How Large is the Contribution of Excited Mesons in Coupled-Channel Effects? Postprint

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Abstract

We study the excited B mesons' contributions to the coupled-channel effects under the framework of the 3P0 model for bottomonium. Contrary to what has been widely accepted, the contributions of P-wave B mesons are generally the largest, and to some extent, this result is independent of the potential parameters. We also push the calculation beyond B(1P) and carefully analyze the contributions of B(2S). A form factor is a key ingredient to suppress the contributions of B(2S) for low-lying bottomonia. However, this suppression mechanism is not efficient for highly excited bottomonia, such as (5S) and (6S). We provide explanations for why this difficulty occurs in the 3P0 model and suggest analyzing the flux-tube breaking model for the full calculation of coupled-channel effects.

Full Text

Preamble

How Large is the Contribution of Excited Mesons in Coupled-Channel Effects?

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Abstract

We investigate the contributions of excited B mesons to coupled-channel effects within the $^3P_0$ model framework for bottomonium. Contrary to widely

held assumptions, we find that P-wave B mesons generally provide the largest contributions, and this result is largely independent of potential parameters to a significant extent. We extend our calculations beyond $B(1P)$ and carefully analyze the contributions from $B(2S)$. A form factor proves essential for suppressing the $B(2S)$ contributions for low-lying bottomonia, yet this suppression mechanism becomes inefficient for highly excited states such as $\Upsilon(5S)$ and $\Upsilon(6S)$. We explain why this difficulty arises in the $^3P_0$ model and suggest that the flux-tube breaking model should be analyzed for a complete calculation of coupled-channel effects.

Introduction

Heavy quarkonium represents a multiscale system encompassing all regimes of quantum chromodynamics (QCD), making it an ideal laboratory for studying strong interactions [**?**]. Despite QCD's success in the high-energy region due to asymptotic freedom, nonperturbative effects dominate at low energies and pose challenges for perturbative calculations. Lattice QCD serves as one important tool for investigating these nonperturbative effects, yet its computational demands remain prohibitive for calculating all physical quantities with current computing power. Consequently, various phenomenological models have been developed, among which the quark model stands out as particularly prominent. Within this framework, numerous groups have proposed different types of interactions, achieving many impressive successes (see e.g., Refs. \cite{2–6}). However, these potential models do not tell the complete story. A crucial missing ingredient is the mechanism for generating quark-antiquark pairs, which expands the Fock space of the initial state—in other words, the initial state contains multiquark components.

These multiquark components modify the Hamiltonian of the potential model, causing mass shifts and mixing between states with identical quantum numbers, or directly contributing to open-channel strong decays when the initial state lies above the corresponding threshold. These phenomena are collectively known as unquenched effects or coupled-channel effects. Through various approaches—including the $^3P_0$ model \cite{7–11}, flux-tube breaking model \cite{12–16}, microscopic decay models \cite{4,17–19}, and S-matrix analysis \cite{20–23}—coupled-channel effects have been extensively studied in numerous works (e.g., Refs. \cite{24–29}).

Despite these pioneering efforts, we identify at least two factors that may jeopardize calculations of coupled-channel effects. The first is the widely used simple harmonic oscillator (SHO) wavefunction, which approximates the realistic wavefunction in overlap integrals. As we previously demonstrated [**?**], the SHO approximation proves inadequate, particularly for states near thresholds, making precise treatment of wavefunctions essential. The second factor is the assumption that contributions from excited meson loops are negligible. Using bottomonium—the system studied in this paper—as an example, two reasons may explain why this approximation is common. First, the thresholds corresponding to excited B mesons are higher than the ground-state $B\bar{B}$ threshold, leading to expectations of small coupled-channel effects. Second, calculating the contributions of excited B mesons is more complicated, and an effective method for performing such calculations is not widely known.

As explained in Ref. [?], the Gaussian expansion method combined with techniques for transforming between Cartesian and spherical bases can, in principle, handle these sophisticated calculations. Another important observation is that the first excited P-wave B meson is only about 450 MeV heavier than the ground-state B meson, a mass difference less than one-tenth of the B meson mass. Given that the quantum numbers of B_1 differ from those of ground-state B mesons and that coupled-channel effects depend critically on quantum numbers, the suppression arising purely from the larger mass may not be as substantial as commonly assumed.

