Decays and productions via bottomonium for $Z_b$ resonances and other $B\bar{B}$ molecules

S. Ohkoda$^1$, Y. Yamaguchi$^1$, S. Yasui$^2$, and A. Hosaka$^1$

$^1$Research Center for Nuclear Physics (RCNP),
Osaka University, Ibaraki, Osaka, 567-0047, Japan and
$^2$KEK Theory Center, Institute of Particle and Nuclear Studies,
High Energy Accelerator Research Organization,
1-1, Oho, Ibaraki, 305-0801, Japan

Abstract

We discuss decays and productions for the possible molecular states formed by bottom mesons $B$ ($B^*$) and $\bar{B}$ ($\bar{B}^*$). The twin resonances found by Belle, $Z_b^{\pm,0} (10610)$ and $Z_b^\pm (10650)$, are such candidates. The spin wave functions of the molecular states are rearranged into those of heavy and light spin degrees of freedom by using the re-coupling formulae of angular momentum. By applying the heavy quark symmetry we derive model independent relations among various decay and production rates, which can be tested in experiments.

PACS numbers: 14.40.Rt, 12.39.Hg, 13.20.Gd
Exotic hadrons are important for the study of hadron physics. Not only multiquarks, but also hadronic molecules, provide an alternative opportunity for the study of structures of hadrons. Hadronic molecules are especially interesting because they can appear in the threshold region of hadron-quark-clusters which saturate the color dependent force of QCD and weakly interact each other. Thus, they emerge as loosely bound states of hadrons with spatially extended property, which should be distinguished from the conventional hadrons where the minimal number of constituent quarks are considered to be rather localized.

One of the good candidates of such a hadronic molecule is a pair of $Z_b(10610)$ and $Z_b(10650)$ which has been recently observed in the processes $\Upsilon(5S) \rightarrow \pi\pi\Upsilon(nS)(n = 1, 2, 3)$ and $\Upsilon(5S) \rightarrow \pi\pi h_b(kP)(k = 1, 2)$ by Belle group \[1, 2\]. The reported masses and widths of the two resonances are $M(Z_b(10610)) = 10607.2 \pm 2.0$ MeV, $\Gamma(Z_b(10610)) = 18.4 \pm 2.4$ MeV and $M(Z_b(10650)) = 10652.2 \pm 1.5$ MeV, $\Gamma(Z_b(10650)) = 11.5 \pm 2.2$ MeV, which appear very close to the $B\bar{B}^*$ and $B^*\bar{B}^*$ thresholds, respectively, and have relatively narrow decay widths. The facts that (1) the minimal constituent quarks must be four because of isospin one in their observed quantum numbers $I^G(J^{PC}) = 1^+(1^{+-})$, (2) the mass difference between $Z_b(10610)$ and $Z_b(10650)$ is so small as 40 MeV, and (3) the decays into a bottomonium both through heavy spin conserved and non-conserved processes \[3\] indicate that $Z_b(10610)$ and $Z_b(10650)$ are likely $B^{(s)}\bar{B}^{(s)}$ molecular states \[5, 6\]. Here, we use a notation $B^{(s)}$ to stand for $B$ or $B^*$. 

In our previous work, we have investigated all possible $B^{(s)}\bar{B}^{(s)}$ bound or resonant states having exotic quantum numbers of the total angular momentum $J \leq 2$ \[7\]. To verify whether these theoretically expected states exist or not in experiments, it is useful to study their production and decay processes. In this paper, we derive relative rates of each transition when we consider that $B^{(s)}\bar{B}^{(s)}$ molecular states obey the heavy quark symmetry. The heavy quark symmetry allows only the processes where heavy quark spin is conserved, leading to selection rules among certain classes of transitions. To derive them, we consider the spin structure of the mesons in terms of spin re-coupling formula which is equivalent to Fierz rearrangement. By rearranging the two heavy quarks in $B^{(s)}$ and $\bar{B}^{(s)}$ mesons of a molecular state, we can separate the heavy quark spin and the spin of light degrees of freedom in heavy quark limit.

