eta^2+\frac{1}{3}P_{\eta'}^2)$
$\Sigma^{*^-}(dds)$	$-P_\pi^2(P_K^2+2)$	$2-P_\pi^2(4+3P_K^2+\frac{2}{3}P_\eta^2+\frac{4}{3}P_{\eta'}^2)$	$1-2P_\pi^2(2P_K^2+\frac{2}{3}P_\eta^2+\frac{1}{3}P_{\eta'}^2)$
$\Xi^{*^o}(uss)$	$1-P_\pi^2(2+3P_K^2+\frac{1}{3}P_\eta^2+\frac{2}{3}P_{\eta'}^2)$	$-P_\pi^2(2P_K^2+1)$	$2-P_\pi^2(5P_K^2+\frac{8}{3}P_\eta^2+\frac{4}{3}P_{\eta'}^2)$
$\Xi^{*^-}(dss)$	$-P_\pi^2(2P_K^2+1)$	$1-P_\pi^2(2+3P_K^2+\frac{1}{3}P_\eta^2+\frac{2}{3}P_{\eta'}^2)$	$2-P_\pi^2(5P_K^2+\frac{8}{3}P_\eta^2+\frac{4}{3}P_{\eta'}^2)$
$\Omega^{-}(sss)$	$-3P_\pi^2P_K^2$	$-3P_\pi^2P_K^2$	$3-6P_\pi^2(P_K^2+\frac{2}{3}P_\eta^2+\frac{1}{3}P_{\eta'}^2)$

TABLE II. Input parameters of the χCQM transition probability parameters used in the analysis of axial-vector form factors.

Parameter →	P_π^2	P_K^2	P_η^2	$P_{\eta'}^2$	M_A
Set 1	0.114	0.202	0.202	0.562	$1.10^{+0.13}_{-0.15}$ GeV
Set 2	0.220	0.202	0.202	0.562	$1.10^{+0.13}_{-0.15}$ GeV

available [41,44]. Following these physical considerations, a through and refined analysis was carried out. The input parameters have been summarized in Table II (set 1). In order to understand the quantitative contribution of the chiral fluctuation, we have taken another set of parameters (set 2), wherein we have varied the transition probability P_π^2.

Using these set of parameters, the quark spin polarizations and the axial-vector form factors for the decuplet baryons at zero momentum transfer have been presented in Table III. After a cursory look at the table, we obtain the following relations for the case of flavor singlet axial-vector form factors for Δ, Σ^*, Ξ^*, and Ω decuplet baryons

$$\begin{aligned} &G^0_{\mathrm{AV},\Delta^{++}}(0)=G^0_{\mathrm{AV},\Delta^{+}}(0)=G^0_{\mathrm{AV},\Delta^{0}}(0)=G^0_{\mathrm{AV},\Delta^{-}}(0),\\ &G^0_{\mathrm{AV},\Sigma^{*+}}(0)=G^0_{\mathrm{AV},\Sigma^{*0}}(0)=G^0_{\mathrm{AV},\Sigma^{*-}}(0),\\ &G^0_{\mathrm{AV},\Xi^{*0}}(0)=G^0_{\mathrm{AV},\Xi^{*-}}(0). \end{aligned} \tag{29}$$

The flavor hypercharge (octet) axial-vector form factors $G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$ also obey the same relations for each multiplet of the decuplet baryons. The flavor isovector (triplet) axial-vector form factors, however, are different for each baryon in the multiplet. We have

$$\begin{aligned} &G^3_{\mathrm{AV},\Delta^{++}}(0)\neq G^3_{\mathrm{AV},\Delta^{+}}(0)\neq G^3_{\mathrm{AV},\Delta^{0}}(0)\neq G^3_{\mathrm{AV},\Delta^{-}}(0),\\ &G^3_{\mathrm{AV},\Sigma^{*+}}(0)\neq G^3_{\mathrm{AV},\Sigma^{*0}}(0)\neq G^3_{\mathrm{AV},\Sigma^{*-}}(0),\\ &G^3_{\mathrm{AV},\Xi^{*0}}(0)\neq G^3_{\mathrm{AV},\Xi^{*-}}(0). \end{aligned} \tag{30}$$