Summing all intermediate meson loops presents more than just a computational challenge—it may also reveal profound physics. In the light sector, a series of works by Geiger et al. \cite{31–33} demonstrated that although different intermediate meson loops contribute to breaking the Okubo-Zweig-Iizuka (OZI) rule, under certain simplifications (such as neglecting mass differences in denominators), they contribute destructively, preserving the OZI rule perfectly. Summing only a subset of meson loops could lead to incorrect conclusions. However, one should not confuse their calculations with our work. They studied flavor-changing processes such as $u\bar{u} \to$ virtual meson pairs $\to d\bar{d}$, whereas we focus on mass shifts, and such cancellation does not occur in our case. For states below threshold, intermediate loops always contribute a negative mass shift, and even for above-threshold cases where mass shifts may add destructively, mass differences matter in realistic calculations.

In Ref. [?], Geiger and Isgur showed that under the flux-tube breaking model, when summing all intermediate states, the mass shift caused by coupled-channel loops fails to converge under two assumptions: (i) the meson wavefunction is the SHO wavefunction, and (ii) the string length between generated quark pairs is zero. In the $^3P_0$ model, this conclusion requires verification or recalculation because the SHO approximation is far from realistic, as previously emphasized, and the vertex for quark generation differs from that in the flux-tube breaking model.

Some studies of excited meson loops for heavy quarkonium exist. In the charmonium sector, a lattice calculation in Ref. [?] revealed that η_c and J/ψ have small but non-negligible $D_1\bar{D}^*$ components. In the bottomonium sector, loops involving B_1 mesons prove critical for explaining the large breaking of heavy quark spin symmetry in $\Upsilon(10860)$ [?]. However, these studies focus on specific states, and systematic investigations of excited meson loops remain lacking.

In summary, two questions require clarification: whether the ground-state approximation is valid, and what are the general properties of excited meson loops. This work addresses these questions within the $^3P_0$ model framework.

This paper is organized as follows. Section 2 briefly describes the ingredients of coupled-channel effects and calculation techniques, including the $^3P_0$ model, Cornell potential model, Gaussian expansion method, and transformations between spherical and Cartesian bases. Section 3 presents results for higher excited B mesons up to $B(2S)$, explaining the necessity of form factors and the limitations of the $^3P_0$ model. Finally, Section 4 provides a brief summary.

2 Theoretical Framework

The calculation methods and tools are described in detail in Ref. [**?**]; therefore, we only outline the key steps in this section.

[Figure 1: see original paper] Sketch of coupled-channel effects in the $^3P_0$ model for bottomonium. i and f denote the initial and final states with the same J^{PC}, and $B\bar{B}$ represents all possible B meson pairs, including excited B mesons.

The coupled-channel effects in the $^3P_0$ model \cite{7–9} are illustrated in Fig. 1. In this model, the generated quark pairs carry vacuum quantum numbers $J^{PC} = 0^{++}$, corresponding to $^3P_0$ in spectroscopic notation, which explains the model's name.

The Hamiltonian for generating quark-antiquark pairs is expressed as

$$H_I = 2m_q\gamma \int d^3x \bar{\psi}_q \psi_q,$$

where m_q is the mass of the produced quark and γ is a dimensionless coupling constant. Since the probability of generating heavier quarks is suppressed, we use the effective strength $\gamma_s = m_q\gamma$ in our calculations, where $m_q = m_u = m_d$ is the constituent quark mass of the up (or down) quark and m_s is the strange quark mass.

As sketched in Fig. 1, the experimentally observed state should be a mixture of a pure quarkonium state $|\psi_0\rangle$ and B meson continuum states $|BC; p\rangle$. The observed state $|A\rangle$ can be expressed as

$$|A\rangle = c_0|\psi_0\rangle + \sum \int d^3p \, c_{BC}(p)|BC; p\rangle,$$

which is an eigenstate of the full Hamiltonian H defined in Eq. (6). The coefficients c_0 and c_{BC} represent normalization constants for the bare state and B meson continuum, respectively.

