I. INTRODUCTION

The BarBar Collaboration recently measured the branching fraction for \( B^0 / \overline{B}^0 \rightarrow a_0^+ \pi^- \) to be \((6.2 \pm 3.0 \pm 1.1) \times 10^{-6}\) \([1]\). This is comparable to \( B(B^0 \rightarrow \pi^+ \pi^-) \approx 4.4 \times 10^{-6}\) \([2]\). The tree interaction does not contribute to \( B^0 \rightarrow a_0^- \pi^- \) in the factorization limit since \( a_0^- \) cannot be produced from the \( V - A \) current \([3]\). Contrary to what was alluded in Ref. \([3]\), however, this decay should be strongly suppressed even for the penguin interaction if the perturbative QCD picture is correct. Therefore the observed branching fraction should consist almost entirely of \( B^0 \rightarrow a_0^- \pi^+ \) and \( \overline{B}^0 \rightarrow a_0^+ \pi^- \). Consequently the time-dependent \( B^0 / \overline{B}^0 \rightarrow a_0^+ \pi^- \) decay will not be suitable for extraction of the angle \( \alpha \) since there is little \( B^0 / \overline{B}^0 \) interference in these channels \([3,4]\).

We shall first show that the \( B^0 \rightarrow a_0^+ \pi^- \) decay amplitude is power suppressed by \( 1/m_B \) for the penguin interaction if short-distance terms dominate. A numerical estimate will be given for the decay amplitude. We then examine other factorization-suppressed decays, in particular, the production of \( b_1^+ \) that shares the same chiral property with the \( a_0^+ \) production. The chance of measuring the \( B^0 / \overline{B}^0 \) interference is even slimmer for \( b_1^+ \pi^- \). The underlying assumption leading to these conclusions is that perturbative QCD is valid even for the final-state interactions below \( m_B \), i.e., the perturbative-QCD-improved factorization \([5,6]\). This assumption need to be tested. For this purpose we discuss helicity conservation of light quarks in the \( B \) decay, in general. It leads us to zero-helicity dominance in the charmed \( B \) decay into two mesons both with spin. This selection rule provides us the simplest test of perturbative QCD in final-state interactions.

II. FACTORIZATION-SUPPRESSED PROCESS

A. \( B^0 \rightarrow a_0^+ \pi^- \)

We study for definiteness the decay amplitudes of the \( B(bq) \) meson instead of the \( B \) meson. The tree interaction \( O_{1,2} \), i.e., \((b_L u_L)(\bar{u}_L d_L)\) and its color-crossing, cannot produce \( \pi^- \) from the factorized current. It cannot produce \( a_0^+ \) either since \( G \)-parity does not match between \( a_0^+ \) and \( \pi^-\mu d \). It appears therefore that the penguin interaction dominates in this decay. If so, we would have an opportunity to extract the weak angle \( \alpha \) from the time-dependent decay of \( B^0 / \overline{B}^0 \rightarrow a_0^+ \pi^- \) \([3,4]\). However, the penguin decay amplitude is suppressed by \( 1/m_B \) so that the \( B^0 / \overline{B}^0 \) interference is too small for this purpose.

The QCD penguin operators, normally referred to as \( O_{3,6} \), generate the scalar density \( \overline{u} \gamma^\mu d + \overline{u} \gamma^\mu u \) by crossing. In the quark model its matrix element for the \( a_0^+ \) production is

\[
\langle u(k)|\overline{d}(-k)|\pi d(0)\rangle = 2k \cdot (\sigma) \tag{1}
\]

where \( \langle \sigma \rangle = \chi^\dagger \chi \gamma' \) with \( \chi \) and \( \chi' \), the Pauli spinors of \( u \) and \( \overline{d} \), respectively. Superposing Eq. \((1)\) with the wave function \( \Phi(k) \) of a \( p \)-wave bound state, we define the \( a_0^+ \) decay constant \( f_{a0} \) by

\[
\langle a_0^+ |\pi d(0)\rangle = f_{a0} m_{a0}. \tag{2}
\]