In general, the total angular momentum $J$ of a hadron is a conserved quantity. The spin of heavy quark $S_H$ is also conserved in the limit of heavy quark mass, $m_Q \rightarrow \infty$. Then we
can define the spin of light degrees of freedom $S_l$ by
\[
S_l \equiv J - S_H,
\]
which is also conserved. Generally speaking, the spin of light degrees of freedom has complex structure. For instance, a $Q\bar{q}$ meson includes an anti-quark $\bar{q}$, gluons, an arbitrary number of $q\bar{q}$ pairs and angular momentum $L$ as the light degrees of freedom. Although they are not conserved separately, the sum of them is conserved in the heavy quark limit. This quantity is referred to simply as “light spin” because it includes all degrees of freedom except for the heavy quark spin. Therefore, we can describe the spin structure of a hadron containing heavy quarks in terms of the good quantum numbers $S_H$ and $S_l$. As an example, the spin structure of $B$ meson is $B(0^-) = (|\frac{1}{2}H \otimes \frac{1}{2}l|)_{J=0}$ and that of $B^*$ meson is $B^*(1^-) = (|\frac{1}{2}H \otimes \frac{1}{2}l|)_{J=1}$.

Now we consider $Z_b(10610)$ and $Z_b(10650)$. We assume that the main component of the wave function of $Z_b(10610)$ is $\frac{1}{\sqrt{2}}(B\bar{B}^* - B^*\bar{B})(^3S_1)$ and that of the $Z_b(10650)$ is $B^*\bar{B}^*(^3S_1)$. Because these masses are close to $B\bar{B}^*$ and $B^*\bar{B}^*$ thresholds respectively, and the rate of $D$-wave mixing is not large as the previous study indicates that the probability of $\frac{1}{\sqrt{2}}(B\bar{B}^* - B^*\bar{B})(^3D_1)$ is about 9% and that of $B^*\bar{B}^*(^3D_1)$ is about 6% in the total wave function of $Z_b(10610)$ [7]. Let us now employ the spin re-coupling formula with 9-j symbols to analyze the spin structure of $\frac{1}{\sqrt{2}}(B\bar{B}^* - B^*\bar{B})(^3S_1)$ and $B^*\bar{B}^*(^3S_1)$. This standard formula is written as
\[
[[l_1, s_1]^{j_1}, [l_2, s_2]^{j_2}]^J = \sum_{L,S} \hat{j}_1 \hat{j}_2 \hat{L}\hat{S} \begin{pmatrix}
    l_1 & s_1 & j_1 \\
    l_2 & s_2 & j_2
\end{pmatrix}
\begin{pmatrix}
    L & S & J
\end{pmatrix}
[[l_1, l_2]^{L}, [s_1, s_2]^{S}]^J,
\]
where $[j_1, j_2]^J$ means that the angular momenta $j_1$ and $j_2$ are coupled to the total angular momentum $J$, and $\hat{J} = \sqrt{2J + 1}$. By using this, the heavy and light spins of $B\bar{B}^*(^3S_1)$ and $B^*\bar{B}^*(^3S_1)$ are re-coupled as
\[
|B\bar{B}^*(^3S_1)\rangle = [[b\bar{q}]^0, [\bar{b}q]^1]^1
\]
\[
= \sum_{H, \hat{l}} \hat{0}\hat{l}\hat{H}\hat{l} \begin{pmatrix}
    1/2 & 1/2 & 0 \\
    1/2 & 1/2 & 1
\end{pmatrix}
\begin{pmatrix}
    [b\bar{b}]^H, [\bar{q}q]^l
\end{pmatrix}
= \frac{1}{2} [[b\bar{b}]^0, [\bar{q}q]^1] + \frac{1}{2} [[b\bar{b}]^1, [\bar{q}q]^0] + \frac{1}{\sqrt{2}} [[b\bar{b}]^1, [\bar{q}q]^1]
\]
\[
= \frac{1}{2} (0_{\hat{H}} \otimes 1_{\hat{l}}) - \frac{1}{2} (1_{\hat{H}} \otimes 0_{\hat{l}}) + \frac{1}{\sqrt{2}} (1_{\hat{H}} \otimes 1_{\hat{l}}),
\]
\[
(|B\bar{B}^*(^3S_1)\rangle, |B^*\bar{B}^*(^3S_1)\rangle).
\]
\[ |B^* \bar{B}^{(3S_1)} \rangle = [[b \bar{q}]^1, [b \bar{q}]^1]^0 \]
\[ = -\frac{1}{2} (0_H^+ \otimes 1^-) + \frac{1}{2} (1_H^- \otimes 0^-) + \frac{1}{\sqrt{2}} (1_H^- \otimes 1^-), \] (4)

