Z$_2$ spin liquids in $S = 1/2$ Heisenberg model on kagome lattice: projective symmetry group study of Schwinger-fermion mean-field states

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(Dated: January 19, 2013)

Due to strong geometric frustration and quantum fluctuation, $S = 1/2$ quantum Heisenberg antiferromagnets on the kagome lattice has long been considered as an ideal platform to realize spin liquid (SL), a novel phase exhibiting fractionalized excitations without any symmetry breaking. A recent numerical study$^1$ of Heisenberg $S = 1/2$ kagome lattice model (HKLM) shows that in contrast to earlier results, the ground state is a singlet-gapped SL with signatures of $Z_2$ topological order. Motivated by this numerical discovery, we use projective symmetry group to classify all 20 possible Schwinger-fermion mean-field states of $Z_2$ SLs on kagome lattice. Among them we found only one gapped $Z_2$ SL (which we call $Z_2[0, \pi|\beta]$ state) in the neighborhood of $U(1)$-Dirac SL state. Since its parent state, i.e. $U(1)$-Dirac SL is found$^2$ to be the lowest among many other candidate $U(1)$ SLs including the uniform resonating-valence-bond states, we propose this $Z_2[0, \pi|\beta]$ state to be the numerically discovered SL ground state of HKLM.

PACS numbers: 71.27.+a, 75.10.Kt

I. INTRODUCTION

At zero temperature all degrees of freedom tend to freeze and usually a variety of different orders, such as superconductivity and magnetism, will develop in different materials. However, in a quantum system with a large zero-point energy, one may expect a liquid-like ground state to exist even at $T = 0$. In a system consisting of localized quantum magnets, we call such a quantum-fluctuation-driven disordered ground state a quantum spin liquid (SL). It is an exotic phase with novel “fractionalized” excitations carrying only a fraction of the electron quantum number, e.g. spinons which carry spin but no charge. The internal structures of these SLs are so rich that they are beyond the description of Landau’s symmetry breaking theory of conventional ordered phases. Instead they are characterized by long-range quantum entanglement encoded in the ground state, which is coined “topological order” in contrast to the conventional symmetry-breaking order.

Geometric frustration in a system of quantum magnets would lead to a huge degeneracy of classical ground state configurations. The quantum tunneling among these classical ground states provides a mechanism to realize quantum SLs. The quest for quantum SLs in frustrated magnets (for a recent review see Ref.$^3$) has been pursued for decades. Among them the Heisenberg $S = 1/2$ kagome lattice model (HKLM)

$$H_{HKLM} = J \sum_{<i,j>} S_i \cdot S_j$$

(1)

has long been thought as a promising candidate. Here $<i,j>$ denotes $i,j$ being a nearest neighbor pair. Experimental evidence of SL$^{10-13}$ has been observed in ZnCu$_3$(OH)$_6$Cl$_2$ (called herbertsmithite), a spin-half antiferromagnet on the kagome lattice. Theoretically, in lack of an exact solution of the two-dimensional (2D) quantum Hamiltonian$^1$ in the thermodynamic limit, in previous studies either a honeycomb valence bond crystal$^{14-15}$ (HVBC) with an enlarged 6 x 6-site unit cell, or a gapless SL$^{19}$ were proposed as the ground state of HKLM. However, recently an extensive density-matrix-renormalization-group (DMRG) study$^{2}$ on HKLM reveals the ground state of HKLM as a gapped SL, which substantially lowers the energy compared to HVBC. Besides, they also observe numerical signatures of $Z_2$ topological order in the SL state.

Motivated by this important numerical discovery, we try to find out the nature of this gapped $Z_2$ SL. Different $Z_2$ SLs on the kagome lattice have been previously studied using Schwinger-boson representation.$^{20,22}$ Here we propose the candidate states of symmetric $Z_2$ SLs on kagome lattice by Schwinger-fermion mean field approach.$^{22-28}$ Following is the summary of our results. First we use projective symmetry group ($^{6}$PSG) to classify all 20 possible Schwinger-fermion mean-field ansatz of $Z_2$ SLs which preserve all the symmetry of HKLM, as shown in TABLE I. We analyze these 20 states and rule out some obviously unfavorable states: e.g. gapless states, and those states whose 1st nearest neighbor (n.n.) mean-field amplitudes must vanish due to symmetry. Then we focus on those $Z_2$ SLs in the neighborhood of the $U(1)$-Dirac SL$^2$. In Ref.$^2$ it is shown that $U(1)$-Dirac SL has a significantly lower energy compared with other candidate $U(1)$ SL states, such as the uniform resonating-valence-bond (RVB) state (or the $U(1)$ SL-[0,0] state in notation of Ref.$^2$). We find out that there is only one gapped $Z_2$ SL, which we label as $Z_2[0, \pi|\beta]$ in the neighborhood of (or continuously connected to) $U(1)$-Dirac SL. Therefore we propose this $Z_2[0, \pi|\beta]$ state as a promising candidate state for the ground state of HKLM. The mean-field ansatz of $Z_2[0, \pi|\beta]$ state is shown in FIG. (1b). Our work also provides guideline for choosing variational
FIG. 1: (color online) (a) kagome lattice and the elements of its symmetry group. $\vec{a}_{1,2}$ are the translation unit vectors, $C_6$ denotes $\pi/3$ rotation around honeycomb center and $\sigma$ represents mirror reflection along the dashed blue line. Here $u_a$ and $u_\beta$ denote 1st and 2nd nearest neighbor (n.n.) mean-field bonds while $u_\alpha$ and $u_\bar{\alpha}$ represent two kinds of independent 3rd n.n. mean-field bonds. (b) Mean-field ansatz of $Z_2[0,\pi]\beta$ state up to 2nd nearest neighbor. Colors in general denote the sign structure of mean-field solutions describe SL states if lattice symmetry is preserved. By construction the mean-field ansatz $\psi_i \rightarrow W_i \psi_i, U_{ij} \rightarrow W_i U_{ij} W_j, W_i \in SU(2)$ leaves the action invariant. This redundancy is originated from representation Eq.(2): this local $SU(2)$ transformation leaves the spin operators invariant and does not change physical Hilbert space. One can try to solve Eq.(1) by mean-field (or saddle-point) approximation. At mean-field level, $U_{ij}$ and $a'_0(i)$ must be chosen such that constraints Eq.(3) are satisfied at the mean field level: $\langle \psi_i^\dagger \tau^\mu \psi_i \rangle = 0$. The mean-field ansatz can be written as:

$$H_{MF} = -\sum_{<ij>} \psi_i^\dagger (i|j) \psi_j + \sum_{i} \psi_i^\dagger a'_0(i) \tau^\mu \psi_i.$$  

where two-component fermion notation $\psi_i \equiv (f_{i\uparrow}, f_{i\downarrow})^\top$ is introduced for reasons that will be explained shortly. We use $\tau^0$ to denote the $2 \times 2$ identity matrix and $\tau^{1,2,3}$ are the three Pauli matrices. $U_{ij}$ is a matrix of mean-field amplitudes:

$$U_{ij} = \left( \frac{\chi_{ij}^\dagger}{\Delta_{ij}}, \frac{\Delta_{ij}}{-\chi_{ij}} \right).$$

$\alpha'_0(i)$ are the local lagrangian multipliers that enforce the constraints Eq.(3).

In terms of $\psi$, Schwinger-fermion representation has an explicit $SU(2)$ gauge redundancy: a transformation $\psi_i \rightarrow W_i \psi_i, U_{ij} \rightarrow W_i U_{ij} W_j, W_i \in SU(2)$ leaves the action invariant. This redundancy is originated from representation Eq.(2): this local $SU(2)$ transformation leaves the spin operators invariant and does not change physical Hilbert space. One can try to solve Eq.(1) by mean-field (or saddle-point) approximation. At mean-field level, $U_{ij}$ and $a'_0(i)$ must be chosen such that constraints Eq.(3) are satisfied at the mean field level: $\langle \psi_i^\dagger \tau^\mu \psi_i \rangle = 0$. The mean-field ansatz can be written as:

$$H_{MF} = -\sum_{<ij>} \psi_i^\dagger (i|j) \psi_j + \sum_{i} \psi_i^\dagger a'_0(i) \tau^\mu \psi_i.$$  

where we defined $(i|j) \equiv \frac{3}{2} J_{ij} U_{ij}$. Under a local $SU(2)$ gauge transformation $(i|j) \rightarrow W_i (i|j) W_j^\dagger$, but the physical spin state described by the mean-field ansatz $(i|j)$ remains the same. By construction the mean-field ansatz does not break spin rotation symmetry, and the mean field solutions describe SL states if lattice symmetry is preserved. Different $(i|j)$ ansatz may be in different SL phases. The mathematical language to classify different SL phases is projective symmetry group (PSG).