In the \( B^0 \) rest frame where \( u \) and \( \overline{d} \) fly fast with total momentum \( p = p_u + p_d \),

\[
\langle u|\overline{d}(p)|\pi d(0)\rangle = 2k_\perp \cdot (\sigma_\perp) + 2E_k \beta(2x - 1)(\sigma_\parallel), \tag{3}
\]

where \( \beta = |p|/E_p \) \((E_p = \sqrt{m^2_{ud} + p^2})\) and \( x = |p_u|/|p| \) with \( \parallel \) and \( \perp \) referring to the parallel and perpendicular direction to the momentum \( p \). The right-hand side of Eq. \((2)\) is \( O(1) \) in \( E_p \) in the fast moving frame, which is consistent with the right-hand side of Eq. \((3)\) after superposition with the light-cone distribution function \( \Psi(x, k_\perp) \). It is one power lower in \( E_p \) than, for instance, the \( \pi^+ \) production from \( \overline{u} \gamma^\mu \gamma_5 d \). This difference is due partly to the Lorentz property, but more importantly to the chiral property of quark fields involved, namely, \( \overline{R}L \) vs \( \overline{L}L \pm \overline{R}R \). We can perform the same calculation for
the $^3P_J$ state, namely $a_2$. The matrix element of quark-pair production $\langle ud|\bar{u}d\gamma_{\mu}d|d0\rangle$ is equal to $2i(k\times\sigma)$ in the $a_1$ rest frame. If we superpose it with the same $p$-wave orbital wave function as $a_0$ and boost it to the $B^0$ rest frame, we obtain

$$\langle a_1^+(p)|\bar{u}d\gamma^\mu\gamma_5d|0\rangle = f_{sa_1m_a}e^\mu(p),$$

(4)

with $f_{sa_1} = \sqrt{2/3}f_{sa_0}$ in the constituent quark model. Since $f_{sa_1} \simeq f_{sa_0} \simeq f_s(\simeq f_3)$, a similar calculation for $a_1^-(p)$ leads us to $\langle a_1^-(p)|\bar{u}d\gamma^\mu\gamma_5d|0\rangle = f_{sa_2m_a}e^\mu(p)$ with $f_{sa_2} = \sqrt{1/2}f_{sa_0}$. The formation of $\pi^-$ through $\mathcal{O}_{5,6}$ in $B^0 \rightarrow a_0^+\pi^-$ is described by the scalar form factor of $B^0 \rightarrow \pi^-$,

$$\langle \pi^-|\bar{b}d(B^0)|m_BF_s(q^2)\rangle.$$

(5)

The relation $i\partial_\mu\langle \bar{B}\gamma^\mu u \rangle = (m_u - m_b)(\bar{u}b)$ relates $F_s(q^2)$ to the vector form factors of $B^0 \rightarrow \pi^-$ as $m_u m_B F_s(q^2) = (m_B^2 - m_\pi^2)F_1(q^2) + q^2F_2(q^2)$ so that $F_s(q^2) \simeq F_1(q^2)$ for $|q^2| \ll m_B^2$. Expressing the penguin amplitude $A_p(B^0 \rightarrow a_0^+\pi^-)$ in terms of the tree amplitude $(A_t)$ of $B^0 \rightarrow \pi^+\pi^-$, we have

$$\frac{A_p(B^0 \rightarrow a_0^+\pi^-)}{A_t(B^0 \rightarrow \pi^+\pi^-)} \sim \frac{2V_{ts}V_{td}^*C_1(m_b f_{sa_0} m_a m_B F_s(m_{sa_0}))}{V_{ub}V_{cd}^*C_1(m_b f_{sa_0} m_a m_B F_s(m_{sa_0}))},$$

$$\approx 0.04 \times \left[ F_s(m_{sa_0})/F_1(m_{sa_0})\right],$$

(6)

where $C_1$ and $C_2$ are from the Wilson coefficients of the tree and penguin operators, respectively. The value $|V_{td}| \approx 0.01$ has easily 50% uncertainty. Since the form factors scale as $1/(1 - q^2/m_B^2)$ with a relevant B-meson mass in the simple pole approximation, we may set $F_s(m_{sa_0}) \simeq F_s(m_{sa_2}) \approx F_1(m_{sa_2})$. Then the right-hand side is less than a tenth. The main source of this suppression is the kinematical factor $1/m_B$ that is traced to the chiral suppression of $a_{0a}$ production in the fast moving frame.