In addition to these, there are some interesting relations between the flavor triplet axial-vector form factors of the baryons from different multiplets, which are important to mention here. We get

$$\begin{aligned} &G^3_{\mathrm{AV},\Delta^{+}}(0)=G^3_{\mathrm{AV},\Xi^{*0}}(0),\\ &G^3_{\mathrm{AV},\Delta^{0}}(0)=G^3_{\mathrm{AV},\Xi^{*-}}(0), \end{aligned} \tag{31}$$

Further, we have

$$G^3_{\mathrm{AV},\Sigma^{*0}}(0)=G^3_{\mathrm{AV},\Omega^{-}}(0)=0. \tag{32}$$

TABLE III. The χCQM results for the quark spin polarizations and the axial-vector form factors at $Q^2=0$ for the Δ, Σ^*, Ξ^*, and Ω decuplet baryons.

Quantity	Δ^{++}	Δ^{+}	Δ^{0}	Δ^{-}	Σ^{*^+}	Σ^{*^0}	Σ^{*^-}	Ξ^{*^0}	Ξ^{*^-}	Ω^-
$\Delta u_{B^{\frac{3}{2}}}$	2.095	1.283	0.470	−0.342	1.374	0.561	−0.251	0.652	−0.160	−0.069
$\Delta d_{B^{\frac{3}{2}}}$	−0.342	0.470	1.283	2.095	−0.251	0.561	1.374	−0.160	0.652	−0.069
$\Delta s_{B^{\frac{3}{2}}}$	−0.069	−0.069	−0.069	−0.069	0.834	0.834	0.834	1.738	1.738	2.641
$G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$	1.684	1.684	1.684	1.684	1.957	1.957	1.957	2.230	2.230	2.503
$G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$	2.437	0.812	−0.812	−2.437	1.625	0.000	−1.625	0.812	−0.812	0.000
$G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$	1.891	1.891	1.891	1.891	−0.545	−0.545	−0.545	−2.983	−2.983	−5.421

These observations, along with isospin symmetry, can be clearly understood from Table I. It is clear that the explicit quark spin polarizations of each baryon can be related to its isospin partner with the interchange of $u \longleftrightarrow d$, $d \longleftrightarrow s$, or $s \longleftrightarrow u$. This is the reason we get

$$\begin{aligned} \Delta u_{\Delta^{++}} &= \Delta u_{\Delta^-}, \\ \Delta u_{\Delta^+} &= \Delta u_{\Delta^0}, \\ \Delta u_{\Sigma^{*+}} &= \Delta u_{\Sigma^{*-}}, \\ \Delta u_{\Xi^{*0}} &= \Delta u_{\Xi^{*-}}. \end{aligned} \tag{33}$$

We also have

$$\Delta s_{\Delta^{++}} = \Delta u_{\Delta^+} = \Delta u_{\Delta^0} = \Delta u_{\Delta^-}. \tag{34}$$

This has important implication toward the contribution of quark sea in the nonstrange baryons. The strange spin polarization is coming entirely from the quark sea, as there are no s quarks in the Δ baryons. Further, for the case of Σ^* and Ξ^* having $(q_1 q_2 s)$ and $(q_1 s s)$ quark content, respectively, we have

$$\begin{aligned} \Delta s_{\Sigma^{*+}} &= \Delta s_{\Sigma^{*0}} = \Delta s_{\Sigma^{*-}}, \\ \Delta s_{\Xi^{*0}} &= \Delta s_{\Xi^{*-}}. \end{aligned} \tag{35}$$

The singlet axial-vector form factor as defined in Eq. (28) gives the sum of explicit spin polarizations $\Delta u_{B^{*\frac{3}{2}}} + \Delta d_{B^{*\frac{3}{2}}} + \Delta s_{B^{*\frac{3}{2}}}$. If we just consider only the constituent quarks of baryons, we should get

$$G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0) = 3. \tag{36}$$

This combination is further related to the total spin of the baryon as

$$\frac{1}{2}\Delta\Sigma = \frac{1}{2} G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0) = \frac{3}{2}. \tag{37}$$

If we see the results from Table III, the results for singlet axial-vector form factor for the case of all the multiplets of baryons are much lower than the constituent quark model results. This is completely in line with the observations from the DIS experiments characterized by proton spin crisis, which have provided the strong evidence for the constituent quarks carrying only about 30% of the total spin for the case of proton. It is possible to describe the missing spin through the angular momentum conservation, where this reduced spin due to the quark sea is compensated by orbital angular momentum carried by the sea [46,48,49]. Even though there is no possibility of any experimental studies for the case of the decuplet baryons, a deep understanding of the dynamics of the constituent quarks will form a basis for formulating any model in the nonperturbative regime.