B. Projective symmetry group (PSG) classification of topological orders in spin liquids

PSG characterizes the topological order in Schwinger-fermion representation: SLs described by different PSGs are different phases. It is defined as the collection of all combinations of symmetry group and $SU(2)$ gauge transformations that leave mean-field ansatz $(i|j)$ invariant (as $a'_0(i)$ are determined self-consistently by $(i|j)$, these transformations also leave $a'_0(i)$ invariant). The invariance under global $SU(2)$ spin rotations. After a Hubbard-Stratonovich transformation, the lagrangian of the spin system can be written as

$$L = \sum_i \psi_i^\dagger \partial_\tau \psi_i + \sum_{<ij>} \frac{3}{8} J_{ij} \left[ \frac{1}{2} \text{Tr}(U_{ij}^\dagger U_{ij}) - (\psi_i^\dagger U_{ij} \psi_j + h.c.) \right] + \sum_i a'_0(i) \psi_i^\dagger \tau^\mu \psi_i.$$  

where $a'_0(i)$ are the local lagrangian multipliers that enforce the constraints Eq.(3).
of a mean-field ansatz \{\langle ij \rangle \} under an element of PSG \(G_U\) can be written as

\[
G_U \langle \langle ij \rangle \rangle = \langle \langle ij \rangle \rangle, \\
U(\langle \langle ij \rangle \rangle) = \langle \langle U^{-1}(i)U^{-1}(j) \rangle \rangle, \\
G_U(\langle \langle ij \rangle \rangle) = \langle \langle G_U(i)ij G_U(j) \rangle \rangle, \\
G_U(i) \in SU(2).
\]

Here \(U \in SG\) is an element of symmetry group (SG) of the corresponding SL. In our case of symmetric SLs on the kagome lattice, we use \((x,y,s)\) to label a site with sublattice index \(s = u,v,w\) and \(x, y \in \mathbb{Z}\). Bravais unit vector are chosen as \(\vec{a}_1 = a\hat{x}\) and \(\vec{a}_2 = \frac{\sqrt{3}}{2}(\hat{x} + \sqrt{3}\hat{y})\) as shown in FIG. Ia. The symmetry group is generated by time reversal operation \(T\), lattice translations \(T_{1,2}\) along \(\vec{a}_{1,2}\) directions, \(\pi/3\) rotation \(G_6\) around honeycomb plaquette center and the mirror reflection \(\sigma\) (for details see Appendix [A]). For example, if \(U = T_1\) is the translation along \(\vec{a}_1\)-direction in FIG.Ia, \(T_1\{(x,y,s)\} = \{x+1, y, s\}\). \(G_U\) is the gauge transformation associated with \(U\) such that \(G_U U\) leave \(\{\langle ij \rangle \}\) invariant. Notice this condition (7) allows us to generate all symmetry-related mean-field bonds from one by the following relation:

\[
\langle ij \rangle = G_U(i)(U^{-1}(i)U^{-1}(j))G_U^\dagger(j)
\]

There is an important subgroup of PSG, the invariant gauge group (IGG), which is composed of all the pure gauge transformations in PSG: \(IGG \equiv \{\{W_i\}|W_i(ij)W_i^\dagger = \langle ij \rangle, W_i \in SU(2)\}\). In other words, \(W_i = G_e(i)\) is the pure gauge transformation associated with identity element \(e \in SG\) of the symmetry group. One can always choose a gauge in which the elements in IGG is site-independent. In this gauge, IGG can be the global \(Z_2\) transformations: \(\{G_e(i) \equiv G_e = \pm 1\}\), the local \(U(1)\) transformations: \(\{G_e(i) \equiv e^{i\theta_j}, \theta \in [0, 2\pi]\}\), or the global \(SU(2)\) transformations: \(\{G_e(i) \equiv e^{i\theta_{ij}}, \theta \in (0, 2\pi), \vec{n} \in S^2\}\), and we term them as \(Z_2\), \(U(1)\) and \(SU(2)\) state respectively.

The importance of IGG is that it controls the low energy gauge fluctuations of the corresponding SL states. Beyond mean-field level, fluctuations of \(\langle ij \rangle\) and \(\theta_{ij}\) need to be considered and the mean-field state may or may not be stable. The low energy effective theory is described by fermionic spinon band structure coupled with a dynamical gauge field of IGG. For example, \(Z_2\) state with gapped spinon dispersion can be a stable phase because the low energy \(Z_2\) dynamical gauge field can be in the deconfined phase. Notice that the condition \(\{G_e(i) \equiv G_e = \pm 1\}\) for a \(Z_2\) SL leads to a series of consistent conditions for the gauge transformations \(\{G_U(i)|U \in SG\}\), as shown in Appendix [A]. Gauge inequivalent solutions of these conditions (A4)-(A11) lead to different \(Z_2\) SLs. Soon we will show that there are 20 \(Z_2\) SLs on the kagome lattice that can be realized by a Schwinger-fermion mean-field ansatz \(\{\langle ij \rangle\}\).

### III. \(Z_2\) Spin Liquids on the Kagome Lattice and \(Z_2[0, \pi/\beta]\) State

Following previous discussions, we use PSG to classify all possible 20 \(Z_2\) SL states on kagome lattice in this section. As will be shown later, among them there is one gapped \(Z_2\) SL labeled as \(Z_2[0, \pi/\beta]\) state in the neighborhood of \(U(1)\)-Dirac SL. This \(Z_2[0, \pi/\beta]\) SL state is the most promising candidate for the SL ground state of HKLM.

#### A. PSG classification of \(Z_2\) spin liquids on kagome lattice

Applying the condition \(G_e(i) \equiv G_e = \pm r^0\) to kagome lattice with symmetry group described in Appendix [A], we obtain a series of consistent conditions for the gauge transformation \(G_U(i)\), i.e. conditions (A4)-(A11). Solving these conditions we classify all the 20 different Schwinger-fermion mean-field states of \(Z_2\) SLs on kagome lattice, as summarized in TABLE II. These 20 mean-field states correspond to different \(Z_2\) SL phases, which cannot be continuously tuned into each other without a phase transition.

As discussed in Appendix [B], from PSG elements \(G_U(i)\) one can obtain all other symmetry-related mean-field bonds from one using symmetry condition [5]. Therefore we use \(u_\alpha = \{0,0, v|0,0, u\}\) to represent 1st nearest neighbor (n.n.) mean-field bonds. \(u_\beta = \{0,1, w|0,0, u\}\) is the representative of 2nd n.n. mean-field bonds. There are two kinds of symmetry-unrelated 3rd n.n. mean-field bonds, represented by \(u_\gamma = \{1,0, u|0,0, u\}\) and \(\tilde{u}_\gamma = \{1,-1, u|0,0, u\}\). The symmetry conditions for these mean-field bonds are summarized in [B13]-[B16]. Besides, the on-site chemical potential terms \(\Lambda(i)\) (which guarantee the physical constraint [3] on the mean-field level) also satisfy symmetry conditions [B12]. We can show that \(\Lambda(x,y,s) = \Lambda_s\) for these 20 \(Z_2\) SL states. The symmetry-allowed mean-field amplitudes/bonds are also summarized in TABLE II.

From TABLE II we can see there are 6 states, i.e. \#7-\#12 that don’t allow nonzero 1st n.n. mean-field amplitudes due to symmetry. Moreover, they cannot realize \(Z_2\) SLs with up to 3rd n.n. mean-field amplitudes. Therefore they are unlikely to be the HKLM ground state. Ruling out these six \(Z_2\) SLs, we can see the other 14 \(Z_2\) SL states fall into 4 classes. To be specific, they are continuously connected to different parent \(U(1)\) gapless SL states on kagome lattice. These parent \(U(1)\) SL states in general have the following mean-field ansatz

\[
H_{U(1)SL} = \chi_1 \sum_{\langle ij \rangle} \nu_{ij} (f_{i\alpha}^\dagger f_{j\alpha} + h.c.)
\]

where \(\nu_{ij} = \pm 1\) characterizes the sign structure of hopping terms with \(\chi_1 \in \mathbb{R}\). Different parent \(U(1)\) SL
states are featured by the flux of $f$-spinon hopping phases around basic plaquette: honeycombs and triangles on the kagome lattice.

