The semileptonic baryonic decay $D_s^+ \rightarrow p\bar{p}e^+\nu_e$

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The decay $D_s^+ \rightarrow p\bar{p}e^+\nu_e$ with a proton-antiproton pair in the final state is unique in the sense that it is the only semileptonic baryonic decay which is physically allowed in the charmed meson sector. Its measurement will test our basic knowledge on semileptonic $D_s^+$ decays and the low-energy $p\bar{p}$ interactions. Taking into account the major intermediate state contributions from $\eta, \eta', f_0(980)$ and $X(1835)$, we find that its branching fraction is at the level of $10^{-9} \sim 10^{-8}$. The location and the nature of $X(1835)$ state are crucial for the precise determination of the branching fraction. We wish to trigger a new round of a careful study with the upcoming more data in BESIII as well as the future super tau-charm factory.

Keywords: Semileptonic baryonic decay, $D_s$ transition form factor, low-energy $p\bar{p}$ interaction, $X(1835)$ meson

I. INTRODUCTION

A great deal of effort has been devoted to the semileptonic decay modes of $B$ mesons $^1$ due to the fact that the $B$ meson is heavy enough to allow a baryon-antibaryon pair production in the final state. Concerning the semileptonic decay involving a baryon-antibaryon pair, $B^- \rightarrow p\bar{p}e^-\nu_e \ (\ell = e, \mu)$ is the only measurement that has been done by the Belle Collaboration in 2014 $^2$. Its branching fraction was reported to be $(5.8^{+2.4}_{-2.2} \pm 0.9) \times 10^{-6}$ with the upper limit $9.6 \times 10^{-6}$ at the 90\% confidence level. In the charmed meson sector, $D_s^+ \rightarrow p\bar{p}$ is the only hadronic baryonic $D$ decay mode which is physically allowed. Its branching ratio is naively expected to be very small, of order $10^{-6}$, due to chiral suppression $^3$. Hence, the observation of this mode by CLEO with $B(D_s^+ \rightarrow p\bar{p}) = (1.30 \pm 0.36^{+0.12}_{-0.10}) \times 10^{-3}$ $^4$ is indeed a surprise. Nevertheless, it can be explained by the final-state rescattering of $\pi^+\eta(\pi)$ and $K^+K^0$ into $p\bar{p}$ $^5$. Besides the channel $D_s^+ \rightarrow p\bar{p}$, we notice that there is another physically allowed one, $D_s^+ \rightarrow p\bar{p}e^+\nu_e$. The mass difference $m_{D_s^+} - 2m_p \approx 82$ MeV prohibits the emission of $\pi^+$ or even the lepton $\mu^+$, thus only the electron mode is permissible. Moreover, the $p\bar{p}$ pair stays in the near-threshold region, i.e., the invariant mass squared $s = (p_p+p_\bar{p})^2$ is not far from $4m_p^2$. The future experimental measurement can rectify the description of $D_s^+ \rightarrow p\bar{p}$ hadronic transition form factors as well as the low-energy $p\bar{p}$ interaction. If this channel can be observed, it renders a preponderant possibility to access the $p\bar{p}$ bound state due to the low-energy $p\bar{p}$ region.

Below we will calculate the branching fraction of the decay channel $D_s^+ \rightarrow p\bar{p}e^+\nu_e$. We first consider the Cabibbo-favored decay $D_s^+ \rightarrow M(s\bar{s}) + e^+\nu_e$ with $M$ being the meson containing a sizable $s\bar{s}$ quark component. Since such meson decayin to a $p\bar{p}$ pair is an OZI suppressed process, we shall focus on the intermediate mesons $M$ with comparable amount of $q\bar{q} = \frac{1}{\sqrt{2}}(u\bar{u} + d\bar{d})$ and $s\bar{s}$ components in order to alleviate the OZI suppression. Combining the existing knowledge on the $D_s^+ \rightarrow M$ transition form factors and $Mp\bar{p}$ couplings fixed by the $p\bar{p}$ scattering data, we are able to take into account the $\eta, \eta', f_0(980)$ and $X(1835)$ meson exchanges, and find that the branching fraction of $D_s^+ \rightarrow p\bar{p}e^+\nu_e$ is at the level of $10^{-9} \sim 10^{-8}$.