For the case of flavor isovector (triplet) axial-vector form factors, the role of the quark sea is clearly evident when we compare the results with the constituent quark model results [50–52], where only the constituent quarks contribute. We have

$$\begin{aligned} G^3_{\mathrm{AV},\Delta^{++}}(0) &= -G^3_{\mathrm{AV},\Delta^-}(0) = 3, \\ G^3_{\mathrm{AV},\Delta^+}(0) &= -G^3_{\mathrm{AV},\Delta^0}(0) = 1, \\ G^3_{\mathrm{AV},\Sigma^{*+}}(0) &= -G^3_{\mathrm{AV},\Sigma^{*-}}(0) = 2, \\ G^3_{\mathrm{AV},\Xi^{*0}}(0) &= -G^3_{\mathrm{AV},\Xi^{*-}}(0) = 1, \\ G^3_{\mathrm{AV},\Sigma^{*0}}(0) &= G^3_{\mathrm{AV},\Omega^-}(0) = 0. \end{aligned} \tag{38}$$

When we compare these results with the χCQM results, we find that because of the quark sea contributing with an opposite magnitude, the results of χCQM are toward the lower side. If we compare it with the case of octet baryons, the results of experiments are much lower than the constituent quark model results for the case of nucleon, and the χCQM results are more or less in agreement with the data [53], making us conclude that the quark sea contribution will play a vital role in the case of decuplet baryons also. The isovector triplet axial-vector form factor is one of the most well-known factors, and it connects the axial charge and the strong coupling constant through the Goldberger-Treiman relation.

The flavor singlet axial-vector form factor and the flavor hypercharge (octet) axial-vector form factor reduce to the Ellis-Jaffe sum rule [54,55] in the limit of $\Delta s = 0$, as evident from Eq. (28). We have

$$G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0) = G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(0)|_{\Delta s=0}. \tag{39}$$

Since the strange quarks and consequently strange spin polarization are the same for each baryon in a particular multiplet, we get the same hypercharge (octet) axial-vector for the case of Δ, Σ^*, Ξ^*, and Ω decuplet baryons.

C. Q^2 dependence in axial-vector form factors

In order to understand the role of Q^2 dependence of the axial-vector form factors, the form factors are often parametrized in terms of dipole type or parametrization or a p–pole type of parametrization. Experimentally, the Q^2 dependence is investigated from the quasielastic neutrino scattering [6,7], as well as from the pion electroproduction [8]. In the present work, to analyze the axial-vector form factors, we consider the most conventionally used dipole form of parametrization, which is given by

$$G^a_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) = \frac{G^a_{\mathrm{AV},B^{*\frac{3}{2}}}(0)}{\left(1+\frac{Q^2}{M_A^2}\right)^2}, \tag{40}$$

where $G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$, $G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$, and $G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$ are the isovector axial-vector coupling constants at zero momentum transfer. Here, M_A is defined as the axial mass, which can be related to the axial radius through this parametrization. As extracted from neutrino scattering experiments, we have the global average of M_A as (1.026 ± 0.021) GeV [56]. A slightly smaller value can also be found at $M_A = (1.001 \pm 0.020)$ GeV [57]. Axial mass can also be fitted to the experiment and taken as free parameter [58]. At present, we have the experimental data only for the case of nucleon axial-vector form factors, but, by convention, the axial masses corresponding to Σ^*, Ξ^*, and Ω are expected to be larger. Even though the large value of axial mass will lead to larger values of the axial-vector form factors in magnitude, the overall behavior of the form factors, however, remains the same. We have considered the most recent value from the MiniBooNE Collaboration $M_A = 1.10^{+0.13}_{-0.15}$ GeV [59] for studying the Q^2 dependence in the present work.