In comparison, the tree-allowed decay $B^0 \rightarrow a_0^\pm\pi^\mp$ is given for $f_{sa_0} \simeq f_s$ by

$$\frac{A_t(B^0 \rightarrow a_0^+\pi^-)}{A_t(B^0 \rightarrow \pi^+\pi^-)} \sim \frac{F_1^a(m_{sa_2})}{F_1^a(m_{sa_0})},$$

(7)

where $F_1^a(q^2)$ is the axial-vector form factor $\propto (p_B + p_a)^a$ of $B^0 \rightarrow a_0^-$. Since $F_1^a(m_{sa_2}) \approx F_1^a(m_{sa_0})$ is not far out of line, the right-hand side is about unity. Then Eqs. (6) and (7) are in line with the measured branching fraction [1] if it consists almost entirely of $B^0 \rightarrow a_0^\pm\pi^\mp$ and $\mathcal{T}_1 \rightarrow a_0^\pm\pi^-$. It means little $B^0-\mathcal{B}^0$ interference in the $a_0^\pm\pi^\mp$ channels.

For the decay $B(B^0 \rightarrow a_{0a}^0\pi^0)$, all of $\mathcal{O}_{1,2,3,4}$, and $\mathcal{O}_{5,6}$ possibly contribute with comparable magnitudes so that their sum involves much a larger uncertainty. Therefore a clean extraction of the weak angles will be difficult from $B^0-\mathcal{B}^0 - a_{0a}^0\pi^0$. It is obvious that Eqs. (6) and (7) apply to the ratios of $B^0 \rightarrow a_{0a}^0\rho^\mp$ to $B^0 \rightarrow \pi^+\rho^\mp$ as well, if we replace the transition form factors appropriately.

The QCD corrections turn the local operator $\bar{u}(x)\Gamma_q(x)$ into the nonlocal operator $\bar{u}(x)\gamma(y)$. But $\bar{u}(x)\gamma(y)$ can be expanded in the Taylor series of local operators in powers of $(x-y)\mu\partial^\mu$. If short-distance interactions dominate, $|x-y|$ is a fraction of $1/m_B$ so that all terms of the Taylor expansion have the same $1/m_B$ dependence as the leading term, i.e., the local operator. Evaluation of the higher derivative terms require knowledge of more than a decay constant. When a meson cannot be produced from a local operator appearing in the effective interactions $\mathcal{O}_i$, it should be noted that the leading contribution is accompanied by $\alpha_s/\pi$ due to a QCD loop. This is the case for $b_1$ and $a_2$. In such a case the amplitude of the leading order in $1/m_B$ depends on the shape of the distribution function of the $p$-wave bound states of which we know less.

If $a_0$ is a four-quark state $qq\bar{q}q$ instead of a $q\bar{q}$ in $^3P_0$, we can show with a dimension argument of perturbative QCD that the $B^0 \rightarrow a_0^\pm\pi^\mp$ decay amplitudes are even more suppressed in $1/m_B$ than $A_p(B^0 \rightarrow a_0^\pm\pi^\mp)$. The positive identification of $B^0/\mathcal{B}^0 \rightarrow a_0^\pm\pi^\mp$ [1] is an evidence against the four-quark assignment of $a_0$ or else for breakdown of perturbative QCD.

**B. $B^0 \rightarrow b_1^-\pi^+$**

Production of other $^3P_J$ states, $a_1^+\pi^-$ and $a_2^+\pi^-$, is different from that of $a_0^-$. They are produced from $(\pi\gamma^\mu d_R) - (\bar{u}\gamma^\mu d_L)$ and $(\pi\gamma^\mu d_R) + (\bar{u}\gamma^\mu d_L)$, respectively. While the axial-vector current is found in the tree interaction, the tensor operator must be generated by gluon corrections. Therefore its derivative $\partial^\nu$ comes with $\alpha_s/\pi$ and also with $1/E \simeq |x-y| = O(2/m_B)$. Since the tree interaction contributes in full strength to production of $a_1^+$ and of $a_2^+$, these decays are simply related to the tree-dominated $B^0 \rightarrow \pi^+\pi^-$ decay as:

$$\frac{B(B^0 \rightarrow a_1^+\pi^-)}{B(B^0 \rightarrow \pi^+\pi^-)} \approx \frac{f_{sa_1} F_1^a(m_{sa_1})}{f_E F_1(m_{sa_1})} \simeq 1,$$

$$\frac{B(B^0 \rightarrow a_2^+\pi^-)}{B(B^0 \rightarrow \pi^+\pi^-)} \approx \frac{\alpha_s(E)}{\pi} \left| \frac{f_{sa_2} E f_2 F_1(m_{sa_2})}{f_E F_1(m_{sa_2})} \right|^2.$$  