II. KINEMATICS AND DECAY RATE

The four-body decay kinematics can be described in terms of five variables: the invariant mass squared of the $p\bar{p}$ pair, $s = (p_p+p_\bar{p})^2 = M_{p\bar{p}}^2$, the invariant mass squared of the dilepton pair, $s_l = (p_e+p_\nu)^2$, the angles $\theta_p, \theta_l$ and $\phi$, where $\theta_p (\theta_l)$ is formed by the proton $p$ ($e^+$) direction in the diproton (dilepton) center-of-mass (CMS) frame with respect to the diproton (dilepton) line of flight in the $D_s^+$ frame, and $\phi$ is the dihedral angle between the diproton and diphoton planes. Their physical ranges are

\[
4m_p^2 \leq s \leq (m_{D_s^+} - m_{u\bar{d}})^2 , \quad m_{u\bar{d}}^2 \leq s_l \leq (m_{D_s^+} - \sqrt{s})^2 , \quad 0 \leq \theta_p, \theta_l \leq \pi, \quad 0 \leq \phi \leq 2\pi .
\]

One may refer to e.g., Ref. $^7$ for an illustration of the four-body decay kinematics.

Instead of the separate momenta $p_p, p_\bar{p}, p_e, p_\nu$, it is more convenient to use the following kinematic variables

\[
P = p_p + p_\bar{p}, \quad Q = p_p - p_\bar{p}, \quad L = p_e + p_\nu, \quad N = p_e - p_\nu.
\]

\[
(P - m_p)^2 = 4m_p^2, \quad Q^2 = 4m_p^2, \quad L^2 = 4m_p^2, \quad N^2 = 4m_p^2.
\]
It follows that
\[ P^2 = s, \quad Q^2 = 4m_p^2 - s, \quad L^2 = -N^2 = s_l, \]
\[ P \cdot L = \frac{1}{2}(m_{D^+_s}^2 - s - s_l) , \quad P \cdot N = X \cos \theta_l, \]  
(3)
where the function \( X \) is defined by
\[ X(s, s_l) = ((P \cdot L)^2 - s s_l)^{1/2} = \frac{1}{2} \lambda^{1/2}(m_{D^+_s}^2, s, s_l) , \]
\[ \lambda(x, y, z) = (x - y - z)^2 - 4yz. \]  
(4)
with \( m_p \) being the proton mass. The term \( P \cdot N \) can be derived by expressing the four momenta of \( p, \bar{p}, e^+, \nu_e \) in the rest frame of \( D^+_s \) via the Lorentz transformation, see e.g., [8]. Note that we have neglected the electron mass over most of the available phase space (although \( p\bar{p} \) sits in the low energy region), i.e., \( m_e^2 / s_l \ll 1 \). This has also been checked numerically.

The decay amplitude of \( D^+_s \to p\bar{p}e^+\nu_e \) can be written as
\[ T = \frac{G_F}{\sqrt{2}} V_{cd} l'_\mu h^\mu , \]
\[ l'_\mu = \bar{u}(p_\nu) \gamma_{\mu}(1 - \gamma^5) v(p_e) , \]
\[ h^\mu = \langle p\bar{p}|V^\mu - A^\mu|D^+_s \rangle , \]  
(5)
where the currents \( V^\mu \) and \( A^\mu \) denote the vector and axial-vector ones, respectively, and their hadronic matrix elements will be discussed in Sec. III. We then have the differential decay rate
\[ d^3\Gamma = \frac{1}{4(4\pi)^3 m_{D_s^+}^3} \sigma(s) X(s, s_l) \sum_{\text{spins}} |T|^2 
\times ds \, ds_l \, d \cos \theta_p \, d \cos \theta_l \, d\phi, \]  
(6)
with
\[ \sigma(s) = \sqrt{1 - 4m_p^2 / s} . \]  
(7)
The four-body phase space was studied very early in 1960s within the context of \( K\bar{K} \) analysis [9], see also Ref. [10] for a modern compilation. More details of derivation can be found in e.g., Refs. [11, 12]. Equation (6) is in agreement with Refs. [13, 14], as has been checked.