The behavior and magnitude of the form factors for decuplet baryons mainly depends on the constituent quark structure of the baryon. Therefore, it would be interesting to mention here that the axial-vector form factors can be defined for the constituent quarks, which are spatially extended particles [58]. The explicit quark flavor axial-vector form factors can be defined in terms of singlet, triplet, and octet axial-vector form factors, and further Q^2 dependence can be studied using the dipole form of parametrization [Eq. (40)]. We have

$$\begin{aligned}
G^u_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) &= \frac{1}{3}G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) + \frac{1}{2}G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) \\
&\quad + \frac{1}{2\sqrt{3}}G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2), \\
G^d_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) &= \frac{1}{3}G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) - \frac{1}{2}G^3_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) \\
&\quad + \frac{1}{2\sqrt{3}}G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2), \\
G^s_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) &= \frac{1}{3}G^0_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2) - \frac{1}{\sqrt{3}}G^8_{\mathrm{AV},B^{*\frac{3}{2}}}(Q^2).
\end{aligned} \tag{41}$$

The explicit values of quark flavor axial-vector form factors at zero momentum transfer have been presented in Table IV using set 1 from Table II. A cursory look at the table reveals the fact that the constituent quarks clearly dominate over the quark sea, and the sea quarks contribute in the opposite direction. Experimental measurements for heavier baryons can perhaps substantiate this fact further.

TABLE IV. The χCQM results for the quark flavor axial-vector form factors at $Q^2 = 0$ for the Δ, Σ^*, Ξ^*, and Ω decuplet baryons.

Baryon	$G^u_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$	$G^d_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$	$G^s_{\mathrm{AV},B^{*\frac{3}{2}}}(0)$
Δ^{++}	2.326	−0.111	−0.530
Δ^{+}	1.513	0.701	−0.530
Δ^{0}	0.701	1.513	−0.530
Δ^{-}	−0.111	2.326	−0.530
Σ^{*^+}	1.307	−0.317	0.967
Σ^{*^0}	0.494	0.494	0.967
Σ^{*^-}	−0.317	1.307	0.967
Ξ^{*^0}	0.288	−0.524	2.466
Ξ^{*^-}	−0.524	0.288	2.466
Ω^-	−0.730	−0.730	3.964

In order to discuss the dependence of the constituent quark form factors on Q^2 $(0 \le Q^2 \le 1)$, in Fig. 1, we have plotted the explicit u, d, and s quark flavor axial-vector form factors for each of the Δ, Σ^*, Ξ^*, and Ω decuplet baryons.

The constituent quark structure is clearly reflected in the plots. For example, in the case of $\Delta^{++}(uuu)$, the form factor $G^u_{\mathrm{AV},B^{*\frac{3}{2}}}$ dominates. The interesting point in this case is the presence of the $G^s_{\mathrm{AV},B^{*\frac{3}{2}}}$. Even though it is small, it clearly reflects the importance of the quark sea in understanding its underlying dynamics. Similarly, in the case of $\Sigma^{*^+}(uus)$, it is very clear that the form factors $G^u_{\mathrm{AV},B^{*\frac{3}{2}}}$ and $G^s_{\mathrm{AV},B^{*\frac{3}{2}}}$ dominate. Here, also, the presence of $G^d_{\mathrm{AV},B^{*\frac{3}{2}}}$ is due to the presence of quark sea. Further, in the case of $\Xi^{*^0}(uss)$, the form factor $G^s_{\mathrm{AV},B^{*\frac{3}{2}}}$ dominates, and for $\Omega^-(sss)$, the quark sea contributes equally to the u and d form factors, and $G^s_{\mathrm{AV},B^{*\frac{3}{2}}}$ has maximum contribution. Another important observation in the plots is the dominance of the role of quark sea at low Q^2. As we move toward higher Q^2, the contribution is still dominated by the constituent quarks, but the constituent quark form factors fall off rapidly with increasing Q^2.