To determine $|A\rangle$, one must first solve for the heavy quarkonium wavefunction $|\psi_0\rangle$. In this work, we obtain it by solving the Schrödinger equation

$$\left(H_0|\psi_0\rangle = 2m_b + \frac{p^2}{m_b} + V(r) + V_s(r)\right)|\psi_0\rangle = M_0|\psi_0\rangle,$$

where m_b and M_0 represent the b quark mass and the bare mass of bottomonium, respectively. In the above equation, $V(r)$ is the well-known Cornell potential [?, ?]

$$V(r) = -\frac{4\alpha_s}{3r} + \lambda r + c,$$

with α_s, λ, and c representing the strength of the color Coulomb potential, linear confinement, and mass renormalization, respectively. $V_s(r)$ denotes spin-dependent interactions that restore the hyperfine and fine structures of bottomonium:

$$V_s(r) = \left(\frac{2\alpha_s}{m_b^2 r^3} - \frac{3\lambda}{2m_b^2 r}\right)\vec{L}\cdot\vec{S} + \tilde{\delta}(r)\vec{S}_b\cdot\vec{S}_{\bar{b}} + \frac{1}{m_b^2}\left[\vec{S}_b\cdot\vec{S}_{\bar{b}} - \frac{3(\vec{S}_b\cdot\vec{r})(\vec{S}_{\bar{b}}\cdot\vec{r})}{r^2}\right],$$

where $\vec{L}$ denotes the relative orbital angular momentum, $\vec{S} = \vec{S}_b + \vec{S}_{\bar{b}}$ is the total spin of the b quark pair, and m_b is the b quark mass. The smeared delta function $\tilde{\delta}(r)$ is written as $\tilde{\delta}(r) = (\sigma/\sqrt{\pi})^3 e^{-\sigma^2 r^2}$ [?, ?].

We treat the spin-dependent term as a perturbation and obtain spatial wavefunctions by numerically solving the Schrödinger equation using Numerov's method [?].

Combining the Cornell potential with the dynamics of quark pair generation yields the full Hamiltonian

$$H = H_0 + E_{BC} + H_I,$$

where $E_{BC} = \sqrt{m_B^2 + p^2} + \sqrt{m_C^2 + p^2}$. As the name suggests, this term is naturally defined as the energy of the B meson pair. The H_I term in Eq. (6) is primarily responsible for the mass shift ΔM, which can be obtained by solving the integral equation

$$\Delta M \equiv M - M_0 = \sum\int \frac{|\langle BC; p|H_I|\psi_0\rangle|^2}{M - E_{BC} - i\epsilon}.$$

Note that the $i\epsilon$ term is included to handle cases where $m_A > m_B + m_C$. In such situations, ΔM acquires an imaginary part

$$\mathrm{Im}(\Delta M) = \sum |\langle BC; P_B|H_I|\psi_0\rangle|^2,$$

which equals one-half of the decay width. P_B and E_B denote the momentum and energy of the B meson, respectively.

The wavefunction overlap integral appears in the term

$$\langle BC; P_B|H_I|\psi_0\rangle = \sum_{\text{polarization}} \int d^3k\, \phi_0(\vec{k}+\vec{P}_B)\phi_B^*(\vec{k}+x\vec{P}_B)\phi_C^*(\vec{k}+x\vec{P}_B)|\vec{k}|Y_1^m(\theta_{\vec{k}}, \phi_{\vec{k}}),$$

where $x = m_q/(m_b + m_q)$, and m_b and m_q denote the b quark and light quark masses, respectively.

Once M is determined, the coefficients of different components can also be calculated. For states below threshold, the probability of the $b\bar{b}$ component is expressed as

$$P_{b\bar{b}} := |c_0|^2 = \left[1 + \sum \int \frac{p^2|MLS|^2}{(M - E_{BC})^2}\right]^{-1},$$

where $|MLS|^2$ is represented as

$$|MLS|^2 = \int d\Omega_B |\langle BC; P_B|H_I|\psi_0\rangle|^2.$$

The main computational effort lies in evaluating Eq. (10). To calculate it precisely, we employ the Gaussian expansion method (GEM) [**?**] and transformations between spherical and Cartesian bases to make GEM automatically applicable to excited B mesons.