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### III. HADRONIC MATRIX ELEMENTS AND RESULTS

We begin with the hadronic matrix elements [13]
\[ \langle p\bar{p}|V^\mu - A^\mu|D_s^+ \rangle = -i\bar{u}(p_\nu) [g_1 \gamma^\mu + ig_2 \sigma_{\mu\nu} L^\nu + g_3 L^\mu + g_4 P^\mu + g_5 Q^\prime \gamma^5 \nu(p_\nu) , \]
\[ \langle p\bar{p}|A^\mu - M^\mu|D_s^+ \rangle = -i\bar{u}(p_\nu) [f_1 \gamma^\mu + if_2 \sigma_{\mu\nu} L^\nu + f_3 L^\mu + f_4 P^\mu + f_5 Q^\prime \gamma^5 \nu(p_\nu) . \]  
(8)
Note the spinors \( u \) and \( v \) have a relative opposite sign under parity transformation. Various form factors \( f_i \) and \( g_i \) will be evaluated below. As mentioned in the Introduction, to alleviate the OZI suppression for the intermediate meson exchange that leads to the decay \( D_s^+ \to p\bar{p}e^+\nu_e \), we shall focus on the intermediate states which have comparable \( q\bar{q} \) and \( s\bar{s} \) components. The two- and multi-meson exchanges are expected to be loop suppressed, and also the direct \( D_s^+ p\bar{p}W \) production vertex without any meson exchange can be safely neglected. We then concentrate on one-meson exchange denoted by \( M \). The combination of the existing knowledge of \( D_s^+ \to M \) transition and the coupling \( p\bar{p}M \) constitutes our basic strategy. In Ref. [16], we have explored the form factors and branching fractions for the semileptonic \( D_s^+ \to M \) transition. As for the \( p\bar{p}M \) part, we shall stick to the Jülich nucleon-antinucleon model [17, 18] which provides a fair description of \( p\bar{p} \) total, elastic, charge-exchange and annihilation cross sections. In such a \( p\bar{p} \) model, the exchanged mesons with mass up to 1.5 GeV were considered. We first include the spin-0 boson, \( \eta, \eta' \), \( f_0(980) \) in our study. The decay mechanism is shown in Fig. 1 where the upper panel describes the mechanism at the quark level with the bulk denoting the meson with the \( q\bar{q} \) and \( s\bar{s} \) components, and the lower one from the viewpoint of effective meson theory with the dashed line denoting the exchanged mesons.

We will calculate the Feynman diagram to single out the contributions of \( \eta, \eta' \), \( f_0(980) \) to the \( D_s^+ \to p\bar{p} \) transition form factors. We have
\[ \langle p\bar{p}|V^\mu - A^\mu|D_s^+ \rangle = \sum_M \langle M|V^\mu - A^\mu|D_s^+ \rangle \frac{i}{p^2 - m^2} V, \]  
(9)

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\(^1\) As a cross check, we may first keep the electron mass and retain the factor of \( z_l = m_e^2 / s_l \). Letting \( z_l \to 0 \), we then recover Eq. (22) below. The numerical results remain stable irrespective of the tiny electron mass.

\(^2\) The \( p\bar{p} \) interaction within the framework of chiral effective field theory involving pion degrees of freedom and contact terms was recently explored in Ref. [19] and Ref. [20]. A similar method has been recently applied to charmed baryon scattering [21].
which amounts to inserting the intermediate meson with the momentum $p$ and mass $m$, and $V$ is the vertex of $p\bar{p}M$ coupling. The $f_0(980)$ has a large width which may remind us of replacing $p^2 - m^2$ by $p^2 - m^2 - i m f_0 \Gamma_f_0$. However, it is not necessary to do so since the mass of $f_0(980)$ is still far from the $p\bar{p}$ invariant mass. The induced difference by including $\Gamma_f_0$ is only of order 0.1%. The $D_s^+ \rightarrow M$ transition can be described by

\[
\langle P | V^\mu | D_s^+ \rangle = \left( p^\mu - \frac{m_{D_s^+}^2 - m_P^2}{q^2} q^\mu \right) F_1^{D_s^+ \rightarrow P}(q^2) \\
+ \frac{m_{D_s^+}^2 - m_P^2}{q^2} q^\mu F_0^{D_s^+ \rightarrow P}(q^2),
\]