Further, to discuss the Q^2 dependence $(0 \le Q^2 \le 1)$ in the axial-vector form factors, we present in Fig. 2, singlet, triplet, and octet axial-vector form factors for Δ, Σ^*, Ξ^*, and Ω decuplet baryons. From the plots, one can easily discuss the variation and sensitivity to Q^2 for the form factors. Broadly speaking, the singlet, triplet, and octet vary with Q^2 in the following order

$$\begin{aligned}
&G^0_{\mathrm{AV},\Omega} > G^0_{\mathrm{AV},\Xi^{*-}} > G^0_{\mathrm{AV},\Sigma^{*+}} > G^0_{\mathrm{AV},\Delta^+}, \\
&G^3_{\mathrm{AV},\Delta^{++}} > G^3_{\mathrm{AV},\Sigma^{*+}} > G^3_{\mathrm{AV},\Delta^+} > G^3_{\mathrm{AV},\Xi^{*-}}, \\
&G^8_{\mathrm{AV},\Delta^+} > G^8_{\mathrm{AV},\Sigma^{*+}} > G^8_{\mathrm{AV},\Xi^{*-}} > G^8_{\mathrm{AV},\Omega}.
\end{aligned} \tag{42}$$

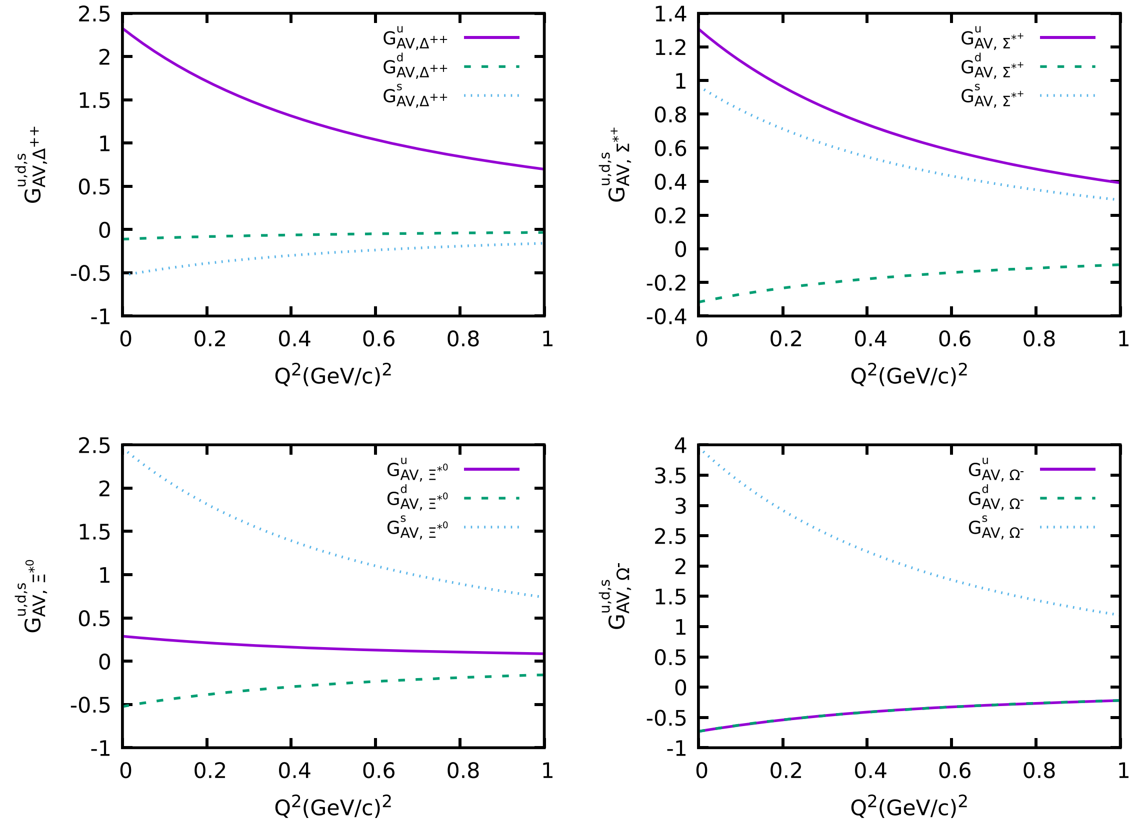

FIG. 1. The explicit quark flavor axial-vector form factors for Δ, Σ^*, Ξ^*, and Ω decuplet baryons plotted as function of Q^2.