which give the spin structure of \( \frac{1}{\sqrt{2}}(B \bar{B}^* - B^* \bar{B})^{(3S_1)} \). For \( B^* \bar{B}^{*(3S_1)} \), we have
\[ |B^* \bar{B}^{*(3S_1)} \rangle = [[b \bar{q}]^1, [b \bar{q}]^1]^1 \]
\[ = \sum_{H', l} \frac{1}{\sqrt{2}} \left[ [b \bar{b}]^{10}, [\bar{q} \bar{q}]^{10} \right] + \frac{1}{\sqrt{2}} \left[ [b \bar{b}]^1, [\bar{q} \bar{q}]^1 \right] \]
\[ = \frac{1}{\sqrt{2}} (0_H^- \otimes 1^-) + \frac{1}{\sqrt{2}} (1_H^- \otimes 0^-). \] (5)

If the structure of \( Z_b \)'s is dominated by \( B^{(*)} \bar{B}^{*(3S_1)} \), their spin configurations are given from (3)-(5) as
\[ |Z_b(10610)\rangle = \frac{1}{\sqrt{2}} (0_H^- \otimes 1^-) - \frac{1}{\sqrt{2}} (1_H^- \otimes 0^-), \] (6)
\[ |Z_b(10650)\rangle = \frac{1}{\sqrt{2}} (0_H^- \otimes 1^-) + \frac{1}{\sqrt{2}} (1_H^- \otimes 0^-). \] (7)

It is important to note that \( Z_b \)'s have the same fraction of a heavy quark spin singlet component and a triplet component. Due to the heavy quark spin symmetry, the former component decays into spin-singlet bottomonium \( h_b(kP) \) (\( |h_b(kP)\rangle = 0_H^- \otimes 1^- \)) with a pion emission in \( S \)-wave. By contrast, the latter decays into a spin-triplet bottomonium \( \Upsilon(nS) \) (\( |\Upsilon(nS)\rangle = 1_H^- \otimes 0^+ \)) with a pion emission in \( P \)-wave. Therefore, the rates of the decays into a spin singlet state \( Z_b \to h_b(kP)\pi \) should be comparable to that of the decays into a spin triplet state \( Z_b \to \Upsilon(nS)\pi \). Hence, the experimental observation of the two pion emission of \( \Upsilon(5S) \) can be explained if the process occurs through the intermediate state of \( Z_b \) \( \square \). These arguments have been already made by Fierz transformation in Ref. \( \square, \square \). Here we have shown the same results in terms of the spin re-coupling formula \( \square \), which is also applied to other processes in a straightforward manner.

Next we consider the neutral resonance \( Z_b^0(10610) \) recently observed in the processes \( \Upsilon(5S) \to \pi^0 \pi^0 \Upsilon(1S, 2S) \) by Belle group \( \square \). It is possible for \( Z_b^0(10610) \) to decay into \( \Upsilon \pi^0, \)
\( h_b \pi^0, \eta_b \gamma(\rho^0) \) and \( \chi_{bJ}\gamma \) (\( J = 1, 2, 3 \)) from the viewpoint of the conservation of quantum numbers and kinematics. In general, the decay \( \Upsilon(5S) \to \eta_b \pi^0 \gamma(\rho^0) \) should be suppressed
because this process requires heavy quark spin flip. However, going through $Z_b$, the decay into the singlet bottomonium state is allowed.

$Z^0_b(10610)$ can decay into $\chi_{bJ}$ ($J = 1, 2, 3$) by a photon emission. The radiative transition into a bottomonium is a new decay pattern for $Z_b$, which cannot be seen in charged $Z^\pm_b$. It can be used to investigate the structure and interaction of $Z_b$. Here, we estimate the ratio of decays $Z^0_b \rightarrow \chi_{bJ}\gamma$. To do this, we derive the spin structure of $\chi_{bJ}\gamma$. Since $Z^0_b(10610)$ has $I(J^PC) = 1(1^-)$ as a possible isospin partner of $Z^\pm_b(10610)$, the photon of the $\chi_{b0}\gamma$ decay channel must carry $J^P = 1^+$, which corresponds to an M1 transition. Therefore, the spin structure of the photon can be written as $|\gamma(M1)\rangle = 0^+_H \otimes 1^+_l$. The spin structure of $\chi_{b0}$ is $|\chi_{b0}\rangle = (1^-_H \otimes 1^-_l)|J=0\rangle$. Applying the re-coupling formula to the spin of $\chi_{b0}$ and photon, we find