(10)

\[
\langle S | A^\mu | D_s^+ \rangle = -i \left[ \left( p^\mu - \frac{m_{D_s^+}^2 - m_S^2}{q^2} q^\mu \right) F_1^{D_s^+ \rightarrow S}(q^2) \\
+ \frac{m_{D_s^+}^2 - m_S^2}{q^2} q^\mu F_0^{D_s^+ \rightarrow S}(q^2) \right],
\]

(11)

where $P$ denotes the pseudoscalars $\eta$ and $\eta'$, $S$ the scalar $f_0(980)$, $p = p_{D_s^+} + p_p + p_{\bar{p}}$ and $q = p_{D_s^+} - p_p - p_{\bar{p}} = L$, thus $q^2 = s_L$. The form factors $F_1(q^2)$ and $F_0(q^2)$ for $D_s^+ \rightarrow \eta(q')$ have been investigated using the covariant light-front quark model [22, 24],

\[
F_1^{D_s^+ \rightarrow \eta}(q^2) = \frac{0.76}{1 - 1.02 \frac{q^2}{m_{D_s^+}^2} + 0.40 \left( \frac{q^2}{m_{D_s^+}^2} \right)^2},
\]

\[
F_0^{D_s^+ \rightarrow \eta}(q^2) = \frac{0.76}{1 - 0.60 \frac{q^2}{m_{D_s^+}^2} + 0.04 \left( \frac{q^2}{m_{D_s^+}^2} \right)^2},
\]

\[
F_i^{D_s^+ \rightarrow \eta}(q^2) = - \sin \phi \phi_i^{D_s^+ \rightarrow \eta}(q^2),
\]

\[
F_i^{D_s^+ \rightarrow \eta}(q^2) = \cos \phi \phi_i^{D_s^+ \rightarrow \eta}(q^2),
\]

(12)

for $i = 0$ or 1. $\phi$ is the mixing angle between $\eta$ and $\eta'$ defined by [27]

\[
|\eta\rangle = \cos \phi |\eta_q\rangle - \sin \phi |\eta_s\rangle,
\]

\[
|\eta'\rangle = \sin \phi |\eta_q\rangle + \cos \phi |\eta_s\rangle.
\]

(13)

It is determined to be 39.3° ± 1.0° in the Feldmann-Kroll-Stech mixing scheme [27], which is consistent with the recent result $\phi = 42° ± 2.8°$ extracted from the CLEO data [28]. In Ref. [16] it has been shown that such a description of form factors gives a rather good description of the branching fraction compared to experiment, and that replacing $m_{D_s^+}$ by $m_D$ in the denominator does not make significant difference for the result. As we have already commented in Ref. 16, the $D_s^+ \rightarrow f_0(980)$ transition form factor cannot be appropriately treated by the covariant light-front model since i) $f_0(980)$ is widely believed to be a tetraquark state (see e.g., [29]) or a $\bar{K}K$ molecular (see e.g. Refs. [30, 31]) rather than a pure quark-antiquark meson; and ii) the decay constant of $f_0(980)$ vanishes due to the charge conjugation invariance and thus there is no reliable constraint on the parameter in its wave function within the light-front quark model. However, the information of $F_1(q^2)$ for $D_s^+ \rightarrow f_0(980)$ is directly accessible by experiment, that is

\[
F_1^{D_s^+ \rightarrow f_0(980)}(q^2) = \frac{0.4}{1 - q^2/M_{\text{pole}}^2},
\]

(14)

with $M_{\text{pole}} = 1.7^{+0.5}_{-0.7}$ GeV from the CLEO Collaboration [32]. This situation is different from Refs. [33, 34], where only the form factor $F_0(q^2)$ enters in the factorization scheme of the two-body nonleptonic decay. In