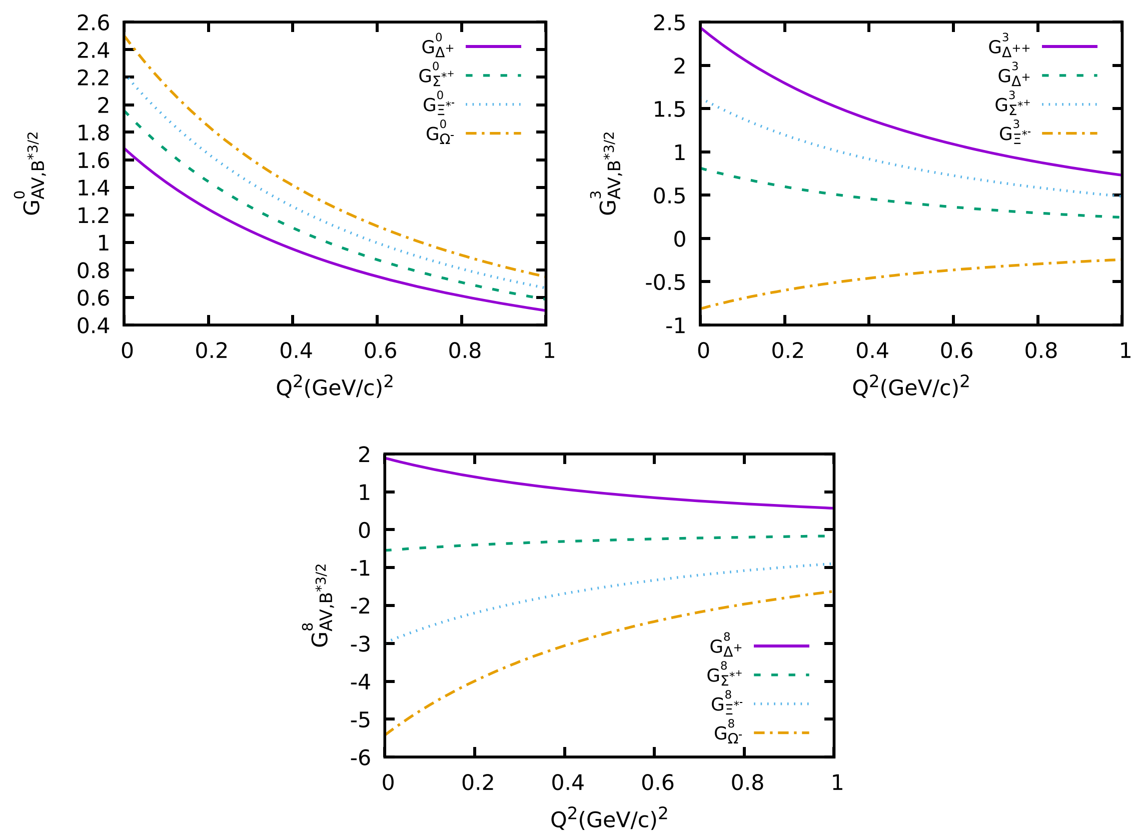

FIG. 2. Form factors for Δ, Σ^*, Ξ^*, and Ω decuplet baryons plotted as function of Q^2.

The singlet axial-vector form factors $G^0_{\mathrm{AV},B^{*\frac{3}{2}}}$ fall off rapidly for all the decuplet baryons. The triplet axial-vector form factors $G^3_{\mathrm{AV},B^{*\frac{3}{2}}}$ fall off rapidly for the cases where u quark dominates ($G^3_{\mathrm{AV},\Delta^{++}}$ and $G^3_{\mathrm{AV},\Sigma^{*+}}$), whereas for other cases, it falls off slowly. The octet axial-vector form factors $G^8_{\mathrm{AV},B^{*\frac{3}{2}}}$, fall off rapidly for the cases having dominant s quarks ($G^8_{\mathrm{AV},\Omega}$ and $G^8_{\mathrm{AV},\Xi^{*-}}$).