(8)

While a considerable uncertainty exists in the relevant value of $\alpha_s$, we reasonably expect from the second line with $E \simeq E_{sa_2}$ that $B(B^0 \rightarrow a_2^+\pi^-) \lesssim 10^{-2} \times B(B^0 \rightarrow \pi^+\pi^-)$. Similarly, $B(B^0 \rightarrow a_1^+\pi^-) \approx B(B^0 \rightarrow \pi^+\pi^-)$ and $B(B^0 \rightarrow a_2^+\pi^-) \lesssim 10^{-2} \times B(B^0 \rightarrow \pi^+\pi^-)$.

Production of $b_1^- (^1P_1)$ has similarity with $a_0$ in the chiral structure and with $a_2$ in the $\alpha_s/\pi$ suppression. The local operator that matches the quantum numbers of $b_1^-$ is $i(\pi\gamma_5 \partial^\nu d)$, whose chiral property is $\mathcal{T}_R - \mathcal{T}_L$. In the $b_1$ rest frame,

$$\langle u(k)|\bar{d}(-k)|\pi\gamma_5 \partial^\nu d|0\rangle = 2E_k k(1).$$

(9)
where $\langle 1 \rangle = \chi^\dagger \chi'$. Superposing Eq. (9) with the same $p$-wave orbital function as the $^{3}P_J$ mesons, we obtain

$$
\langle b_{0}(p)\bar{m}^{\dagger}_{\gamma_{5}}\partial^{\mu}d|0\rangle = f_{b_{0}}m_{b_{0}}^{2}\varepsilon^{\mu}(p)
$$

with $f_{b_{0}} = \frac{1}{2}f_{b_{1}}$ in the quark model. In the $B^{0}$ rest frame,

$$
\langle u\bar{d}(p)\bar{m}^{\dagger}_{\gamma_{5}}\partial^{\mu}d|0\rangle = \begin{cases} 2\gamma\beta E_{k}k_{\mu}(1), & (\mu = 0) \\ 2\gamma E_{k}k_{\mu}(1), & (\mu = \parallel) \\ 2E_{k}k_{\mu}(1), & (\mu = \perp), \end{cases}
$$

where $\gamma = (1 - \beta^{2})^{-1/2}$ so that $\gamma E_{k} \simeq E_{p}$, the $b_{0}$ energy in the $B^{0}$ frame. The right-hand side of Eq. (11) is therefore $O(E_{p})$ for the time and longitudinal components in the fast moving frame, and $O(1)$ for the transverse components, which is consistent with Eq. (10).

As pointed out above, the derivative $\partial^{\mu}$ is accompanied by $\alpha_{s}(E)/\pi E$ so that it does not enhance the high-energy behavior when short distances dominate, i.e., $E = O(1/m_{b})$. As in the case of $a_{0}^{+}$, the high-energy behavior of $LR - RL$ is lower by one power of energy in the fast moving frame. Consequently the decay branching fraction $B(B^{0} \to b_{1}^{+}\pi^{-})$ just like $B(B^{0} \to a_{0}^{+}\pi^{-})$, namely $1/m_{B}^{2}$ down relative to the allowed-tree branching fraction:

$$
\frac{B(B^{0} \to b_{1}^{+}\pi^{-})}{B(B^{0} \to a_{0}^{+}\pi^{-})} \simeq \left(\frac{\alpha_{s}(E)}{\pi}\right)^{2} \left(\frac{F_{1}(m_{b}^{2})}{F_{1}(m_{a_{0}}^{2})}\right)^{2} \simeq \left(\frac{\alpha_{s}(E)}{\pi}\right)^{2},
$$

(12)

where $f_{b_{1}} \approx f_{b_{0}}$ has been used. Eq. (12) is the prediction of short-distance dominance. We expect the same relation for the tree-allowed decay modes: $B(B^{0} \to b_{1}^{+}\pi^{+}) \simeq (\alpha_{s}/\pi)^{2}B(B \to a_{0}^{+}\pi^{+})$. We should keep in mind that the prediction on $b_{1}^{+}\pi^{+}$ involves the same uncertainties as we have mentioned for $a_{2}^{+}\pi^{+}$ at the end of the preceding section. Nonetheless, it is safe to state with Eqs. (6), (7), and (12) that $b_{1}^{-}\pi^{+}$ will not be produced at the level of $10^{-7}$ or higher if the perturbative QCD picture is correct for final-state interactions. The $b_{1}^{-}\pi^{+}$ decay should not be seen at any level ($<10^{-9}$) in that case.