\footnote{The value $F_1(0) = 0.4$ is not shown explicitly in Ref. 32, but can be obtained using the masses $m_{f_0}$, $M_{\text{pole}}$, and $M_{D_s^+} \rightarrow f_0(980)e^+\nu_e \approx 0.4\%$ reported there. The slope of $F_1(q^2)$, namely, $M_{\text{pole}}$, is fitted to the measured event distribution, which differs from the $d\Gamma/dq^2$ only by an overall constant, so $F_1(0)$ cannot be constrained by the event distribution and is left as a float in Ref. 33.}
the decay rate of the semileptonic decay for $D$ or $D_s^+$ to spin-0 boson, the form factor $F_0(q^2)$ is accompanied by the electron mass and thus negligible. In other words, $F_0(q^2)$ can be constrained by the corresponding nonleptonic decay rate based on factorization, but not from a direct experimental measurement. That is \[ F_0^{D_c \to f_0(980)}(q^2) = \frac{0.52}{1 - q^2/m_{D_c}^2} \] . (15)

For the part of the $p\bar{p}$ interaction, we have the Lagrangian \[ \mathcal{L}_P = g_P \bar{\psi}(x) i \gamma^5 \psi(x) \phi(x) \] (16) for the nucleon-nucleon-pseudoscalar (NNS) coupling, and \[ \mathcal{L}_S = g_S \bar{\psi}(x) \psi(x) \phi(x) \] (17) for the nucleon-nucleon-scalar coupling (NNS), with $\psi(x)$ denoting the nucleon field and $\phi(x)$ the meson field. The dimensionless couplings read \[ g_\eta = 2.87, \quad g_{\eta'} = 3.72, \quad g_{f_0} = 8.48. \] (18)

Note for the NNP coupling there is another form, namely, the so-called pseudovector coupling, \[ \mathcal{L}_{pv} = \frac{f}{m_\pi} \bar{\psi}(x) \gamma^5 \gamma^\mu \psi(x) \cdot \partial_\mu \phi(x). \] (19)

The pseudoscalar coupling and the pseudovector one are related by \[ \frac{f}{m_\pi} = \frac{g_P}{2m_N} \] (20) for free nucleon satisfying the Dirac equation. One may refer to Ref. [30] for more details. The pseudoscalar coupling was used in Ref. [17], although the pseudovector form of the Lagrangian appeared in the appendix of the paper.

Another essential and “dominant” piece should be the $X(1835)$ ($J^{PC} = 0^{-+}$) exchange since i) it locates near the $p\bar{p}$ threshold such that the propagator can enhance the contribution, and ii) the strong connection/relation between $X(1835)$ and the $p\bar{p}$ state. The first observation of $X(1835)$ (denoted by the $X$ particle below) was reported by BESII from the channel $J/\psi \to \gamma p\bar{p}$ \[ [38], \] where the mass reads $1859^{+31+5}_{-10-25}$ MeV with the statistic and systematic errors, in order, by using the $S$-wave Breit-Wigner function. The huge enhancement of the event distribution near the $p\bar{p}$ threshold was interpreted as the effect due to the $p\bar{p}$ final-state interaction (FSI) \[ [39], \] where the Watson-Migdal approach is exploited, i.e., the amplitude for $J/\psi \to \gamma p\bar{p}$ is expressed by a normalization constant multiplied by the $p\bar{p}$ scattering $T$-matrix. A refit with the inclusion of the fixed FSI factor introduced in Ref. \[ [39] \] has been carried out in a subsequent publication \[ [40]. \] The resulting mass is slightly changed and reads $1826.5^{+13.0}_{-3.4}$ MeV \[ [1]. \] From the state-of-the-art viewpoint, such FSI treatment has been superseded by the outcome of Ref. \[ [18], \] where the total amplitude $A$ is written as $A = A_0 + A_0G_0T$ with $A_0$, $G_0$, $T$ denoting bare production amplitude without FSI, free Green function and $p\bar{p}$ scattering $T$-matrix, respectively. One can refer to the review in Ref. \[ [20] \] for more details. It has been shown that the threshold enhancement could be indeed a $p\bar{p}$ bound state \[ [18]. \] However, one should be cautious that the pure FSI explanation proposed in Ref. \[ [39] \] reproduces the data very well and thus cannot be excluded. To date, the nature or even its existence of $X(1835)$ still remains mysterious. However, $X(1835)$ can be viewed, at least, as a poor man’s approach or an effective way to incorporate the strong FSI of $p\bar{p}$, and in this respect, we include it as a subthreshold resonance in our meson-exchange model calculation.