The simplest example is the so-called uniform RVB state with $\psi_j = +1$ for all 1st n.n. mean-field bonds. The hopping phase around any plaquette is $1 = \exp[i0]$, and the corresponding flux is $[0, 0]$ for [triangle,honeycomb] motifs. The 4 possible $Z_2$ spin liquids in the neighborhood of uniform RVB states (i.e. $U(1)$ SL-[0, 0] state in Ref. [2] are classified in Appendix D. They are #1, #5, #15, #13 in TABLE[1] and TABLE[11]. We label them as $Z_2[0, 0]A$, $Z_2[0, 0]B$, $Z_2[0, 0]C$ and $Z_2[0, 0]D$ states. They all have gapped spectra of spinons.

The ansatz of two other parent $U(1)$ SLs are shown in FIG. 4. They both have $\pi$-flux piercing through a triangle basic plaquette. Following the above notations of hopping phase in [triangle,honeycomb] motifs, with either $\pi$-flux or 0-flux through the honeycomb plaquette, they are called $U(1)$ SL-[\pi, \pi] state and $U(1)$ SL-[\pi, 0] state. There are three $Z_2$ SLs in the neighborhood of both $U(1)$ SL states, i.e. #3, #17, #19 around $U(1)$ SL-[\pi, \pi] state and #4, #18, #20 around $U(1)$ SL-[\pi, 0] state. All these six $Z_2$ SLs have gapless spinon spectra, inherited from the two parent gapless $U(1)$ SLs. To be precise, the spinon band structure of these six $Z_2$ SL states are featured by a doubly-degenerate flat band and a Dirac cone at Brillouin-zone center. This is in contrast to the numerically observed gap in two-spinon spectrum, thus we can also rule out these 6 $Z_2$ SLs for the HKLM ground state.

Another $U(1)$ SL state is the so called $U(1)$-Dirac SL or $U(1)$ SL-[0, $\pi$] state. Its mean-field ansatz is shown by the 1st n.n. bonds in FIG. 4(b). Clearly $\pi$-flux pierces through certain triangle plaquette with no flux through the honeycomb plaquette. According to variational Monte Carlo studies [32], this $(1)$-Dirac SL have substantially lower energy compared to many other competing phases, including the uniform RVB state. Therefore we shall focus on those $Z_2$ SLs in the neighborhood of the $U(1)$-Dirac SL in our search of the HKLM ground state. We need to mention that although unlikely, the four $Z_2$ SLs in the neighborhood of uniform RVB state, or $U(1)$ SL-[0, 0] state are potentially possible to be the HKLM ground state.

In a previous study using PSG in Schwinger-boson representation [2], it was shown that there are 8 different Schwinger-boson mean-field ansatz of $Z_2$ SLs on the kagome lattice which preserve all lattice symmetry. However, these 8 $Z_2$ SLs may or may not preserve time-reversal symmetry. One can show that requiring all lattice symmetry and time-reversal symmetry, there are 16 different Schwinger-boson $Z_2$ SLs on the kagome lattice. The relation between the 20 $Z_2$ SLs in Schwinger-
fermion representation (see TABLE I) and the 16 $Z_2$ SLs in Schwinger-boson representation are not clear. To clarify the relation between SL states in these two different representations, one can compare the neighboring (ordered) phases of the SLs, e.g. by computing the vison quantum numbers$^{33}$ of SL states.

B. $Z_2[0, \pi]\beta$ state as a promising candidate for the HKLM ground state

How to find those $Z_2$ SLs in the neighborhood of (or continuously connected to) the $U(1)$-Dirac SL? Naively, we expect the mean-field ansatz of these $Z_2$ SLs can be obtained from that of $U(1)$-Dirac SL by adding an infinitesimal perturbation. To be specific, we require an infinitesimal spinon pairing term on top of the $U(1)$-Dirac SL mean-field ansatz, or to break the IGG from $U(1)$ to $Z_2$ through Higgs mechanism. Mathematically, we need to find those $Z_2$ SL states whose PSG is a subgroup of the $U(1)$-Dirac SL’s PSG$^{33}$. Such $Z_2$ SL states are defined to be in the neighborhood of $U(1)$-Dirac SL. Similar criterion applies to the neighboring $Z_2$ SL states of any parent $U(1)$ or $SU(2)$ SL state.

We find out all four $Z_2$ SLs in the neighborhood of $U(1)$-Dirac SLs in Appendix C. They are states #6, #2, #14, #16 in TABLE I labeled as $Z_2[0, \pi]\alpha$, $Z_2[0, \pi]\beta$, $Z_2[0, \pi]\gamma$ and $Z_2[0, \pi]\delta$ states respectively. Since the effective theory of a $U(1)$-Dirac SL is an 8-component Dirac fermion coupled with dynamical $U(1)$ gauge fields,$^{33}$ we can find out all symmetry-allowed mass terms that can open up a gap in the Dirac-like spinon spectrum. Following detailed calculations in Appendix C, we can see that among the four $Z_2$ SLs around the $U(1)$-Dirac SL, only one state, i.e. $Z_2[0, \pi]\beta$ (state #2 in TABLE I) can generate a mass gap in the spinon spectrum. In other 3 states the Dirac cone in spinon spectrum is protected by symmetry. The mean-field ansatz of $Z_2[0, \pi]\beta$ SL state up to 2nd n.n. is shown in FIG. (1b):

\[ H_{MF} = \sum_i (\lambda_3 \sum_{\alpha} f_{i\alpha}^\dagger f_{i\alpha} + \lambda_1 f_{i\uparrow}^\dagger f_{i\downarrow}^\dagger + h.c.) + \chi \sum_{<ij>} \nu_{ij} (f_{i\alpha}^\dagger f_{j\alpha} + h.c.) + \Delta_2 \sum_{\alpha\beta} \epsilon^{\alpha\beta} f_{i\alpha}^\dagger f_{i\beta}^\dagger + h.c.) \]

where $\epsilon^{\alpha\beta}$ is the completely anti-symmetric tensor. We only list up to 2nd n.n. mean-field amplitudes because as shown in TABLE I (see also Appendix C), this $Z_2[0, \pi]\beta$ state only needs 2nd n.n. pairing terms to realize a $Z_2$ SL. We can always choose a proper gauge so that mean-field parameters $\chi_{1,2}$ and $\Delta_2$ are all real. The sign structure of $\nu_{ij}$ are $\pm 1$ are shown in FIG. (1b), with red denoting $\nu_{ij} = +1$ and other colors representing $\nu_{ij} = -1$. As discussed in Appendix C the 2nd n.n. singlet-pairing term $\Delta_2 \neq 0$ not only break the $U(1)$ gauge symmetry down to $Z_2$, but also opens up a mass gap in the spinon spectrum. The on-site chemical potential $\lambda_{1,3}$ are self-consistently determined by the following constraint:

\[ \sum_i \langle f_{i\uparrow}^\dagger f_{i\uparrow}^\dagger \rangle = \sum_i \langle f_{i\uparrow} f_{i\downarrow} \rangle = 0, \]
\[ \sum_i \langle \sum_{\alpha=\uparrow,\downarrow} f_{i\alpha}^\dagger f_{i\alpha} - 1 \rangle = 0. \]  

For further n.n. mean-field ansatz see discussions in Appendix C.

IV. CONCLUSION

To summarize, motivated by the strong evidence of a $Z_2$ SL as the HKLM ground state in recent DMRG study, we classify all possible $Z_2$ SL states in Schwinger-fermion mean-field approach using PSG. We found 20 different Schwinger-fermion mean-field states of $Z_2$ SLs on kagome lattice, among which 6 states are unlikely due to vanishing 1st n.n. mean-field amplitude. In other 14 $Z_2$ SLs only 5 possess a gapped spinon spectrum, which is observed in the DMRG results. These five symmetric $Z_2$ SL states are all in the neighborhood of certain parent $U(1)$ gapless SLs. To be precise, four are in the neighborhood of gapless uniform RVB (or $U(1)$ SL-[0,0]) state, while the other one, i.e. $Z_2[0, \pi]\beta$ is in the neighborhood of gapless $U(1)$-Dirac SL (or $U(1)$ SL-[0,0,0]) state. Previous variational Monte Carlo study$^{34}$ showed that gapless $U(1)$-Dirac SL has a substantially lower energy in comparison to the uniform RVB state. This suggests $Z_2$ SLs in the neighborhood of $U(1)$-Dirac SL should have lower energy compared to those in the neighborhood of uniform RVB state. Therefore we propose this $Z_2[0, \pi]$$\beta$ state with mean-field ansatz (10) shown in FIG. (1b) as the HKLM ground state numerically detected in Ref. 1. Our work provides important insight for future numeric study, e.g. variational Monte Carlo study of Gutzwiller projected wavefunctions.