### III. HELICITY CONSERVATION AND SPIN STRUCTURE

The predictions in the preceding sections lead us to conclude that the $B^{0}/\bar{B}^{0}$ interference can be observed in the $a_{0}^{+}\pi^{+}$ or the $b_{1}^{+}\pi^{+}$ channel only if a very strong enhancement occurs by long-distance interactions in an otherwise suppressed mode. In such a case, the classification of amplitudes by the tree and the penguin interaction becomes a bit blurry.

One powerful test exists for the short-distance dominance. For a two-body decay where both final mesons have spin, $J$ and $J'(\geq J)$, we have $2J + 1$ independent amplitudes. If short-distance interactions dominate, the decay into the zero-helicity state should dominate over all other helicity states in the charmless $B$ decay. Let us explain it briefly since this is a robust prediction of the Standard Model and provides a simple experimental test of short-distance dominance independent of the rate measurement.

For the tree operators in which all quark fields are left-chiral, one of the final mesons must be formed with the fast $q_{L}$ and $\overline{q}_{L}$ without involving the spectator. Since these quarks remember their helicities throughout interactions with hard gluons, the resulting $q_{L}\overline{q}_{L}$ meson state must be in helicity zero ($h = 0$) in the approximation of ignoring the quark mass and the higher configuration such as $q_{L}q_{R}$ and $q_{R}q_{R}$. Then the other meson that picks up the spectator is forced to have $h = 0$ by overall angular momentum conservation. This argument applies to the penguin operators $O_{3},O_{4}$ too. In the case of $O_{5,6}$, the argument is the same when $q_{R}$ and $q_{L}$ form one meson. If instead $\overline{q}_{L}(\overline{q}_{L})$ and $q_{R}$ form a meson, this meson would have $h = +1$, and the other meson, being formed with $\overline{q}_{R}$ and the spectator, can only be in $h = 0$ or $-1$. The overall angular conservation therefore forbids this decay. No matter which interaction causes a decay, therefore, the dominant final helicity state is zero.\(^1\) Production of $h = \pm 1$ is allowed to the extent of the nonzero quark masses and of the transverse motion of quarks inside a meson, i.e., the meson mass.

A remark is in order concerning the quark mass effect in the $a_{0}$ production. The $a_{0}^{+}$ meson is produced by the operator $\overline{q}_{R}q_{L}$ of $O_{5,6}$ which produces $h = +1$, while the spinless $a_{0}^{0}$ cannot have nonzero helicity. This means that, if we ignore the quark mass, production of $a_{0}^{+}$ is forbidden. The $a_{0}^{+}$ production occurs through the small $h = 0$ component of $O(\sqrt{m_{a_{0}}^{2} + k_{L}^{2}/E_{p}})$ that is contained in $\overline{q}_{L}q_{R}$. This remark applies to $b_{1}^{0}$ too. That is the reason why the production amplitudes of $a_{0}^{+}$ and $b_{1}^{0}$ are suppressed by one power of $1/m_{B}$.

When the mass and $k_{\perp}$ corrections are included, the $h = +1$ amplitude is generated with $O(m/E)$ for the $B(b_{0})$ decay while the $h = -1$ amplitude is generated only with $O(m^{2}/E^{2})$. The reason is as follows: The spectator quark can be either in $h = +\frac{1}{2}$ or $-\frac{1}{2}$ with a 50/50 chance. Therefore the $h = +1$ decay amplitude can be realized with the small opposite helicity component of a single fast $q_{L}$ (out of $q_{L}q_{L}\overline{q}_{L}$ from $O_{1,4}$) or $q_{R}$ (out of $q_{R}q_{R}\overline{q}_{R}$ from $O_{5,6}$). On the other hand the $h = -1$ amplitude needs small components of two $q_{L}$’s (from $O_{1,4}$) or $q_{R}$ and $\overline{q}_{R}$ (from $O_{5,6}$). Since each small component costs $m_{T}/E_{p} \sim 2m_{T}/m_{B}$, we obtain the following hier-