$$|\chi_{b0}\gamma(M1)\rangle = (1^-_H \otimes 1^-_l)|J=0\rangle \otimes (0^+_H \otimes 1^+_l) = \frac{1}{3}(1^-_H \otimes 0^-_l) - \frac{1}{\sqrt{3}}(1^-_H \otimes 1^-_l)|J=1\rangle + \frac{\sqrt{5}}{3}(1^-_H \otimes 2^-_l)|J=1\rangle. \quad (8)$$

The same consideration is applied to the decay $Z^0_b(10610) \rightarrow \chi_{b1}\gamma$, where M1 and E2 transitions are possible. The spin structures of these decay channels are given by

$$|\chi_{b1}\gamma(M1)\rangle = -\frac{1}{\sqrt{3}}(1^-_H \otimes 0^-_l) + \frac{1}{2}(1^-_H \otimes 1^-_l)|J=1\rangle + \frac{15}{6}(1^-_H \otimes 2^-_l)|J=1\rangle, \quad (9)$$

$$|\chi_{b1}\gamma(E2)\rangle = -\frac{1}{2}(1^-_H \otimes 1^-_l)|J=1\rangle + \frac{\sqrt{3}}{2}(1^-_H \otimes 2^-_l)|J=1\rangle. \quad (10)$$

The decay $Z^0_b(10610) \rightarrow \chi_{b2}\gamma$ can also occur in M1 and E2 transitions and the spin structures of these states are given by

$$|\chi_{b2}\gamma(M1)\rangle = \frac{\sqrt{5}}{3}(1^-_H \otimes 0^-_l) + \frac{\sqrt{15}}{6}(1^-_H \otimes 1^-_l)|J=1\rangle + \frac{1}{6}(1^-_H \otimes 2^-_l)|J=1\rangle, \quad (11)$$

$$|\chi_{b2}\gamma(E2)\rangle = \frac{\sqrt{3}}{2}(1^-_H \otimes 1^-_l)|J=1\rangle + \frac{1}{2}(1^-_H \otimes 2^-_l)|J=2\rangle. \quad (12)$$

Because $Z^0_b(10610)$ has the structure (6), its radiative decay is possible only through M1 transitions, and (8)-(12) imply that the decay ratio is given as

$$\Gamma(Z^0_b \rightarrow \chi_{b0}\gamma) : \Gamma(Z^0_b \rightarrow \chi_{b1}\gamma) : \Gamma(Z^0_b \rightarrow \chi_{b2}\gamma) = 1 : 3 : 5. \quad (13)$$

Note that the differences of the phase spaces are neglected, and furthermore overlaps of the meson wave functions are assumed to be the same, which is the case when the spatial wave
functions of $\chi_{b,J}$ have the same node quantum number. Considering the phase space factors proportional to the cube of the photon energy $\omega_J$, we find the relation for $Z_b^0(10610) \to \chi_{b,J}(1P)$

$$
\Gamma(Z_b^0 \to \chi_{b0}(1P)\gamma) : \Gamma(Z_b^0 \to \chi_{b1}(1P)\gamma) : \Gamma(Z_b^0 \to \chi_{b2}(1P)\gamma) = 1 : 2.6 : 4.1,
$$

(14)

and for $Z_b^0(10610) \to \chi_{b,J}(2P)$

$$
\Gamma(Z_b^0 \to \chi_{b0}(2P)\gamma) : \Gamma(Z_b^0 \to \chi_{b1}(2P)\gamma) : \Gamma(Z_b^0 \to \chi_{b2}(2P)\gamma) = 1 : 2.5 : 3.8,
$$

(15)

which can be compared with experimental data. So far we have assumed that $Z_b(10610)$ wave function is dominated by $^3S_1$ of $\frac{1}{\sqrt{2}}(B\bar{B}^* - B^*\bar{B})$. It should be noted that a $^3D_1$ component exists with a small fraction in the $\frac{1}{\sqrt{2}}(B\bar{B}^* - B^*\bar{B})$ molecular state. The spin structure of $\frac{1}{\sqrt{2}}(B\bar{B}^* - B^*\bar{B})(^3D_1)$ is $\frac{1}{\sqrt{2}}(0^H \otimes 1^1_1) + \frac{1}{\sqrt{2}}(1^-_H \otimes 1^-_1)|_{J=1}$, which modifies slightly the relations (14) and (15). Moreover, it implies that E2 transition in the processes $Z_b^0 \to \chi_{b,J}\gamma$ must occur mediated by the $D$-wave component. This is also an interesting point to be studied in experiments.