Acknowledgments

YML thank Prof. Ziqiang Wang for support under DOE Grant DE-FG02-99ER45747. YR is supported by the startup fund at Boston College. PAL acknowledges the support under NSF DMR-0804040.

Appendix A: Symmetry group of kagome lattice and algebra conditions for $Z_2$ spin liquids

As shown in FIG. (1a), we label the three lattice sites in each unit cell with sublattice index $\{s = u, v, w\}$. Choosing Bravais unit vector as $\vec{a}_1 = a\hat{x}$ and $\vec{a}_2 = \frac{a}{\sqrt{3}}(\hat{x} + \sqrt{3}\hat{y})$, the positions of the three atoms in a unit cell labeled by indices $i = (x, y, s)$ are

\[ \vec{r}(x, y, u) = (x + \frac{1}{2})\vec{a}_1 + (y + \frac{1}{2})\vec{a}_2, \]
\[ \vec{r}(x, y, v) = (x + \frac{1}{2})\vec{a}_1 + y\vec{a}_2, \]
\[ \vec{r}(x, y, w) = x\vec{a}_1 + (y + \frac{1}{2})\vec{a}_2. \]
The symmetry group of such a two-dimensional kagome lattice is generated by the following operations

\[ T_1 : (x, y, s) \rightarrow (x + 1, y, s); \]  
\[ T_2 : (x, y, s) \rightarrow (x, y + 1, s); \]  
\[ \sigma : (x, y, u) \rightarrow (y, x, u), \]  
\[ (x, y, v) \rightarrow (y, x, w), \]  
\[ (x, y, w) \rightarrow (y, x, v); \]  
\[ C_6 : (x, y, u) \rightarrow (-y - 1, x + y + 1, v), \]  
\[ (x, y, v) \rightarrow (-y, x + y, w), \]  
\[ (x, y, w) \rightarrow (-y - 1, x + y, u). \]

together with time reversal \( T \).

The symmetry group of a kagome lattice is defined by the following algebraic relations between its generators:

\[ T^2 = \sigma^2 = (C_6)^6 = e, \]  
\[ g^{-1}T^{-1}gT = e, \quad \forall g = T_{1,2}, \sigma, C_6, \]

\[ T_1^{-1}T_2T_1 = e, \]
\[ \sigma^{-1}T_2^{-1}\sigma T_2 = e, \]
\[ \sigma^{-1}C_6\sigma C_6 = e, \]
\[ C_6^{-1}T_2^{-1}C_6T_1 = e, \]
\[ C_6^{-1}T_2^{-1}T_1C_6T_2 = e, \]
\[ \sigma^{-1}C_6\sigma C_6 = e. \]

where \( e \) stands for the identity element in the symmetry group. Therefore the consistent conditions for a generic \( Z_2 \) PSGs on a kagome lattice is written as

\[ [G_T(i)]^2 = \eta_T \tau^0, \]  
\[ G_\sigma(i)G_\sigma(i) = \eta_\sigma \tau^0, \]  
\[ G_{T_1}(i)G_{T_1}(i)G_{T_2}(i)G_{T_2}(i)G_{T_2}(i) = \eta_{T_1} \tau^0, \]  
\[ G_{T_2}(i)G_{T_2}(i)G_{T_1}(i)G_{T_1}(i)G_{T_1}(i) = \eta_{T_2} \tau^0, \]  
\[ G_{C_6}(i)G_{C_6}(i)G_{C_6}(i)G_{C_6}(i)G_{C_6}(i) = \eta_{C_6} \tau^0, \]  
\[ G_{T_1}(i)G_{T_1}(i)G_{T_2}(i)G_{T_2}(i)G_{T_2}(i)G_{T_2}(i) = \eta_{T_1} \tau^0, \]  
\[ G_{C_6}(i)G_{C_6}(i)G_{C_6}(i)G_{C_6}(i)G_{C_6}(i)G_{C_6}(i) = \eta_{C_6} \tau^0, \]

where \( \eta \) are \( Z_2 \) integers characterizing different \( SL \)s: different (gauge inequivalent) choices of these \( Z_2 \) integers (different \( Z_2 \) PSGs) correspond to different \( Z_2 \) SLs. Notice that under a local gauge transformation \( W(i) \in SU(2) \) the PSG element \( G_\sigma(i) \) transforms as

\[ G_\sigma(i) \rightarrow W(i)G_\sigma(i)W^\dagger(U^{-1}(i)) \]  

Appendix B: Classification of all \( Z_2 \) spin liquids on kagome lattice

1. Classification of \( Z_2 \) algebraic PSGs on kagome lattice

In this section we classify all possible \( Z_2 \) spin liquids on a kagome lattice. Mathematically we need to find out all gauge-equivalent solutions of algebraic conditions \[ [A1] - [A15] \] for \( Z_2 \) PSGs.

First from condition \[ [A10] \] we can always choose a proper gauge so that

\[ G_{T_1}(x, y, s) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^{xy}, \]  
\[ G_{T_2}(x, y, s) = \tau^0. \]  

From \[ [A12] \] and \[ [A13] \] we can see \( G_\sigma(x, y, s) = \eta_{\sigma T_1}^y \eta_{\sigma T_2}^x \eta_{C_6}^{xy} \eta_\sigma(s) \). Condition \[ [A5] \] further determine \( \eta_{\sigma T_1} = \eta_{\sigma T_2} \) and therefore we have

\[ G_\sigma(x, y, s) = \eta_{\sigma T_1}^y \eta_{\sigma T_2}^x \eta_{C_6}^{xy} \eta_\sigma(s) \]  

where \( SU(2) \) matrices \( g_\sigma(s) \) satisfy

\[ g_\sigma(w)g_\sigma(v) = [\eta_\sigma(w)^{-1}]^2 = \eta_\sigma \tau^0 \]

Notice that we can always choose a proper global \( Z_2 \) gauge on \( G_{T_1}(x, y, s) \) (which doesn’t change the mean-field ansatz) so that \( \eta_{C_6}T_2 = 1 \) in \[ [A10] \]. From \[ [A11] \] and \[ [A16] \] it’s straightforward to show that

\[ G_{C_6}(x, y, u/v) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^y \eta_{C_6}^x \eta_{C_6}^w \]  
\[ G_{C_6}(x, y, u) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^y \eta_{C_6}^x \eta_{C_6}^w \]  
\[ G_{C_6}(x, y, w) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^y \eta_{C_6}^x \eta_{C_6}^w \]  
\[ G_{C_6}(x, y, v) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^y \eta_{C_6}^x \eta_{C_6}^w \]

where \( SU(2) \) matrices \( g_{C_6}(s) \) satisfy

\[ [g_{C_6}(u)g_{C_6}(v)g_{C_6}(u)]^2 = \eta_{2\eta_{C_6} \tau^0}, \]

\[ [\eta_{C_6}(s)]^2 = \eta_{\sigma C_6} \tau^0. \]

Now through a gauge transformation \( W(x, y, s) = \eta_{T_1}^y \eta_{T_2}^x \) we can fix \( \eta_{T_1,2} = 1 \) and the PSG elements become

\[ G_{C_6}(x, y, s) = \eta_{T_1}^y \eta_{T_2}^x \eta_\sigma(s), \]

\[ G_{C_6}(x, y, u/v) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^y \eta_{C_6}^x \eta_{C_6}^w \]

\[ G_{C_6}(x, y, w) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^y \eta_{C_6}^x \eta_{C_6}^w \]

\[ G_{C_6}(x, y, v) = \eta_{T_1}^y \eta_{T_2}^x \eta_{C_6}^y \eta_{C_6}^x \eta_{C_6}^w \].
According to (A4), (A6) and (A7) we can see that

\[ G_T(x, y, s) = \eta_T^T \eta_T^T g_T(s), \]

further determines \( \eta_T^T g_T = 1 \) and by choosing a proper gauge we have

\[ G_T(x, y, s) = g_T(s) \equiv \begin{cases} \tau^0, & \eta_T = 1, \\ i\tau^1, & \eta_T = -1. \end{cases} \]  

(B7)

which satisfy

\[ g_T(u)g_T(u) = \eta_T^T g_T(u)g_T(u), \]

\[ g_T(v)g_T(w) = \eta_T^T g_T(v)g_T(w), \]

\[ g_T(w)g_T(v) = \eta_T^T g_T(w)g_T(v), \]

\[ g_C(u)g_T(u) = g_C(u)g_T(u), \]

\[ g_C(u)g_T(v) = g_C(u)g_T(v), \]

\[ g_C(u)g_T(w) = g_C(u)g_T(w), \]

\[ g_C(w)g_T(v) = g_C(w)g_T(v), \]

\[ g_C(w)g_T(w) = g_C(w)g_T(w). \]

(B8)

according to (A9) and (A8).