3.1 B_1's Contributions with Traditional $^3P_0$ Model

In this work, we do not aim to reproduce the bottomonium spectrum and decay widths but rather to study the general contributions of excited B mesons to coupled-channel effects. Our calculation involves two simplifications. First, since the constituent quark masses of u and d quarks are set equal for obtaining wavefunctions, we further approximate $m(B^0) \approx m(B^\pm)$. Despite this simplification, 42 channels remain to be calculated: $B^{(*)}\bar{B}_s^{(1P)}$ and $\bar{B}^{(s)}(1P)\bar{B}^{(s)}(1P)$. In principle, one could treat their masses precisely; however, based on our experience, this simplification is quite accurate and saves considerable computational effort.

The second simplification concerns B_1 meson multiplets. From the quark model perspective, four $1P$-wave B mesons exist. However, the predicted $B_0(1P)$ and one of the $B_1(1P)$ states remain unobserved experimentally. For coupled-channel calculations, the absence of these two states poses no real problem. Because the b quark mass is very large, heavy quark spin symmetry is expected to hold approximately. Under this limit, it is reasonable to set the masses of $B_0(^3P_0)$ and the missing partner of $B_1(5721)$ equal to that of $B_1(5721)$.

Mixing occurs between $B(^3P_1)$ and $B(^1P_1)$ to form the experimentally observed $B_1(5721)$ [?]. However, since we sum over all possible B meson combinations, mixing between the two 1^+ B mesons can be neglected. For intermediate states with well-measured masses, we use values from the Particle Data Group [?].

The coupled channels are considerably more complex than cases including only ground-state B mesons. For illustration, we sum contributions from states within the same multiplet and flavor (u, d, s) along with their charge-conjugate partners, classifying channels into three groups (see Fig. 2 [Figure 2: see original paper]). For example, $B(1S)$ denotes ground-state B, B^*, B_s, and B_s^* mesons. Throughout this paper, we define $\Delta M(i, j)$ to represent the sum of mass shifts due to all possible combinations of i- and j-wave B mesons.

The parameters are listed in Table 1 , identical to those in our previous work [?].

Table 1: Parameters used in our calculation.

Parameter	Value
α	0.34
λ	0.22 GeV^2
c	0.435 GeV
m_b	4.5 GeV
$m_u = m_d$	0.33 GeV
m_s	0.5 GeV
γ	0.205
σ	3.838 GeV

The mass shifts ΔM for both $B(1S)$ and $B(1P)$ mesons are depicted in Fig. 2.

In the remainder of this section, we discuss several interesting features of these plots. When states lie far below threshold, the mass shift is always negative [as indicated in Eq. (8)], with $|\Delta M|$ increasing as the threshold is approached. At higher masses, the magnitude of ΔM no longer increases monotonically but oscillates, reflecting the nodal structure of the wavefunctions. This phenomenon occurs both when B_1 meson contributions are included and when they are omitted.

A specific example clarifies this point. For $\Upsilon(10860)$, which we treat as $\Upsilon(5S)$, $\Delta M(1S, 1P)$ is generally the largest, but exhibits a significant dip around 11.14 GeV with the parameters in Table 1. However, using the parameters of Ref. [?], this channel's contribution always increases. In such cases, drawing firm conclusions about spectral behavior becomes difficult—the spectrum is sensitive to potential parameters, a conclusion that agrees with Ref. [?].

The most notable impact of the B_1 family is its unexpectedly large contributions to the mass shift; generally, $\Delta M(1S, 1P)$ is the largest. Even for $\Upsilon(1S)$, which lies far below threshold, $\Delta M(1S, 1P) + \Delta M(1P, 1P)$ is 14 times larger than $\Delta M(1S, 1S)$. Compared with calculations considering only ground-state B mesons, threshold effects become more clearly resolved when excited B mesons are included because open channels from different multiplets are well separated.

The first open-bottom threshold is $2m_B \approx 10.56$ GeV. As the mass approaches this value, the denominator $M - E_{BC}$ in Eq. (8) approaches zero, causing enhancement of ΔM from the $B(1S)\bar{B}(1S)$ channel. In the $\Upsilon(3S)$ case shown in Fig. 2, one can clearly observe the sharp increase in $\Delta M(1S, 1S)$ between 10.55 and 10.6 GeV, while the slopes of other channels change little because their thresholds are 450 MeV higher.

Threshold effects are partially mitigated by wavefunction nodes. For higher excited states, peaks and valleys in the wavefunction are more likely to cancel each other, leaving a relatively small slope for ΔM. Comparing the results for $\Upsilon(4S)$ and $\Upsilon(3S)$ verifies this conclusion.