Note that the $X(1835)$ has also been observed in $\gamma \eta' \pi \pi$ channel with a statistical significance of $7.7\sigma$ \[ [40]. \] So, it could be most likely a mixing state of $p\bar{p}$ and $ss$ and this idea has been investigated in e.g., Ref. \[ [41. \] Then one may write \[ |X(1835)) = c_1|p\bar{p}\rangle + c_2|ss\rangle, \] (21) with $|c_1|^2 + |c_2|^2 = 1$. The maximum production for $D_s^+ \to p\bar{p}e^+\nu_e$ corresponds to $c_1 = c_2 = 1/\sqrt{2}$. The Lagrangian for the $X(1835)p\bar{p}$ coupling is of the same form as $p\eta\eta$, and we will take the coupling constant $g_{Xp\bar{p}} \approx 3.5$ \[ [42] \] provided that $X(1835)$ is a pure baryonium. This value agrees with the one given in Ref. \[ [43] \] after applying the Weinberg compositeness theorem \[ [12, 44–46] \], i.e., the coupling $g_{Xp\bar{p}}$ is obtained from the vanishing wave function renormalization. In the case of Eq. \[ (21) \], the true $Xp\bar{p}$ coupling will be multiplied by a factor of $1/\sqrt{2}$, while the transition $D_s^+ \to X(1835)$ will proceed via the form factors $F_1^{D_s^-\eta}/\sqrt{2}$ and $F_0^{D_s^-\eta}/\sqrt{2}$ \[ [16, 23, 28] \].

Combining all these ingredients together, we obtain
respectively, with the convention $\epsilon^{0123} = 1$. The amplitude modulus squared reads

$$|T|^2 = G_F^2 |V_{cs}|^2 \left( |f_4|^2 (s - 4m_p^2) + s |g_4|^2 \right) + \left( m_{D_s^+}^2 + (s + s_l)(s - s_l - 2m_{D_s^+}^2) \right) - 4X^2 (s, s_l) \cos^2 \theta_l, \tag{29}$$

and all other form factors vanish, where

$$\Delta S[p] = \frac{m_{D_s^+}^2 - m_{S[p]}^2}{s_l}. \tag{26}$$

To evaluate the amplitude modulus squared, we introduce the hadronic and leptonic tensor currents given by

$$\mathcal{H}^{\mu\nu} = h^{\mu} h^{\nu*} = 2 \left[ (s - 4m_p^2) (f_3^2 f_4^\dagger P^\mu L^\nu + f_3 f_4^\dagger P^\nu L^\mu + s |g_3|^2 L^\mu L^\nu + s (g_3^* g_4 P^\mu L^\nu + g_3 g_4^* P^\nu L^\mu) \right), \tag{27}$$

and

$$\mathcal{L}_{\mu\nu} = l_{\nu} l_{\mu}^* = 4 \left[ (L^\mu L^\nu - N^\mu N^\nu) - g^{\mu\nu}(s_l - m_l^2) + \frac{1}{2} \epsilon_{\mu\nu\alpha\beta} L^\alpha N^\beta \right], \tag{28}$$

respectively, with the convention $\epsilon^{0123} = 1$. The amplitude modulus squared reads

$$\mathcal{B}(D_s^+ \to p\bar{p}\epsilon^+\nu_e) = \frac{1}{\Gamma_{D_s^+}} \int d\Gamma^5 = 3.5 \times 10^{-9}, \tag{30}$$

based on the mass of 1826.5 MeV for the X(1835) reported in PDG [1]. The uncertainties arise from various sources, for example, the coupling constants and the mass of X(1835). The dominant uncertainty should be ascribed to the precision on the X(1835) mass. If we use the mass 1859 MeV, which corresponds to the fit without the primitive treatment of FSI as we have already discussed above, the branching fraction will become

$$\mathcal{B}(D_s^+ \to p\bar{p}\epsilon^+\nu_e) = 1.5 \times 10^{-8}. \tag{31}$$