In the following we find out all the gauge-inequivalent solutions of \( SU(2) \) matrices \( T, \sigma, C \) satisfying the above conditions. They are summarized in TABLE .

(I) \( g_T(s) = \tau^0 \) and therefore \( \eta_T = \eta_T^T = \eta_C = 1 \):

Conditions (B3) and (B4) are automatically satisfied.

(i) \( \eta_T = 1 \):

Notice that under a global gauge transformation \( W(x, y, s) = W_s \in SU(2) \) the PSG elements transform as

\[ g_T(u) \rightarrow W_u g_T(u) \]

\[ g_T(v) \rightarrow W_v g_T(v) \]

\[ g_T(w) \rightarrow W_w g_T(w) \]

\[ g_C(u) \rightarrow W_u g_C(u) \]

\[ g_C(v) \rightarrow W_v g_C(v) \]

\[ g_C(w) \rightarrow W_w g_C(w). \]

Thus from (B2) and (B4) we can always have \( g_T(s) = \tau^0 \) and \( g_C(u) = \tau^0, g_C(v) = \eta_C \tau^0 \) by choosing a proper gauge.

(A) \( \eta_C \tau^0 = \eta_C \tau^0 = 1 \):

from (B3) we have \( g_C(w) = \tau^0 \).

(B) \( \eta_C \tau^0 = \eta_C \tau^0 = -1 \):

from (B3) we have \( g_C(w) = i\tau^3 \) by gauge fixing.

(ii) \( \eta_T = -1 \):

from (B2) we have \( g_T(s) = -g_T(s) = \tau^0 \) and \( g_T(u) = i\tau^1 \) by gauge fixing. Also from (B4) we can choose a gauge so that \( g_C(u) = \tau^0 \) and \( g_C(v) = -i\eta_C \tau^3 \).

(A) \( \eta_C \tau^0 = -1 \):

In this case (B4) requires \( g_C(w) = \tau^0 \) and thus

\[ \eta_C \tau^0 = -1 \]  

(B) \( \eta_C \tau^0 = 1 \):

Now from (B4) and (B3) we have \( g_C(w) = i\tau^1 \) by gauge fixing.

(b) \( \eta_C \tau^0 = 1 \):

by (B4) and (B3) we must have \( g_C(w) = i\tau^3 \).

To summarize there are \( 2 \times (2 + 3) = 10 \) different algebraic PSGs with \( \eta_T = 1 \) and \( g_T(s) = \tau^0 \).

(II) \( g_T(s) = i\tau^1 \) and \( \eta_T = -1 \):

(i) \( \eta_T = 1 \):

According to (B2) and (B3), by choosing a proper gauge we can have \( g_T(s) = \tau^0 \) and \( \eta_T = 1 \). From (B3) and (B4) we also have \( [g_C(w)]^2 = g_C(w)g_C(w) = \eta_C \tau^0 \).

(A) \( \eta_C \tau^0 = \eta_C \tau^0 = 1 \):

from (B3) and (B4), by choosing gauge we have

\[ g_C(s) = \tau^0 \]  

and \( \eta_C \tau^0 = 1 \).

(B) \( \eta_C \tau^0 = \eta_C \tau^0 = -1 \):

(a) \( \eta_C \tau^0 = 1 \):

In this case we have \( g_C(u) = -g_C(v) = \tau^0 \) and \( g_C(w) = i\tau^1 \) by choosing a proper gauge.

(b) \( \eta_C \tau^0 = -1 \):

In this case we can have \( g_C(s) = i\tau^3 \) by choosing a proper gauge.

(ii) \( \eta_T = -1 \):

(A) \( \eta_T = 1 \):

From (B3) and (B2) we have \( g_T(s) = i\tau^1 \) and \( g_T(u) = -g_T(u) = \tau^0 \) by proper gauge fixing. Also from (B4) we know \( [g_C(w)]^2 = \eta_C \tau^0 \) and \( g_C(u)g_C(v) = -i\eta_C \tau^1 \).

(a) \( \eta_C \tau^0 = -1 \):
from (9) and (11), it’s clear that \( \eta_{n,T} = 1 \), \( g_{c_0}(u) = g_{c_0}(v) = \tau^0 \) and \( g_{c_0}(v) = i\tau^1 \) through gauge fixing. Also we have \( \eta_{12} \eta_{c_0} = -1 \).

(b) \( \eta_{c_0} = 1 \):

(b1) \( \eta_{c_0} T = 1 \):

In this case \( \eta_{12} \eta_{c_0} = 1 \), and we can always choose a proper gauge so that \( g_{c_0}(u) = \tau^0 \), \( g_{c_0}(w) = -g_{c_0}(v) = i\tau^1 \).

(b2) \( \eta_{c_0} T = -1 \):

In this case \( \eta_{12} \eta_{c_0} = -1 \), and we can always choose a proper gauge so that \( g_{c_0}(v) = -i\tau^2 \), \( g_{c_0}(v) = g_{c_0}(w) = i\tau^3 \).

(B) \( \eta_{c_0} T = -1 \):

Conditions (8) and (9) assert that \( g_\sigma(s) = i\tau^3 \) by proper gauge choosing.

(a) \( \eta_{c_0} = 1 \):

In this case from (11) we know \( g_{c_0}(w) = i\tau^3 \), hence \( \eta_{c_0} T = -1 \). Then we can always choose a gauge so that \( g_{c_0}(u) = g_{c_0}(v) = i\tau^3 \) and so \( \eta_{12} \eta_{c_0} = -1 \) from (9).

(b) \( \eta_{c_0} = 1 \):

(b1) \( \eta_{c_0} T = 1 \):

In this case from (8), (10) we have \( g_{c_0}(u) = g_{c_0}(v) = \tau^0 \) by a proper gauge choice. Meanwhile conditions (9) and (10) become \( [g_{c_0}(w)]^2 = \eta_{12} \eta_{c_0} \tau^0 \) and \( [i\tau^3 g_{c_0}(w)]^2 = -\tau^0 \).

(b1.1) \( \eta_{12} \eta_{c_0} = 1 \):

Here we have \( g_{c_0}(w) = \tau^0 \).

(b1.2) \( \eta_{12} \eta_{c_0} = -1 \):

Here we have \( g_{c_0}(w) = i\tau^1 \).

(b2) \( \eta_{c_0} T = -1 \):

In this case from (8) and (10) we can always choose a proper gauge so that \( g_{c_0}(u) = -g_{c_0}(v) = i\tau^3 \). We also have \( g_{c_0}(w) = i\tau^2 \) and \( \eta_{12} \eta_{c_0} = -1 \) from (10).

To summarize there are \( 2 \times (3 + 7) = 20 \) different algebraic PSGs with \( \eta_T = -1 \) and \( g_T(s) = i\tau^1 \).

So in summary we have \( 10 + 20 = 30 \) different \( Z_2 \) algebraic PSGs satisfying conditions (4)-(10). Among them there are at most 20 solutions that can be realized by a mean-field ansatz, since those PSGs with \( g_T(s) = \tau^0 \) would require all mean-field bonds to vanish due to (11). As a result there are \( 20 \) different \( Z_2 \) spin liquids on a kagome lattice.