In order to understand the role of chiral fluctuations in terms of the transition probabilities, we consider a different set of input parameters (set 2 from Table II). We consider here the variation of only the transition probability P_π^2, where the strange and nonstrange quark masses are considered to be nondegenerate $M_s > M_{u,d}$. Since the hierarchy of the transition probabilities is based on the scaling of the quark contributions as $\frac{1}{M_q^2}$, it is clearly evident from Eq. (9) that P_π^2 quantitatively dominates over the other transition probabilities to understand the extent of quark sea contribution in the baryon structure. Further, in order to understand in depth the role of the transition parameters for each of the flavor singlet, flavor isovector (triplet), and the flavor hypercharge axial (octet) axial-vector form factors, we present in Fig. 3 the case of $G^0_{\mathrm{AV},\Delta^+}(Q^2)$, $G^3_{\mathrm{AV},\Delta^+}(Q^2)$, and $G^8_{\mathrm{AV},\Delta^+}(Q^2)$ for the two set of parameters given in Table II. From the figure, we find that the singlet axial-vector form factor, which gives the sum of explicit spin polarizations, is highly sensitive to the parameter P_π^2. This is again in line with the experimental observations for the case of nucleons, which indicate that the quark sea contributes to the baryon spin with a reversed sign, or, in other words, it reduces the total contribution of the constituent quarks. The lowering of $G^0_{\mathrm{AV},\Delta^+}(Q^2)$ at $Q^2 = 0$ for $P_\pi^2 = 0.220$ as compared to $P_\pi^2 = 0.114$ endorses the fact that more of the probability of constituent quark fluctuating to a sea of quarks is more so the contribution of the quark sea in the baryon spin. Future experimental measurements for decuplet baryons will fix the transition probabilities so that the constituent and sea quark distributions add up in the right direction to give an excellent overall fit to the axial-vector form factors. Further, flavor triplet axial-vector form factor $G^3_{\mathrm{AV},\Delta^+}(Q^2)$ is also important. In this case, as evident from Eq. (28), only the u and d spin polarizations contribute. As Δ^+ has uud as the quark content, so changing the parameter P_π^2 does not affect it to a large extent. In this case, the strange quarks contribute only through the parameters P_K^2, P_η^2, and $P_{\eta'}^2$, which have values of second order as compared to the value of P_π^2. This reduces to the fact that this term is dominated by the constituent quark structure. Any refinement for the