Now, we consider the production and the decay for recently predicted isor triplet $B^{(*)}\bar{B}^{(*)}$ molecular states having positive $G$-parity. These states can be produced by one pion emission in $P$-wave from $\Upsilon(5S)$. We introduce new notations $W_{bJ}^{PC}$ for the lowest state and $W_{bJ}^{PC}$ for the first excited state. The quantum numbers and the main components of these $B^{(*)}\bar{B}^{(*)}$ molecular states are

$$
W_{b0}^- \ 1^+(0^-) : \frac{1}{\sqrt{2}}(B\bar{B}^* + B^*\bar{B})(^3P_0),
$$

(16)

$$
W_{b1}^- \ 1^+(1^-) : \frac{1}{\sqrt{2}}(B\bar{B}^* + B^*\bar{B})(^3P_1),
$$

(17)

$$
W_{b1}^- \ 1^+(1^-) : B\bar{B}(^1P_1),
$$

(18)

$$
W_{b2}^- \ 1^+(2^-) : B^*\bar{B}^* (^5P_2),
$$

(19)

$$
W_{b2}^- \ 1^+(2^-) : \frac{1}{\sqrt{2}}(B\bar{B}^* + B^*\bar{B})(^3P_2).
$$

(20)

We note that, in principle, $B^*\bar{B}^*(^1P_1)$ and $B^*\bar{B}^* (^5P_2)$ in $I^G(J^{PC}) = 1^+(1^-)$ states can be mixed. However, the previous study indicates that $W_{b1}^- \ 1^+(1^-)$ and $W_{b2}^- \ 1^+(2^-)$ are close to each threshold of $B\bar{B}$ and $B\bar{B}^*$, which means $B^*\bar{B}^*$ components are suppressed. There could be also coupled channels $\frac{1}{\sqrt{2}}(B\bar{B}^* + B^*\bar{B})(^3F_2)$ and $B^*\bar{B}^* (^5F_2)$ components in $1^+(2^-)$ states.


However, we expect that the probabilities of these components are small due to high angular momentum. Therefore, from Eqs. (16)-(20) the corresponding spin structures of $W_{bJ}^{-}$ states are given as

$$W_{b0}^{-} : (1_H^0 \otimes 1_l^+)|_{J=0}, \quad (21)$$

$$W_{b1}^{-} : - \frac{1}{\sqrt{3}} (1_H^- \otimes 0_l^+) + \frac{1}{2} (1_H^0 \otimes 1_l^+)|_{J=1} + \frac{\sqrt{5}}{2\sqrt{3}} (1_H^- \otimes 2_l^+)|_{J=1}, \quad (22)$$

$$W_{b2}^{-} : \frac{1}{2} (0_H^- \otimes 1_l^+) + \frac{1}{2\sqrt{3}} (1_H^- \otimes 0_l^+ - \frac{1}{2} (1_H^0 \otimes 1_l^+)|_{J=1} + \frac{\sqrt{5}}{2\sqrt{3}} (1_H^- \otimes 2_l^+)|_{J=1}, \quad (23)$$

$$W_{b2}^{-} : - \frac{1}{2} (1_H^- \otimes 1_l^+)|_{J=2} + \frac{\sqrt{3}}{2} (1_H^- \otimes 2_l^+)|_{J=2}, \quad (24)$$

$$W_{b1}^{-} : - \frac{1}{2} (1_H^- \otimes 1_l^+)|_{J=2} + \frac{\sqrt{3}}{2} (1_H^- \otimes 2_l^+)|_{J=2}. \quad (25)$$