2. Symmetry conditions on mean-field ansatz

Let’s denote the mean-field bonds connecting sites \((0,0,u)\) and \((x,y,s)\) as \([x,y,s] = (x,y,s) = (0,0,u)\). Using (8) we can generate any other mean-field bonds through symmetry operations (such as translation \( G_{T_1}, G_{T_2} \) and mirror reflection \( G_{\sigma} \)) from \([x,y,s]\). However these mean-field bonds cannot be chosen arbitrarily since they possess symmetry relation (8):

\[
\langle ij \rangle = G_U(i) \langle U^{-1}(i) U^{-1}(j) \rangle G_U^\dagger(j)
\]

(B10)

where \( U \) is any element in the symmetry group. Notice that for time reversal \( T \) we have

\[
G_T(i) \langle ij \rangle G_T^\dagger(j) = -\langle ij \rangle
\]

(B11)

We summarize these symmetry conditions on the mean-field bonds here:

(i) For \( s = u \)

\[
T : \ g_T[x,y,u]g_T^\dagger = -[x,y,u],
\]

\[
T_1^x T_2^y \sigma : \ [x,-x,u] \rightarrow [x,-x,u]^\dagger,
\]

\[
T_1^{x+1} T_2^{y+1} C_6^\dagger : \ [x,y,u] \rightarrow [x,y,u]^\dagger,
\]

\[
\sigma : \ [x,y,u] \rightarrow [x,x,u].
\]

(ii) For \( s = v \)

\[
T : \ g_T[x,y,v]g_T^\dagger = [x,y,v],
\]

\[
T_2^{y+1} \sigma C_6^\dagger \ : \ [0,y,v] \rightarrow [0,y,v]^\dagger,
\]

\[
T_1^{x+1} T_2^{2y-1} C_6^{-1} \ : \ [1-2y,y,v] \rightarrow [1-2y,y,v]^\dagger.
\]

(iii) For \( s = w \)

\[
T : \ g_T[x,y,w]g_T^\dagger = [x,y,w],
\]

\[
T_1^{x-1} T_2^{2-x} \sigma C_6^{-1} \ : \ [x,1-2x,w] \rightarrow [x,1-2x,w]^\dagger,
\]

\[
T_1^{x+1} \sigma C_6^{-2} \ : \ [x,0,w] \rightarrow [x,0,w]^\dagger.
\]

Now let’s consider several simplest examples. At first, on-site chemical potential terms \( 0(x,y,s) = \Lambda_s \) satisfy the following consistent conditions:

\[
\tau^1 \Lambda_s \tau^1 = -\Lambda_s; \quad \sigma \Lambda_s \sigma = \Lambda_s; \quad \Lambda_s = \Lambda_s
\]

(B12)

\[
g_{\sigma}(u)\Lambda_u g_{\sigma}(v) = \Lambda_u,
\]

\[
g_{\sigma}(u)\Lambda_u g_{\sigma}(v) = \Lambda_v,
\]

\[
g_{\sigma}(w)\Lambda_v g_{\sigma}(v) = \Lambda_u;
\]

\[
g_{c_0}(u)\Lambda_u g_{c_0}(v) = \Lambda_u,
\]

\[
g_{c_0}(v)\Lambda_v g_{c_0}(v) = \Lambda_w,
\]

\[
g_{c_0}(w)\Lambda_v g_{c_0}(w) = \Lambda_w.
\]

In fact in all 20 \( Z_2 \) spin on a kagome lattice we all have \( \Lambda_u = \Lambda_v = \Lambda_w = \Lambda_s \) with a proper gauge choice.

All the 1st n.n. mean-field bonds can be generated from \( u_\alpha = [0,0,v] \). For a generic \( Z_2 \) spin liquid with PSG elements \( G_T(x,y,s) = i\tau^1 \) and (11)-(13), the bond \( u_\alpha = [0,0,v] \) satisfies the following consistent conditions:

\[
\tau^1 u_{\alpha} \tau^1 = -u_{\alpha},
\]

(B13)

\[
g_{\sigma}(u)g_{c_0}(u)g_{c_0}(w)u_{\alpha} g_{c_0}(v)g_{c_0}(w)g_{\sigma}(v) = u_{\alpha}.
\]

It follows immediately that for six \( Z_2 \) spin liquids, i.e. \#7 - 12 in TABLE (11) all n.n. mean-field bonds must vanish since \( u_\alpha = 0 \) as required by (11). Therefore it’s unlikely that the \( Z_2 \) spin liquid realized in kagome Hubbard model would be one of these 6 states. In the following we study the rest 14 \( Z_2 \) spin liquids on the kagome lattice.
All 2nd n.n. mean-field bonds can be generated from \( u_{\beta} \equiv [0, 1, w] \) which satisfies the following symmetry conditions

\[
\tau_1 = -u_{\beta}, \quad g_\sigma(u)g_{\sigma}(w)u_{\beta} = u_{\beta}.
\]

(B14)

There are two kinds of 3rd n.n. mean-field bonds: the first kind can all be generated by \( u_{\gamma} \equiv [1, 0, u] \) which satisfies

\[
\tau_1 = -u_{\gamma}, \quad g_\sigma(u)g_{\sigma}(w)u_{\gamma}g_{\sigma}(v)g_{\sigma}(w) = u_{\gamma}.
\]

(B15)

the second kind can all be generated by \( \tilde{u}_{\gamma} \equiv [1, -1, u] \) which satisfies

\[
\tau_1 = -\tilde{u}_{\gamma}, \quad g_\sigma(u)\tilde{g}_\sigma(u)\tilde{u}_{\gamma}g_{\sigma}(v)g_{\sigma}(w) = \eta_{\tilde{u}}\tilde{u}_{\gamma}.
\]

(B16)

\[ g_{\sigma}(u)g_{\sigma}(w)g_{\sigma}(v)g_{\sigma}(w)u_{\gamma}g_{\sigma}(v)g_{\sigma}(w) = \eta_{\tilde{u}}\tilde{u}_{\gamma}. \]

App. C: \( Z_2 \) spin liquids in the neighborhood of \( U(1) \) SL-0, \( \pi \) state

1. Mean-field ansatz of \( U(1) \) SL-0, \( \pi \) state

Following \( SU(2) \) Schwinger fermion formulation with \( \psi_i \equiv (f_{1i}, f_{2i})^T \), we focus on those \( Z_2 \) spin liquids (SLs) in the neighborhood of \( U(1) \) SL-0, \( \pi \) state with the following mean-field ansatz:

\[
(x, y, u|x, y, v) = -(x, y, u|x, y, w) = (-1)^x \chi \tau^3, \quad (1)
\]

\[
(x + 1, y, w|x, y, u) = (x + 1, y + 1, v|x, y, u) = -(x, y, v|x, y, w), \quad (2)
\]

\[
= -(x + 1, y - 1, w|x, y, v) = \chi \tau^3. \quad (3)
\]

where \( \chi \) is a real hopping parameter. We define mean-field bonds \( (x, y, s|x', y', s') \) in the following way

\[
H_{MF} = \sum_{i,j} \psi_i^\dagger(i|j)\psi_j + \text{h.c.} \quad (C2)
\]

For convenience of later calculation we implement the following gauge transformation

\[
\psi_{x,y,u} \rightarrow \tau^3_1 \psi_{x,y,u} \quad (C3)
\]

and the original mean-field ansatz transforms to be

\[
\langle x, y, u|x, y, v \rangle = -(x, y, u|x, y, w) = \chi \tau^0, \quad (C4)
\]

\[
\langle x + 1, y, w|x, y, u \rangle = (x + 1, y + 1, v|x, y, u) = \chi \tau^0,
\]

\[
= -(x, y, v|x, y, w) = \chi \tau^3. \quad (C5)
\]

so that the mean-field ansatz satisfy (3).

2. Classification of \( Z_2 \) spin liquids around \( U(1) \) SL-0, \( \pi \) state

Plugging (C5) into algebraic consistent conditions \( (A4), (A15) \) yields four algebraic solutions of \( Z_2 \) PSGs around the \( U(1) \) SL-0, \( \pi \) state. Choosing a proper gauge they all satisfy

\[
g_T = i\tau_1, \quad g_T = g_T = \eta_T = \eta_{\tau} = \eta_{\tau} = 1, \quad (C6)
\]

The four \( Z_2 \) PSGs near the \( U(1) \) SL-0, \( \pi \) state are featured by

\[
\langle 0, \pi\rangle \alpha: \quad \gamma = g_{C_0} = \eta_{C_0} = 1, \quad (C7)
\]

\[
\langle 0, \pi\rangle \beta: \quad \gamma = \eta_{C_0} = -1; \quad (C8)
\]

\[
\langle 0, \pi\rangle \gamma: \quad \gamma = i\tau^3, \quad \eta_{\tau} = \eta_{\tau} = -1; \quad (C9)
\]

\[
\langle 0, \pi\rangle \delta: \quad \gamma = g_{C_0} = i\tau^3 \quad (C10)
\]

Of course they belong to the 20 \( Z_2 \) spin liquids summarized in Table [1].
3. Four possible $Z_2$ spin liquids around $U(1)$ SL-$[0, \pi]$ state: mean-field ansatz

a. Consistent conditions on mean-field bonds

Implementing the generic conditions mentioned earlier on several near neighbor mean-field bonds with PSG \([C9]-[C10]\), we obtain the following consistent conditions:

(0) For on-site chemical potential terms $\Lambda_s(x,y,s) = \hat{\lambda} (x,y,s) \cdot \vec{r}$, translations operations $G_{T_x} T_{1,2}$ in PSG guarantee that $\Lambda_s(x,y,s) = \Lambda_s(0,0,s) \equiv \Lambda_s$, $s = u, v, w$. They satisfy

$$g_T \Lambda_s g_T^\dagger = -\Lambda_s; \quad (C11)$$

$$g_s \Lambda_u g_s^\dagger = \Lambda_u, \quad g_s \Lambda_v g_s^\dagger = \Lambda_v; \quad (C12)$$

$$g_{C_u} \Lambda_u g_{C_u}^\dagger = \Lambda_v, \quad (g_{C_v} \tau^3) \Lambda_u (g_{C_v} \tau^3)^\dagger = \Lambda_u; \quad (C13)$$

$$g_{C_v} \Lambda_w g_{C_v}^\dagger = \Lambda_u. \quad (C14)$$

(1) For 1st neighbor mean-field bond $u_a \equiv [0,0,v]^\dagger$ (there is only one independent mean-field bond, meaning all other 1st neighbor bonds can be generated from $[0,0,v]$ through symmetry operations)

$$u_a = i a_0 \tau^0 + a_1 \tau^1, \quad w_b = i b_0 \tau^0 + b_1 \tau^1, \quad (C15)$$

$$u_{c1} = c_2 \tau^2 + c_3 \tau^3, \quad u_{c2} = 0, \quad (C16)$$

$$\Lambda_u = \lambda_2 \tau^2 + \lambda_3 \tau^3, \quad \Lambda_{s,w} = -\lambda_2 \tau^2 + \lambda_3 \tau^3. \quad (C17)$$

b. Mean-field ansatz of the four $Z_2$ spin liquids near $U(1)$ SL-$[0, \pi]$ state

For $Z_2[0, \pi] \alpha$ state with $g_s = g_{C_u} = \tau^0$ the mean-field ansatz are (up to 3rd neighbor mean-field bonds)

$$u_a = i a_0 \tau^0 + a_1 \tau^1, \quad u_b = i b_0 \tau^0 + b_1 \tau^1, \quad (C18)$$

$$u_{c1} = c_2 \tau^2 + c_3 \tau^3, \quad u_{c2} = 0, \quad \Lambda_u = \lambda_3 \tau^3, \quad (C19)$$

$$s = u, v, w. \quad (C20)$$

Since we are considering a phase perturbed from the $U(1)$ SL-$[0, \pi]$ state, we shall always assume $a_0 \neq 0$ (1st neighbor hopping terms) in the following discussion. A $Z_2[0, \pi] \alpha$ spin liquid can be realized by 1st neighbor mean-field singlet pairing terms with $a_0 \neq 0$.

For $Z_2[0, \pi] \beta$ state with $g_s = \tau^0$, $g_{C_u} = i \tau^3$ the mean-field ansatz are (up to 3rd neighbor mean-field bonds)

$$u_a = i a_0 \tau^0 + a_1 \tau^1, \quad u_b = i b_0 \tau^0 + b_1 \tau^1, \quad (C21)$$

$$u_{c1} = c_2 \tau^2 + c_3 \tau^3, \quad u_{c2} = 0, \quad \Lambda_u = \lambda_3 \tau^3, \quad s = u, v, w. \quad (C22)$$

A $Z_2[0, \pi] \gamma$ spin liquid can be realized by 2nd neighbor pairing terms with $a_0 \neq 0$.

For $Z_2[0, \pi] \delta$ state with $g_s = \tau^0$, $g_{C_u} = \tau^3$ the mean-field ansatz are (up to 3rd neighbor mean-field bonds)

$$u_a = i a_0 \tau^0, \quad u_b = i b_0 \tau^0 + b_1 \tau^1, \quad (C23)$$

$$u_{c1} = c_2 \tau^2 + c_3 \tau^3, \quad u_{c2} = 0, \quad \Lambda_u = \lambda_3 \tau^3, \quad s = u, v, w. \quad (C24)$$

4. Low-energy effective theory

The reciprocal unit vectors (corresponding to unit vectors $\mathbf{a}_1, \mathbf{a}_2$) on a kagome lattice are $\mathbf{b}_1 = \frac{1}{a}(\hat{x} - \frac{\sqrt{3}}{a}\hat{y})$ and $\mathbf{b}_2 = \frac{1}{a}(\frac{2}{\sqrt{3}}\hat{y})$, satisfying $\mathbf{a}_1 \cdot \mathbf{b}_j = \delta_{i,j}$. In the mean-field ansatz \([C1]\) of $U(1)$ SL-$[0, \pi]$ the unit cell is doubled whose translation unit vectors are $\mathbf{A}_1 = 2\mathbf{a}_1$ and $\mathbf{A}_2 = \mathbf{A}_2$. Accordingly the 1st BZ for such a mean-field ansatz is only half of the original 1st BZ with new reciprocal unit vectors being $\mathbf{A}_1 = \mathbf{b}_1/2$ and $\mathbf{A}_2 = \mathbf{b}_2$. Denoting the momentum as $k \equiv (k_x, k_y)/a = k_1 \mathbf{B}_1 + k_2 \mathbf{B}_2$ with $|k_{1,2}| \leq \pi$, we have

$$k_1 = 2k_x, \quad k_2 = (k_x + \sqrt{3}k_y)/2. \quad (C25)$$

The two Dirac cones in the spectra of $U(1)$ SL-$[0, \pi]$ state \([C4]\) are located at $\pm \mathbf{Q}$ with

$$\mathbf{Q} = (0, \frac{\pi}{\sqrt{3}}) = \frac{\pi}{2} \mathbf{b}_2 \quad (C26)$$

with the proper chemical potential $\Lambda(i) = |i| = \chi(\sqrt{3} - 1) \tau^3$ added to mean-field ansatz \([C4]\).
For convenience we choose the following basis for Dirac-like Hamiltonian obtained from expansion around \( \pm \mathbf{Q} \):

\[
\phi_{\pm,\uparrow,A} = \frac{1}{\sqrt{6}} e^{-i \frac{\pi}{3}} e^{-i \frac{\pi}{3}},
\]

\[
(e^{-i \frac{\pi}{3}}, 0, e^{i \frac{\pi}{3}}, 0, 0, 0, e^{-i \frac{\pi}{3}}, 0, e^{i \frac{\pi}{3}}, 0, e^{-i \frac{\pi}{3}}, 0, \sqrt{2}, 0)^T,
\]

\[
\phi_{\pm,\uparrow,B} = \frac{1}{\sqrt{6}} e^{-i \frac{\pi}{3}},
\]

\[
(1, 0, e^{-i \frac{\pi}{3}}, 0, \sqrt{2} e^{-i \frac{\pi}{3}}, 0, -1, 0, e^{-i \frac{\pi}{3}}, 0, 0, 0)^T,
\]

\[
\phi_{-\pm,\uparrow,b} = R_{T_1}(k_1 = 0, k_2 = -\frac{\pi}{6}) \phi_{\pm,\uparrow,b},
\]

\[
\phi_{\pm,\uparrow,b} = R_{T_2} \phi_{\pm,\uparrow,b}.
\]

(C22)

where \( \pm \) are valley index for two Dirac cones at \( \pm \mathbf{Q} \) with Pauli matrices \( \mu \) and \( b = A, B \) are band indices (for the two bands forming the Dirac cone) with Pauli matrices \( \nu \). Pseudospin indices \( \Sigma = \uparrow, \downarrow \) are assigned to the two degenerate bands related by time reversal, with Pauli matrices \( \sigma \). The corresponding creation operators for these modes are \( \Psi_{\pm,\Sigma,b}^\dagger = \psi_{\pm,\Sigma}^\dagger \phi_{\pm,\Sigma,b} \) in the order of \((0, 0, u), (0, 0, v), (0, v), (0, u), (1, 0, u), (1, 0, v), (1, 0, w)\) for the six sites per doubled new unit cell. Notice that in terms of \( f \)-spinons we have \( \psi_{\pm}^\dagger = (f^\dagger, f_\uparrow) \).