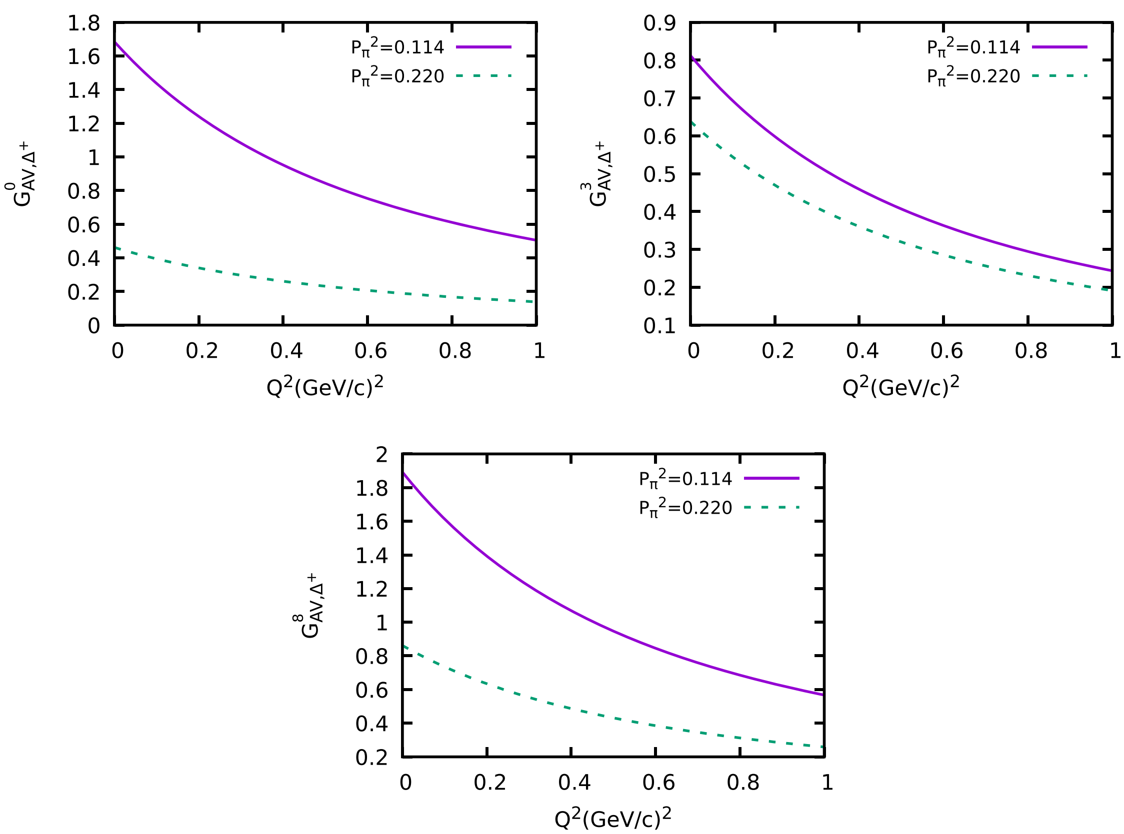

FIG. 3. Axial-vector form factors for the case of Δ^+ plotted as function of Q^2 for two set of parameters.

strangeness dependent quantities would have important implications for the basic features of χCQM. As for the case of $G^0_{\mathrm{AV},\Delta^+}(Q^2)$, the flavor hypercharge axial (octet) axial-vector form factor $G^8_{\mathrm{AV},\Delta^+}(Q^2)$ also is sensitive to the transition probability P_π^2.

IV. SUMMARY AND CONCLUSIONS

To summarize, the axial-vector form factors of the spin $-\frac{3}{2}^+$ decuplet baryons have been investigated in the chiral constituent quark model (χCQM) using the explicit quark spin polarizations, which have important implications for the quark sea arising from the chiral symmetry breaking. The Goldstone bosons mediating the interaction between constituent quarks play an important role to phenomenologically estimate the extent of quark sea contribution in understanding the internal structure of the baryons. Even though this model does not solve any dynamical equations, the symmetry and phenomenological parametrization not only make it easier to handle, but they are also able to explain the underlying dynamics of quarks inside the baryons. In particular, we have studied the axial-vector form factors corresponding to the flavor singlet current, flavor isovector (triplet) current, and the flavor hypercharge axial (octet) current at zero momentum transfer, which are, respectively, represented as $G^0_{\mathrm{AV},B^*\frac{3}{2}}(0)$, $G^3_{\mathrm{AV},B^*\frac{3}{2}}(0)$, and $G^8_{\mathrm{AV},B^*\frac{3}{2}}(0)$. The form factors have been investigated for the case of Δ, Σ^*, Ξ^*, and Ω spin $-\frac{3}{2}^+$ decuplet baryons. Some important relations for the cases for each baryon in the multiplet and for baryons in different multiplets have been discussed. The singlet axial-vector form factor is related to the total spin of the baryon, and the present results are completely in line with the observations from the DIS experiments. The quark sea contributes to the total spin with opposite sign as that of the constituent quark spin, and this reduced spin is further compensated by its orbital angular momentum, which has the same sign as that of the constituent quarks. The isovector triplet axial-vector form factor connects the axial charge and the strong coupling constant through the Goldberger-Treiman relation. Further, the flavor singlet axial-vector form factor and the flavor hypercharge (octet) axial-vector form factor reduce to the Ellis-Jaffe sum rule in the limit of $\Delta s = 0$.

The role of Q^2 dependence of the axial-vector form factors has been parametrized using the most conventionally used dipole form of parametrization. The constituent quark structure is clearly reflected in the study, and the dominance of the role of quark sea at low Q^2 is also clear. As we move toward higher Q^2, the contribution is still dominated by the constituent quarks; however, the dominant constituent quark form factors fall off rapidly. High precision measurements over a wide Q^2 region in the near future will impose important constraints on the parity-violating asymmetries in different kinematical regions.

The quantitative contribution of the quark sea has also been investigated by varying the transition probability of the chiral fluctuation, and it is found that it's more probable that constituent quark fluctuating to a sea of quarks more is the reduction in the total contribution of the constituent quarks. The strange quarks do not contribute significantly as the parameters P_K^2, P_η^2, and $P_{\eta'}^2$ have values of second order as compared to the value of P_π^2. Any refinement for the strangeness dependent quantities would have important implications for the basic tenents of χCQM, as well as the important role of chiral symmetry breaking and SU(3) symmetry breaking in the nonperturbative regime of QCD, where the constituent quarks and the weakly interacting Goldstone bosons constitute the appropriate degrees of freedom at the leading order.

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