\(^1\)Actually, it is sufficient to prove this selection rule for the fundamental weak interaction ($\sim O_{2}$) since all other decay operators are generated from it through hard gluon-loop corrections.
where \( m \) and \( m' \) are the meson masses. In the first line, \( m \) stands for the mass of the meson formed with the fast \( q \) and \( \overline{q} \) emitted from the current, not with the spectator. Because the small opposite helicity of this meson forces the other meson to have \( h = +1 \) state.

The \( h = 0 \) dominance can be easily tested in experiment by measuring the angular distribution of the decay products of either of final mesons [8] although distinguishing between \( h = +1 \) and \( -1 \) is much harder [9]. Actually, the zero-helicity dominance holds for some processes involving a charm quark such as \( B^0 \to D^{*-} \rho^+ \) for which the tree interaction completely dominates. Since \( \rho^+ \) is produced from the \( V - A \) current of the tree interaction, its helicity must be zero. Therefore the \( h = 0 \) dominance should hold to \( O(m_\rho/m_B) \) in this case. The percentage of the \( h = 0 \) branching was found in experiment as \( \Gamma_L = (93 \pm 5 \pm 5)/6 \% \) for \( B^0 \to D^{*-} \rho^- \) [10]. The selection rule holds in a modified form for \( B \to J/\psi K^* \) [12]. A significant correction may arise to the \( h = 0 \) dominance because of the large c-quark mass. However, the \( K^* \) can only be in \( h = 0 \) or \( +1 \) for \( m_{K^*} \ll m_B \) and therefore \( |A_{+1}| \gg |A_{-1}| \) should hold well. Experiment [11] is consistent with \( |A_{+1}| \gg |A_{-1}| \), but cannot distinguish it from \( |A_{-1}| \gg |A_{+1}| \) until more a sophisticated cascade decay measurement is made. The zero-helicity dominance rule is yet to be tested in the charmless \( B \) decay.

The zero-helicity dominance can be found in the early literature. Ali et al [13] computed the \( B \to VV \) decay in the factorization of the tree interaction with the \( U(6) \times U(6) \) form factors which happen to incorporate helicity conservation. One can read off the suppression pattern of Eq. (13) in their results, as Chen et al., [14] recently pointed out. Körner and Goldstein [15] discussed it for charm decays. If the final-state interaction is included through hadron rescattering, the helicity selection rule breaks down in general. We can show, for instance, that if one computes \( B \to \pi \pi \to \rho \rho \) with \( \omega \)-exchange for \( \pi \pi \to \rho \rho \), the final \( \rho \)'s are polarized dominantly in \( h = \pm 1 \). This illustrates that the \( h = 0 \) dominance easily breaks down by long-distance final-state interactions. If on the other hand one computes the \( \pi \pi \to pp \) rescattering with \( \pi \)-exchange, the \( \rho \)'s are polarized mostly in \( h = 0 \). Whether short-distances dominate or not should be determined eventually by experiment.

Test of the zero-helicity dominance has one clear advantage over the branching fraction measurement. There is a limitation in accuracy even in the short-distance calculation when one evaluates the magnitudes and phases of amplitudes and sum them up. In the hadron picture, it is nearly impossible to compute final-state interactions and predict decay amplitudes reliably. Therefore the branching fraction measurement alone will be inconclusive in determining how much long-distance contributions exist in a given decay process. If we combine it with the zero-helicity dominance test, we can be more confident with our conclusions since the rule should hold to all orders of QCD as long as they are of short distances.

**IV. SUMMARY**

If short-distance physics dominates in final states, it will be nearly impossible to obtain the weak angle from the time-dependent decay \( B^0 / \overline{B}^0 \to a_0^\pm \pi^\mp (\rho^\mp) \) or \( b_1^\pm \pi^\mp (\rho^\pm) \). While the flavor-tagged measurement will tell us about short-distance dominance in these decays, more a general and simple test is to examine zero-helicity dominance in the charmless \( B \) decay with spins.

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