Here \( R_T \equiv I_{2 \times 2} \otimes \begin{bmatrix} 0 & 0 & 0 \\ 0 & 0 & 1 \\ 1 & 0 & 0 \end{bmatrix} \otimes g_T \), \( R_{T_2}(k) = e^{-i k_2} I_{6 \times 6} \otimes g_{T_2} \) and \( R_{T_1}(k) = \begin{bmatrix} 0 & -e^{-i k_2} \\ 1 & 0 \end{bmatrix} \otimes g_{T_1} \) are transformation matrices on 12-component eigenvectors for time reversal \( T \) and translation \( T_{1,2} \) operations. By definition of PSG the eigenvectors \( \phi_k \) with momentum \( k = k_1 \mathbf{B}_1 + k_2 \mathbf{B}_2 \equiv (k_1, k_2) \) and energy \( E \) have the following symmetric properties:

\[
T : \quad \hat{\phi}_{(k_1, k_2)}(E) = R_T \hat{\phi}_{(k_1, k_2)}(-E),
\]

\[
T_1 : \quad \hat{\phi}_{(k_1, k_2)}(E) = R_{T_1}(k_1, k_2) \hat{\phi}_{(k_1, k_2)}(E),
\]

\[
T_2 : \quad \hat{\phi}_{(k_1, k_2)}(E) = R_{T_2}(k_1, k_2) \hat{\phi}_{(k_1, k_2)}(E).
\]

\( \hat{\phi} \) and \( \phi \) are the basis after and before the symmetry operations.

In such a set of basis the Dirac Hamiltonian obtained by expanding the \( U(1) \) SL-\([0, \pi]\) mean-field ansatz \( \text{(C4)} \) around the two cones at \( \pm \mathbf{Q} \) is

\[
H_{\text{Dirac}} = \sum_k \frac{\sqrt{2}}{2} \Psi_{k} \mu \sigma \delta \left( -k_x \nu_1 + k_y \nu_2 \right) \Psi_{k} \quad \text{(C23)}
\]

\( \mathbf{k} \) should be understood as small momenta measured from \( \pm \mathbf{Q} \). Possible mass terms are \( \mu^{0,1,2,3} \sigma^{1,2,3} \nu^0 \) and \( \mu^{1,2,3} \sigma^{0,3} \nu^3 \). However not all of them are allowed by symmetry. Here we numerate all symmetry operations and associated operator transformations:

Spin rotation along \( \hat{z} \)-axis by angle \( \theta \):

\[
\Psi_{k} \rightarrow \Psi_{k} e^{i \frac{\theta}{2}}
\]

Spin rotation along \( \hat{y} \)-axis by angle \( \pi \):

\[
\Psi_{k} \rightarrow \Psi_{k}^T \mu^2 \sigma^2 \nu^2
\]

Time reversal \( T \):

\[
\Psi_{k} \rightarrow \Psi_{-k} \left( -i \sigma^2 \right)
\]

Translation \( T_1 \):

\[
\Psi_{k} \rightarrow \Psi_{k} \left( -i \sigma^3 \right)
\]

Translation \( T_2 \):

\[
\Psi_{k} \rightarrow \Psi_{k}^\dagger \left( -i \mu^3 \right)
\]

Considering the above conditions, the only symmetry-allowed mass terms are \( \sum_k \Psi_{k} \mu m_{1,2} \Psi_{k} \) with \( m_1 = \mu^0 \sigma^1 \nu^0 \) and \( m_2 = \mu^3 \sigma^3 \nu^3 \).

The transformation rules for mirror reflection \( \sigma \) and \( \pi/3 \) rotation \( C_6 \) depend on the choice of \( g_\sigma, g_\nu \) in the PSG. In general we have

\[
\sigma: \quad \Psi_{k} \rightarrow \Psi_{-k} M_{\sigma}(g_\sigma),
\]

\[
C_6: \quad \Psi_{k} \rightarrow \Psi_{C_6k} M_{C_6}(g_\nu).
\]

Using the basis \( \text{(C22)} \) the \( 8 \times 8 \) matrices \( M_{\sigma}, M_{C_6} \) can be expressed in terms of Pauli matrices \( \mu \otimes \sigma \otimes \nu \). For the \( Z_2 \) spin liquid we have

\[
M_{\sigma}(g_\sigma = \tau^0) = \mu^3 \otimes \sigma^0 \otimes \left( \begin{array}{cc} 0 & e^{-i \frac{\pi}{3}} \\ e^{-i \frac{\pi}{3}} & 0 \end{array} \right),
\]

\[
M_{\sigma}(g_\sigma = i \tau^3) = \mu^3 \otimes \sigma^3 \otimes \left( \begin{array}{cc} 0 & e^{i \frac{\pi}{3}} \\ e^{i \frac{\pi}{3}} & 0 \end{array} \right),
\]

\[
M_{C_6}(g_\nu = \tau^0) = \left( \begin{array}{cc} 1 & 0 \\ 0 & i \end{array} \right) \otimes \sigma^0 \otimes e^{i \frac{\pi}{3} \nu^0},
\]

\[
M_{C_6}(g_\nu = i \tau^3) = \left( \begin{array}{cc} 1 & 0 \\ 0 & -1 \end{array} \right) \otimes \sigma^0 \otimes e^{i \frac{\pi}{3} \nu^0}.
\]

It turns out in \( Z_2[0, \pi] \beta \) state only the 1st mass term \( m_1 = \mu^0 \sigma^1 \nu^0 \) is invariant under \( \sigma \) and \( C_6 \) operations. In other 3 states neither mass terms \( m_{1,2} \) are symmetry-allowed. As a result we only have one gapped \( Z_2 \) spin liquid, \( \text{i.e.} \ Z_2[0, \pi] \beta \) state in the neighborhood of \( U(1) \) Dirac SL-\([0, \pi]\) state.

Let’s consider mean-field bonds up to 2nd neighbor for ansatz \( Z_2[0, \pi] \beta \). Perturbations to the two Dirac cones of \( U(1) \) SL-\([0, \pi]\) with \( \lambda_3 = (\sqrt{3} - 1)a_0 \) in general has the following form

\[
\delta H_0 = [\lambda_3 - (\sqrt{3} - 1)a_0 - (\sqrt{3} + 1)b_0] \mu^0 \sigma^3 \nu^0 + [(\sqrt{3} + 1)b_1 - \lambda_2 - (\sqrt{3} - 1)a_1 \mu^0 \sigma^1 \nu^0 \quad \text{(C24)}
\]

This means we need either 1st neighbor \( (a_1) \) or 2nd neighbor \( (b_1) \) pairing term to open up a gap in the spectrum. Meanwhile these pairing terms break the original \( U(1) \) symmetry down to \( Z_2 \) symmetry.
Appendix D: $Z_2$ spin liquids in the neighborhood of uniform RVB state

The mean-field ansatz of the uniform RVB state is simple:

$$ H_{MF} = \chi \sum_{<ij>\sigma} f_i^\dagger f_j, \quad \sigma $$

(D1)

where $\chi$ is a real parameter and $<ij>$ represents sites $i,j$ being nearest neighbor (n.n.) of each other. It's straightforward to show the PSG elements of such a mean-field ansatz are

$$ G_U(x,y,s) = g_U, \quad U = T_{1,2}, \sigma, C_6. $$

(D2)

and $SU(2)$ matrices $g_U$ satisfy

$$ g_U T^{\sigma} g_U^\dagger = -T^{\sigma}, \quad g_U T^{\sigma} g_U^\dagger = T^{\sigma}, \quad U = T_{1,2},\sigma, C_6. $$

(D3)

It turns out these four $Z_2$ PSGs as solutions to (A4)-(A16) with the form (D2). In other words, there are only 4 different $Z_2$ in the neighborhood of a uniform RVB states. Choosing a proper gauge they all satisfy $g_T = i r_1$, $g_{T_{1,2}} = r_0$, and $\eta_\sigma T = \eta_{\sigma T} = \eta_\sigma C_{T_{1,2}} = \eta_\sigma T_{1,2} = 1$, $\eta_T = -1$. These four states are characterized by:

(#1) $Z_2[0,0]|A\rangle$: $g_\sigma = g_{C_6} = r_0$, $\eta_\sigma T = \eta_{C_6} = \eta_\sigma C_{C_6} = 1.$  

(D4)

(#5) $Z_2[0,0]|B\rangle$: $g_\sigma = r_0$, $g_{C_6} = r_3$, $\eta_\sigma T = \eta_\sigma C_{C_6} = -1.$

(D5)

(#15) $Z_2[0,0]|C\rangle$: $g_\sigma = r_3$, $g_{C_6} = r_0$, $\eta_\sigma T = \eta_\sigma C_{C_6} = 1.$

(D6)

(#13) $Z_2[0,0]|D\rangle$: $g_\sigma = g_{C_6} = r_3$, $\eta_\sigma T = \eta_{C_6} = -1.$

(D7)

It turns out these four $Z_2$ SLs around uniform RVB state are all gapped as shown in TABLE II.