We emphasize that the $^3S_1$ family is not the only one that couples strongly to $B\bar{B}(1P)$ loops. As shown in Fig. 3 [Figure 3: see original paper], all families—$^1S_0, ^1P_1, ^3P_J$, and $^3D_1$—share these general properties.

A direct consequence of the large $\Delta M(1S, 1P)$ is that parameters fitted considering only ground-state B meson contributions are incomplete or even misleading. This conclusion is largely parameter-independent, as both parameter sets (from Refs. [?] and [?]) yield large $\Delta M(1S, 1P)$.

It might seem feasible to fit the spectrum by tuning parameters such as the mass renormalization term c in Eq. (4) to absorb this unexpectedly large mass shift, or alternatively, by weakening the confinement potential to reflect the general suppression of mass near thresholds. This could be a viable approach if one focuses solely on the spectrum. However, as pointed out in our previous work [?], the mass shift reveals only one aspect of coupled-channel effects. Specifically, the large $B(1P)$ contribution brings not only a substantial mass shift ΔM but also a large meson-pair fraction (or equivalently, a small $P_{b\bar{b}}$). This wavefunction renormalization cannot be accommodated within the potential model framework.

3.2 Beyond B_1

The large B_1 contributions underscore the necessity of performing ab initio calculations of coupled-channel effects. However, before undertaking such calculations, one naturally asks: "How large are the contributions from states beyond B_1, and how should one evaluate all intermediate meson loops?" Providing a complete answer is extremely difficult. Nevertheless, we can estimate contributions up to $B(2S)$ and draw some model-independent conclusions through careful analysis of mass shift behavior.

First, we attempt to precisely evaluate these loops. Since the $B(2S)$ mass and wavefunctions are unknown, the calculation inevitably becomes model-dependent. To date, the $B(2S)$ meson has not been experimentally well determined, and its mass remains model-dependent. For our purposes, it suffices to adopt an average of several theoretical estimates from Refs. \cite{42–46}, while $B(2S)$ wavefunctions are still derived from our parameters for consistency.

We define the ratio

$$R = \frac{\Delta M(2S, 2S) + \Delta M(2S, 1P) + \Delta M(2S, 1S)}{\Delta M(1S, 1S) + \Delta M(1S, 1P) + \Delta M(1P, 1P)}$$

to quantify the magnitude of $B(2S)$ contributions. As explained below, when $B(2S)$ states are involved, we adopt a Gaussian form factor that characterizes the size of the generated quark pairs (see, e.g., Refs. \cite{47–49}). This form factor modifies the matrix element to

$$\langle BC; P_B|H_I|\psi_0\rangle = \sum_{\text{polarization}} \int d^3k\, e^{-2r^2k^2/3} \phi_0(\vec{k}+\vec{P}_B)\phi_B^*(\vec{k}+x\vec{P}_B)\phi_C^*(\vec{k}+x\vec{P}_B)|\vec{k}|Y_1^m(\theta_{\vec{k}}, \phi_{\vec{k}}).$$

In this work, we fit $r = 0.408$ fm, which minimizes R for $\Upsilon(1S)$. This value is slightly larger than $r = 0.335$ fm used in Ref. [**?**].

Fig. 4 [Figure 4: see original paper] clearly demonstrates that the form factor is crucial for suppressing $B(2S)$ contributions. Without it, $B(2S)$ can contribute an additional 25% to ΔM—a non-negligible deviation for precise spectral fits. However, with a form factor of $r = 0.408$ fm, the ΔM from $B(2S)$ is suppressed to less than 5%. This partially answers our earlier question: a form factor is indispensable for suppressing $B(2S)$ contributions. However, as n increases in the $\Upsilon(nS)$ family, the form factor becomes increasingly inefficient at suppressing R.

Another effect of the form factor is that it introduces additional peaks in the ΔM oscillations. This has a straightforward explanation: since the form factor effectively acts as a cutoff, it generally suppresses every channel's contribution, but different channels are suppressed by different magnitudes. As the number of nodes in the wavefunctions of bottomonia and B mesons increases, the relative contributions from $B(2S)$ may increase at specific energy points.