Remarkably, the heavy quark spin singlet state exists only in $W_{b1}^{-}$. Therefore decays of $W_{bJ}^{-}$ into a heavy quark spin singlet state of bottomonium and light hadrons are forbidden except for $W_{b1}^{-}$, although their quantum numbers and kinematics allow them. Only $W_{b1}^{-}$ can decay into $h_1 \pi$ or $\eta_b \rho(\gamma)$. In contrast, all $W_{bJ}^{-}$ can decay into a heavy quark spin triplet state, e.g. $\Upsilon \pi$ in $P$-wave ($|\Upsilon \pi|_{P\text{-wave}} = (1_H^- \otimes 1_l^+)|_{J=0,1,2}$). The decay ratio of the processes $W_{bJ}^{-} \rightarrow \Upsilon \pi$ is obtained as

$$\Gamma(W_{b0}^{-} \rightarrow \Upsilon \pi) : \Gamma(W_{b1}^{-} \rightarrow \Upsilon \pi) : \Gamma(W_{b1}^{-} \rightarrow \Upsilon \pi) : \Gamma(W_{b2}^{-} \rightarrow \Upsilon \pi) : \Gamma(W_{b2}^{-} \rightarrow \Upsilon \pi) = 4 : 1 : 1 : 3 : 1. \quad (26)$$

The decay processes $\Upsilon(5S) \rightarrow W_{bJ}^{-} \rightarrow \Upsilon(nS)\pi\pi$ are not yet observed, though they are similar to $\Upsilon(5S) \rightarrow Z_b \pi \rightarrow \Upsilon(nS)\pi\pi$. However we expect that the production rates of $W_{bJ}^{-}$ is small compared with that of $Z_b$, since the $W_{bJ}^{-}$ transition is mediated by a pion emission in $P$-wave. High statistics and refined analysis of experiments will establish the presence or absence of $W_{bJ}^{-}$ and their production and decay properties. Finally, the production ratio of $W_{bJ}^{-}$ from $\Upsilon(5S)$ is obtained as

$$f(W_{b0}^{-}\pi) : f(W_{b1}^{-}\pi) : f(W_{b1}^{-}\pi) : f(W_{b2}^{-}\pi) : f(W_{b2}^{-}\pi) = 2 : 9 : 4.5 : 9 : 12, \quad (27)$$

where we find the production rate of $W_{b2}^{-}$ is favored, while the production of $W_{b0}^{-}$ is suppressed.

In summary, we have derived the model independent relations among various decay and production rates for possible $B^{(*)}\bar{B}^{(*)}$ molecular states under the heavy quark symmetry.
Part of decay properties of $Z_b(10610)$ and $Z_b(10650)$ are well explained and the possible decay patterns for neutral $Z^0_b(10610)$ are discussed in the present framework. We have shown that the $W_{bJ}^{-}^{-}$ decay into a spin singlet bottomonium is forbidden except for $W_{bJ}^{1-}$. We have also predicted the production rate of various $W_{bJ}^{-}^{-}$ through the one pion emission of $\Upsilon(5S)$. All of them can be tested experimentally and will provide important information to further understand the exotic structure of the new particles.

Acknowledgements

This work is partly supported by the Grant-in-Aid for Scientific Research on Priority Areas titled “Elucidation of New Hadrons with a Variety of Flavors” (E01: 21105006).

[1] I. Adachi [Belle Collaboration], arXiv:1105.4583 [hep-ex].
[2] A. Bondar et al. [Belle Collaboration], Phys. Rev. Lett. 108, 122001 (2012) arXiv:1110.2251 [hep-ex].
[3] I. Adachi et al. [Belle Collaboration], Phys. Rev. Lett. 108, 032001 (2012) arXiv:1103.3419 [hep-ex].
[4] I. Adachi et al. [Belle Collaboration], arXiv:1207.4345 [hep-ex].
[5] A. E. Bondar, A. Garmash, A. I. Milstein, R. Mizuk and M. B. Voloshin, Phys. Rev. D 84, 054010 (2011) arXiv:1105.4473 [hep-ph].
[6] M. B. Voloshin, Phys. Rev. D 84, 031502 (2011) arXiv:1105.5829 [hep-ph].
[7] S. Ohkoda, Y. Yamaguchi, S. Yasui, K. Sudoh and A. Hosaka, Phys. Rev. D 86, 014004 (2012) arXiv:1111.2921 [hep-ph].