As noticed in passing, the X(1835) exchange should dominate due to its proximity to the $p\bar{p}$ state. The X(1835) alone will contribute to $1.42 \times 10^{-8}$ for the branching fraction after turning off the $\eta, \eta'$, $f_0(980)$ effects. If the mass of X(1835) is closer to the $p\bar{p}$ threshold, the branching fraction will be further increased. Considering the width of X(1835) (around 80 MeV) [18], the branching fraction will become smaller by a few times. In this sense, we prefer to emphasize the importance of the “precision” on the X(1835) mass measurement. By the end of 2018, around $10^{10} J/\psi$ data samples are going to be accumulated within the one-year running period [47] and both $J/\psi \to \gamma p\bar{p}$ and $J/\psi \to \gamma\eta'\pi\pi$ can be re-examined to improve the accuracy. The experimental situation will be further improved in the case of super tau-charm factory [19, 34] with the planned luminosity of 100 times as much as BESIII.

In principle, the diagram with one-baryon exchange can be also considered. There is the process $D_s^+ \to p\bar{n}$ followed by the neutron beta decay, $\bar{n} \to \bar{p}\epsilon^+\nu_e$ as depicted in Fig. [2]. As noticed before, the decay rate of
Experimental observations, e.g., the strange quark spin content of the nucleon has also been revealed in several studies. The strangeness of the proton \cite{52} or the rescattering of kaons \cite{53, 54}, low-energy processes, this gives a picture of how the general form factors (constructed from Lorentz structure) emerge.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{diagram2.png}
\caption{The diagram for baryon exchange: $D_s^+ \rightarrow p\bar{n}$ followed by the neutron beta decay, $\bar{n} \rightarrow \bar{p}e^+\nu_e$. It involves two weak vertices and thus the contribution is negligible.}
\end{figure}

Here we comment on the possible OZI violation. The question may arise from the large $\phi$ production rate in $p\bar{p}$ collisions compared to the $\omega$ one, which is attributed to either the intrinsic $s\bar{s}$ component in the wave function of the proton \cite{52} or the rescattering of kaons \cite{53, 54}, see also the reviews in Refs. \cite{55, 56}. The strangeness content of the nucleon has also been revealed in several experimental observations, e.g., the strange quark spin polarization, $\sigma_{xN}$ term, magnetic moment of the proton and the ratio of strange and non-strange quark flavor distributions. However, the weight of the strange content is still small such that it is not expected to make large influence on the current results. On the other hand, the low-energy $p\bar{p}$ scattering data can be fairly well reproduced without the inclusion of the $\phi$ meson exchange, as e.g., done in Ref. \cite{17}, for which we stick to the $p\bar{p}M$ coupling.

In the $B$ meson sector, the semileptonic baryonic decay $B^- \rightarrow p\bar{p}\ell^-\nu_\ell$ has been studied in \cite{77} where the form factors $f_1$ and $g_i$ with $i = 1, \cdots, 5$ defined in analog to Eq. (8) were obtained by fitting them to the available data of $B \rightarrow p\bar{p}M$ in conjunction with the pQCD counting rule for form factors. However, this pQCD argument is not applicable to our case as the energy release in $D_s^+ \rightarrow p\bar{n}$ transition is rather small. Moreover, we notice that the predicted branching fraction $B(B^- \rightarrow p\bar{p}\ell^-\nu_\ell) = 1.04 \times 10^{-4}$ in \cite{77} is too large compared to the experimental observation of order $6 \times 10^{-6}$.

\section{IV. CONCLUSION}

In this work we have discussed the unique decay $D_s^+ \rightarrow p\bar{p}e^+\nu_e$ with a proton-antiproton pair in the final state. It is the only semileptonic baryonic decay which is physically allowed in the charmed meson sector, besides the hadronic baryonic decay $D_s^+ \rightarrow p\bar{n}$. There is abundant physics in this channel. Its measurement will test our knowledge on baryonic weak decays and the low-energy $p\bar{p}$ interactions. Taking into account the contributions from the intermediate states $\eta, \eta', f_0(980)$ and $X(1835)$, we find that its branching fraction is $10^{-9} \sim 10^{-8}$. Our prediction can be tested by BESIII/BEPCII data and its measurement is ongoing.

\section{ACKNOWLEDGMENTS}

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