The sharp peaks for $\Upsilon(5S)$ and $\Upsilon(6S)$ deserve special analysis. Both occur around 11.18 GeV. Given that the wavefunctions of $\Upsilon(5S)$ and $\Upsilon(6S)$ are quite different, their sharing the same sharp peaks at identical energy points appears more than coincidental.

While R conveniently estimates $B(2S)$ relative contributions, it cannot distinguish whether peaks arise from denominator suppression or numerator enhancement. The complete information is encoded in ΔM itself. Fig. 5 [Figure 5: see original paper] shows ΔM values for $\Upsilon(5S)$ and $\Upsilon(6S)$ with different form factors.

Fig. 5 reveals that the enhancement of R around 11.18 GeV actually results from greater suppression of $\Delta M(1S,1S)+\Delta M(1S,1P)+\Delta M(1P,1P)$ [or more precisely, suppression of $\Delta M(1S,1P)$]. We have scanned r over a wide range (0.17–1.08 fm), and all cases show that the form factor does not shift the positions of ΔM peaks or valleys. The form factor's failure to suppress $B(2S)$ contributions around this region is somewhat coincidental, as this suppression does not occur for $\Upsilon(5S)$ and $\Upsilon(6S)$ using the parameters of Ref. [?], where $\Delta M(1S,1P)$ always increases.

Even in the low-energy region where all coupled channels contribute negative ΔM, $B(2S)$ contributes approximately 20% of the total mass shift for $\Upsilon(5S)$ and $\Upsilon(6S)$, regardless of whether the form factor is included. In other words, the form factor fails to suppress $B(2S)$ contributions for these highly excited states. This strongly indicates that merely adding a Gaussian form factor is inadequate for complete coupled-channel calculations within the $^3P_0$ framework.

One might argue that coupled-channel effects could be absorbed into a smeared potential that produces broader wavefunctions than the Cornell model. Such broad wavefunctions would indeed somewhat suppress higher excited-state contributions. However, based on our experience, modifying wavefunction sizes also cannot resolve the convergence issue.

We conclude this section with comments on the $^3P_0$ model. In the classical $^3P_0$ model, suppression from the natural cutoff of wavefunctions and the increase in the denominator of Eq. (8) are too weak. Even with the remedy of a Gaussian form factor, unless the cutoff r is chosen carefully, one may enhance rather than suppress the $B(2S)$ contribution. Compared with other models such as microscopic decay models or the flux-tube breaking model—which contain rich microscopic details of decay vertices—the $^3P_0$ model replaces these fine structures with an overall coupling constant γ. This approximation may be inappropriate (see, e.g., Ref. [?]). In our case, we demonstrate that this approximation leads to poor convergence of the sum over excited meson loops for highly excited bottomonia.

The rich structure of quark pair generation vertices may help resolve the convergence issue. For example, the Hamiltonian of the flux-tube breaking model suppresses the generation of distant quarks, making it more difficult to produce excited intermediate states.

[Figure 6: see original paper] Sketch of coupled-channel effects in the flux-tube breaking model. Red lines represent the flux tube. Regions outside the flux tube are shaved off, i.e., excited-state contributions in loops are suppressed.

As illustrated in Fig. 6, contributions from regions farther from the flux-tube line can be dropped, naturally suppressing excited meson contributions. Without such dynamical suppression, it becomes more difficult to suppress $B(2S)$ contributions even with modified versions of the $^3P_0$ model where form factors are added.

4 Summary and Outlook

In this paper, within the $^3P_0$ model framework, we explicitly calculate the contributions of excited B mesons to coupled-channel effects in bottomonium. We reveal that $B(1P)$ meson contributions are generally the largest compared to ground-state B mesons. Up to this partial wave, ab initio calculations of coupled-channel effects are necessary.

When we extend calculations beyond $B(1P)$, we encounter fundamental difficulties with the $^3P_0$ model. Even with carefully chosen form factors, the model cannot efficiently suppress contributions from higher partial-wave intermediate states. Since we do not fit our parameters to experimental data and have employed several different parameter sets, we have strong reason to believe these difficulties are independent of wavefunctions or potential models.

We suggest that an efficient suppression mechanism, such as dynamical suppression, is needed to properly evaluate coupled-channel effects. How to effectively sum over all intermediate loops remains an open issue.

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Note: Figure translations are in progress. See original paper for figures.